

Plasma-flow interfaces for instability control

Srikar Yadala Venkata

▶ To cite this version:

Srikar Yadala Venkata. Plasma-flow interfaces for instability control. Fluid mechanics [physics.class-ph]. Université de Poitiers, 2020. English. NNT: 2020POIT2292 . tel-03162436

HAL Id: tel-03162436 https://theses.hal.science/tel-03162436

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THÈSE

Pour l'obtention du grade de

DOCTEUR DE L'UNIVERSITÉ DE POITIERS

Faculté des Sciences Fondamentales et Appliquées (Diplôme National - Arrêté du 25 mai 2016)

École Doctorale : Sciences et Ingénierie en Matériaux, Mécanique, Energétique

Domaine de recherche : Mécanique des Fluides

Présentée par

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Interface plasma-fluide pour le contrôle d'instabilité Plasma-flow interfaces for instability control

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Thèse soutenue le 17 Décembre 2020 devant la Commission d'Examen

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FUNDING ORGANISATION

This research is funded by the French Government program *Investissements d'Avenir* (Labex INTERACTIFS, reference ANR-11-LABX-0017-01) and the author is grateful for their support during the course of this thesis.

This research is part of the Labex INTERACTIFS project carried out by Université de Poitiers in partnership with ISAE-ENSMA and CNRS. It is a multidisciplinary research project combining mechanics, materials and energy. It is built around Institut PPrime's fields of expertise and aims to promote research actions through a network of international collaborations. In keeping with this theme, part of the current thesis was carried out at the AWEP Department, Faculty of Aerospace Engineering, Delft University of Technology, The Netherlands and the second part was conducted at Institut PPrime, Université de Poitiers, France.

RÉSUMÉ

La recherche présentée dans cette thèse se concentre sur la conception et l'utilisation d'actionneurs plasma à décharge à barrière diélectrique (DBD) de faible épaisseur et à géométrie complexe afin d'exercer un contrôle d'instabilité sur deux configurations d'écoulement dont la dynamique est régie par des mécanismes d'instabilités primaires et/ou secondaires.

Le cas d'une couche limite tridimensionnelle telle que rencontrée sur une aile en flèche est étudié à l'aide de deux stratégies de forçage permettant de manipuler la transition induite par un phénomène d'instabilité stationnaire. Ici, un réseau d'éléments de rugosité discrets (DRE) est installé en amont du forçage par DBD afin de verrouiller l'origine et l'évolution des tourbillons stationnaires transversaux de la couche limite. La première approche de forçage consiste à modifier l'écoulement amont par déformation (UFD). Une seconde approche par modification directe de l'écoulement de base est également introduite (BFM). Un retard de transition est observé indépendamment du forcage réalisé. Cependant, comme les tourbillons transverses sont fortement amplifiés en raison de l'utilisation de DRE, l'action par approche UFD peut conduire à la fois à une atténuation directe des structures fluidiques transverses telle qu'envisagée mais aussi à une action non intentionnelle sur la nature inflectionnelle de l'écoulement de base. La méthode BFM résulte en une interaction directe sur les tourbillons transverses, interactions confirmées par une étude théorique de l'instabilité sous l'effet d'un modèle simplifié d'actionneur DBD. Il s'agit de la première démonstration expérimentale du retard de transition sur une aile en flèche grâce à l'effet d'un actionneur plasma et également à la première preuve de concept expérimentale de la stratégie BFM.

Le sillage d'une couche de mélange plane à bord épais et les phénomènes d'instabilité primaire et secondaire responsables pour l'expansion spatio-temporelle du sillage sont également étudiés. Des conditions de forçage fréquentiel puis spatial sont successivement testées et analysées par approche spectrale (décomposition orthogonale spectrale, SPOD) sur des données expérimentales de PIV multi-champs résolues en temps. L'instabilité primaire est excitée par un forçage spatialement homogène pulsé à la fréquence naturellement la plus amplifiée. Il est montré que la composante moyenne de l'écoulement n'est pas modifiée tandis que le contenu spectral de la couche de mélange est largement affecté. Ce forçage entraîne notamment l'inhibition de l'appariement des structures tourbillonnaires en raison de l'atténuation des instabilités aux sous-harmoniques. A l'inverse, le forçage direct des instabilités aux sous-harmoniques résulte en un renforcement des phénomènes d'appariement conduisant à un fort taux d'épanouissement de la couche de mélange. Enfin, l'effet d'un forçage modulé spatialement se traduit par un taux d'accroissement variant selon la position transverse et qui traduit à la fois le renforcement et la modulation spatiale des structures à grande échelle. La segmentation du forçage selon l'envergure de la couche de mélange permet toujours de modifier les structures transverses mais en sus, la coalescence des structures longitudinales et transversales est favorisée.

Les travaux de recherche réalisés confirment la capacité des actionneurs plasma de type DBD à exercer un forçage modulé à la fois temporellement et spatialement. Les actionneurs proposés ne permettent qu'un contrôle partiel des phénomènes d'instabilités dans une couche limite tridimensionnelle tandis que la forte réceptivité de la région initiale d'une couche de mélange a conduit à des résultats significatifs à la fois sur la dynamique des structures cohérentes transverses et longitudinales. Grace à une large réduction de la puissance électrique consommée dans le cas d'un forçage modulé spatialement, l'efficacité du système de contrôle s'en trouve largement amélioré.

SUMMARY

The research presented in this thesis focuses on the design and use of dielectric barrier discharge (DBD) plasma actuators with thin and complex geometry electrodes to exert instability control on two flow configurations whose dynamics are governed by primary and/or secondary instability mechanisms.

The case of a three-dimensional boundary layer as encountered on a swept wing is studied using two forcing strategies to manipulate the transition induced by a stationary instability phenomenon. Here, an array of discrete roughness elements (DRE) is installed upstream of the DBD forcing in order to lock the origin and evolution of the stationary cross-flow (CF) vortices in the boundary layer. The first forcing approach is upstream flow deformation (UFD). The second approach based on direct modification of the base flow is also introduced (BFM). Independent of the forcing applied, a transition delay is observed. However, as the CF vortices are strongly amplified due to the use of DRE, the action by UFD approach can lead both to a direct attenuation of the CF vortices as envisaged but also to an unintentional action on the inflectional nature of the base flow. The BFM method results in a direct attenuation of the CF velocity component, which is also confirmed by a theoretical study of instability under the effect of the DBD actuator through a simplified model. This is not only the first experimental demonstration of transition delay on a swept wing using plasma actuators, but also the first experimental proof of concept of the BFM strategy.

The wake of a plane mixed layer with a thick edge and the primary and secondary instability phenomena responsible for the spatio-temporal expansion of the wake are also studied. Frequency and then spatial forcing conditions are successively tested and analysed by spectral approach (spectral proper orthogonal decomposition, SPOD) on experimental data from multifield time-resolved particle image velocimetry. The primary instability is excited by a spatially uniform forcing pulsed at the naturally most amplified frequency. It is shown that the mean component of the flow is not modified while the spectral content of the mixing layer is largely affected. This forcing leads, in particular, to the inhibition of the pairing of vortical structures due to the attenuation of sub-harmonic instabilities. Conversely, direct forcing of sub-harmonic instabilities results in a reinforcement of the pairing phenomena, leading to a higher growth rate of the mixing layer. Finally, spatially modulated forcing results in a growth that varies according to the spanwise position, which reflects both the reinforcement and the spatial modulation of large-scale spanwise structures. The modulation of the forcing according to the scale of the mixing layer always allows the modification of the spanwise structures but in addition, the coalescence of the streamwise and spanwise structures is favoured.

The research work carried out confirms the ability of DBD plasma actuators to exert a forcing modulated both temporally and spatially. The proposed actuators allow only a partial control of the instability phenomena in the three-dimensional boundary layer while the high receptivity of the initial region of a mixing layer has led to significant results on the dynamics of both spanwise and streamwise vortices. Thanks to a large reduction of the power consumed in the case of spatially modulated forcing, the efficiency of the control system is greatly improved.

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ABBREVIATIONS & NOMENCLATURE

Abbreviations

DBD	Dielectric Barrier Discharge
AC	Alternating Current
EFD	Electro-Fluid-Dynamic
PMMA	Poly(methyl methacrylate)
PET	Poly-Ethylene Terephthalate
EE	Exposed electrode
CE	Covered electrode
ES	Ensemble Size
PSD	Power spectral Density
LST	Linear Stability Theory
SPOD	Spectral Proper Orthogonal Decomposition
TKE	Turbulent Kinetic Energy
NS	Navier-Stokes
OS	Orr-Sommerfeld
IR	Infra-Red
PIV	Particle Image Velocimetry
FOV	Field of View
HWA	Hot-Wire Anemometry
\mathbf{CF}	Cross-Flow
ISL	Inviscid Stream Line
DRE	Discrete Roughness Elements
UFD	Upstream Flow Deformation
BFM	Base-Flow Modification
UML	Unforced Mixing Layer
KH	Kelvin-Helmholtz

Mathematical entities

	Time-averaged quantity
\dots' or $\tilde{\dots}$	Fluctuating quantity
$\cdots rms$	Root-mean-square
σ	Standard deviation
$\cdots max$	Maximum value
$\cdots N$	Normalized quantity
$\cdots R$	Reconstructed quantity
∞	Free-stream quantity
Δ	Increment or resolution
*	Complex conjugate
$\cdots r$	Real part of complex number
····i	Imaginary part of complex number
ϵ	Measurement error

t	Time
f	Frequency
Re	Reynold's number
p	Pressure
Tu	Turbulence intensity
ρ	Density of air
ν	Kinematic viscosity
c_{μ}	Momentum coefficient
<i>f</i> #	Aperture opening (focal-STOP) of camera lens
I_1, I_2	Light intensity distribution in consecutive PIV frames
C	Cross-correlation operator to compute vectorial data from PIV acquisition
DBD p	lasma actuators
V_{ac}	AC voltage amplitude
f_{ac}	AC carrier frequency
T_{ac}	Time-period of AC waveform $(1/f_{ac})$
f_b	Burst-modulation frequency
V_0	Plasma onset voltage amplitude
Ι	Discharge current
R_S	Resistance of shunt resistor
V_S	Voltage drop across shunt resistor
C_m	Capacitance of monitor capacitor
V_m	Voltage across monitor capacitor
Q_m	Charge in monitor capacitor
P_{avg}	Average power consumed per AC cycle
A_P	Fitting parameter for power-voltage curve
x_p	Axis along the induced flow/wall-parallel direction
y_p	Axis perpendicular to the dielectric surface
z_p	Axis along the span of the actuator
w_E	Air-exposed electrode width
w_C	Covered/encapsulated electrode width
l_S, w_S	Length and width of covered-electrode stems
λ_z	Distance between adjacent stems along the spanwise direction
0	Electrode overlap
z_{PA}	Spanwise extent of DBD plasma actuator
x_{PA}	Distance of plasma-generating edge from dielectric foil edge
l_p	Plasma discharge length along the span
l_{CE}	Spanwise-length of covered electrode that results in a plasma discharge
	$(= w_S \times \text{ number of stems})$
$(uvw)_p$	Velocity components along $(xyz)_p$ within the induced flow
U_p	Induced velocity magnitude
δ_p	Location of $u_{p,max}$ from dielectric surface
f'	Differential of induced mass-flux (or velocity distribution along y_p)
F_{EMF}	External momentum flux
r_x	LFD body force along the wall-parallel direction (x_p)
I_x	Infust-generated along the wall-parallel direction (x_p)
τ_w	Incurred skin-iriction Contribution of EED hody force (E) to the threat respect of
I_{x}^{-}	Continuation of EFD body force (F_x) to the thrust generated

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Common variables

Linear stability theory

- ψ Perturbation eigenfunction
- α Complex streamwise wavenumber
- α_i Amplification rate
- β Complex spanwise wavenumber
- $\omega \qquad \qquad \text{Angular frequency} \qquad \qquad$
- N Amplification N factor

x Axis normal to the swept-wing leading edge y Axis orthogonal to the chord plane z Axis parallel to the leading edge $(xyz)_{ISL}$ Local coordinate system along the ISL x_{tr} Transition location c Streamwise chord of the swept-wing model c_x Chord length of swept-wing model perpendicular to leading edge α Angle of attack uvw Velocity components along the xyz axes $(uvw)_{ISL}$ Velocity components along the ISL u_e Boundary layer edge velocity θ_e Boundary layer momentum thickness δ^* Boundary layer displacement thickness U_∞ Free-stream velocity M_∞ Free-stream Mach number Re_c Chord Reynold's number CP Coefficient of pressure λ_z Spanwise wavelength of stationary CF instability $\lambda_z,crit$ Spanwise wavelength of DRE C_f Skin-friction coefficient $C_f,lumbulent$ Skin-friction coefficient of laminar and turbulent boundary layers	Swept-wing boundary	layer
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$\begin{array}{llllllllllllllllllllllllllllllllllll$	Re_c	Chord Reynold's number
$\begin{array}{ll} \lambda_z & & \text{Spanwise wavelength of stationary CF instability} \\ \lambda_{z,crit} & & \text{Spanwise wavelength of critical/most unstable stationary} \\ CF instability \\ d_{DRE}, k_{DRE} & & \text{Diameter and height of DRE} \\ C_f & & \text{Skin-friction coefficient} \\ C_{f,laminar}, C_{f,turbulent} & & \text{Skin-friction coefficient of laminar and turbulent boundary layers} \end{array}$	C_P	Coefficient of pressure
$\begin{array}{lll} \lambda_{z,crit} & & \text{Spanwise wavelength of critical/most unstable stationary} \\ CF & \text{instability} \\ d_{DRE}, k_{DRE} & & \text{Diameter and height of DRE} \\ C_f & & \text{Skin-friction coefficient} \\ C_{f,laminar}, C_{f,turbulent} & & \text{Skin-friction coefficient of laminar and turbulent boundary layers} \\ \end{array}$	λ_z	Spanwise wavelength of stationary CF instability
CF instability d_{DRE}, k_{DRE} Diameter and height of DRE C_f Skin-friction coefficient $C_{f,laminar}, C_{f,turbulent}$ Skin-friction coefficient of laminar and turbulent boundary layers	$\lambda_{z,crit}$	Spanwise wavelength of critical/most unstable stationary
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C_f Skin-friction coefficient $C_{f,laminar}, C_{f,turbulent}$ Skin-friction coefficient of laminar and turbulent boundary layers	d_{DRE}, k_{DRE}	Diameter and height of DRE
$C_{f,laminar}, C_{f,turbulent}$ Skin-friction coefficient of laminar and turbulent boundary layers	C_f	Skin-friction coefficient
$\alpha + 1 $ α :	$C_{f,laminar}, C_{f,turbulent}$	Skin-friction coefficient of laminar and turbulent boundary layers
η Control efficiency	η	Control efficiency

Plane mixing layer		
x	Axis along the streamwise direction	
y	Axis orthogonal to the streamwise direction	
z	Axis parallel to the splitter-plate trailing edge	
uvw	Velocity components along the xyz axes	
U_1	Streamwise velocity of low-velocity (bottom) stream	
U_2	Streamwise velocity of high-velocity (top) stream	
ΔU	Velocity difference between the two streams $(= U_2 - U_1)$	
\overline{U}	Average velocity of two streams $(=(U_2+U_1)/2)$	
u_c	Large-scale vortex convection velocity	
r	Velocity ratio $(=U_1/U_2)$	
R	Velocity ratio $(=\Delta U/(2\overline{U}))$	
δ_{99}	Boundary layer thickness based on 99% threshold	
H_1, H_2	Shape factor of low and high-velocity side splitter-plate boundary layers	
$ heta_i$	Momentum thickness of high-velocity boundary layer at the splitter-plate	
	trailing edge	
heta	Momentum thickness of mixing layer	
δ_w	Vorticity thickness of mixing layer	
$y_{0.05}, y_{0.5}, y_{0.95}$	y-location where $u_N = 0.05$, $u_N = 0.5$ and $u_N = 0.95$ respectively	
C	Growth rate fitting parameter	
R_{uu}	Streamwise component of TKE $(=\overline{u'u'})$	
R_{vv}	Cross-stream/transverse component of TKE $(=\overline{v'v'})$	
R_{ru}	Revnold's shear stress $(=-\overline{u'v'})$	
λ_r	Streamwise wavelength	
λ_z	Spanwise wavelength	
\tilde{St}_{θ}	Strouhal number $(= f\theta_i/\overline{U})$	
St_0, f_0, λ_0	Strouhal number, frequency and wavelength of fundamental KH instability	
St _f	Strouhal number of forced primary (KH) instability	
$P_{ava,f}$	Average power consumed by DBD actuator when operated with a	
	burst-modulation cycle	
$T_{x,f}$	Thrust generated by DBD actuator when operated with burst-modulation cycle	
Φ^{γ}	PSD of fluctuating quantity γ	
ϕ	Phase along burst-modulation cycle	
v_{ϕ}	Cross-stream velocity at a particular phase	
ω	Vorticity	
S	Shear strain rate	
Q	Second invariant of velocity field for vortex identification	
SPOD		
\widetilde{Q}	Stochastic variable representing fluctuations	
Ω	Spatial domain	
S	Cross-spectral tensor for SPOD computation	
W	Positive-definitive Hermitian weight function	
N_{fr}	Number of vector fields in each block or window size	
N_k	Number of blocks	
k	Eigenmode number	
Λ	Eigenvalue	
c_k	SPOD mode convergence	
Ψ	Eigenfunction	
Θ	Eigenvector	
A	Mode amplitude	

RÉSUMÉ ÉTENDU

1. INTRODUCTION

La discipline du contrôle d'écoulement a suscité une attention considérable de la part des chercheurs et des ingénieurs, en tant que moyen d'améliorer l'efficacité de nombreux systèmes fluide-mécaniques. Le plus souvent, le contrôle d'écoulement consiste à apporter de petites modifications à une configuration existante, par l'utilisation de certains dispositifs, afin de réaliser un grand avantage technique. En fonction du résultat souhaité, le contrôle de l'écoulement peut être utilisé pour éviter/forcer la séparation de la couche limite dans les écoulements proche paroi, pour retarder/avancer la transition du régime laminaire-turbulente ou pour diminuer/augmenter la turbulence. Elle peut effectivement conduire à la réduction de la traînée, à l'amélioration de la portance, à l'amélioration du mélange ou à la suppression du bruit induit par l'écoulement. Les différentes méthodes et dispositifs utilisés dans le passé dans le domaine du contrôle d'écoulemnet ont été synthétisés dans un certain nombre de travaux de synthèse (Gad-el-Hak et al., 1993; Joslin, 1998; Gad-el-Hak, 2007; Cattafesta III & Sheplak, 2010).

Le contrôle d'écoulement peut être grossièrement classé en deux catégories : passif et actif. Le contrôle passif comprend des dispositifs tels que les éléments de rugosité, les générateurs de tourbillons, les turbulateurs, etc. qui sont des installations permanentes à la surface d'un corps aérodynamique. En tant que tels, ces dispositifs utilisent une approche de "forçage brute" stationnaire qui ne nécessitent aucune source d'énergie auxiliaire ou dépense énergétique. Cependant, un inconvénient majeur des dispositifs passifs est qu'ils sont toujours opérationnels. En outre, étant donné que les méthodes passives impliquent presque toujours des modifications géométriques, le contrôle exercé ne peut pas être régulé en fonction des conditions d'écoulement et peut donc potentiellement avoir un effet préjudiciable sur les effets recherchés. En revanche, le contrôle d'écoulement dit actif implique l'utilisation d'actionneurs qui injectent de l'énergie dans le flux à contrôler et nécessitent donc une certaine dépense d'énergie pour les faire fonctionner. Au minimum, ces actionneurs peuvent être mis en marche en fonction des besoins et le plus souvent certains paramètres de fonctionnement peuvent être modifiés par une commande électrique. Ce type de contrôle actif peut à son tour être classé en deux sous-catégories : le contrôle en boucle fermée et en boucle ouverte. Dans le contrôle en boucle fermée, l'actionneur utilisé fait partie d'une boucle de rétroaction qui comprend également certains capteurs placés dans le fluide. La boucle de contrôle devient ainsi un système dynamique. Dans la commande en boucle ouverte, il n'y a pas de signal de retour et l'actionneur est commandé par le réglage manuel de ses paramètres de fonctionnement.

Par le passé, un certain nombre de dispositifs ont été utilisés afin de réaliser des expériences de contrôle actif d'écoulement. Ainsi, des réchauffeurs de surface peuvent être utilisés pour influencer le flux en faisant varier localement la viscosité ou en favorisant les gradient de densités (voir Liepmann, Brown, et al., 1982). Les actionneurs fluidiques qui permettent une injection (soufflage) ou une aspiration de fluides peuvent avoir une influence significative sur le champ d'écoulement en modifiant le profil de la vitesse moyenne de l'écoulement. De même, il existe

des actionneurs à débit nul (jet synthétique) qui ingèrent et expulsent le fluide consécutivement (Glezer & Amitay, 2002). Les ondes acoustiques peuvent être aussi utilisées pour influencer la dynamique des couches de cisaillement. Des rubans vibrants, des volets oscillants ou des actionneurs piézoélectriques peuvent également être utilisés pour introduire certaines fluctuations périodiques dans le champ d'écoulement. Bien que ces actionneurs (en particulier les dispositifs mécaniques) aient démontré leur efficacité dans diverses configurations d'écoulement, ils présentent des inconvénients inhérents. Ces actionneurs sont composés de pièces mobiles qui s'usent et se déchirent et qui sont donc susceptibles de tomber en panne. Leur installation sur des corps aérodynamiques est souvent compliquée par la nécessité de modifications géométriques afin d'assurer le logement des pièces motrices. Enfin, ils ajoutent du poids et peuvent être une source de bruit.

Dans cette thèse le système de contrôle retenu est de type électrofluidodynamique. Il s'agit d'un dispositif qui ne présente pas les inconvénients listés précédemment. Ici, des actionneurs plasma à décharge à barrière diélectrique (DBD) sont utilisés. La plupart des avantages des dispositifs basés sur la technologie DBD découlent de la nature électrique du phénomène de transfert d'énergie par phénomène electro-fluid-dynamique (EFD) et par conséquent sur l'absence de pièces mobiles. En effet, ces actionneurs sont simplement constitués d'électrodes conductrices fines et d'une barrière diélectrique de quelques millimètres. Les actionneurs de type DBD exercent un forçage fortement instationnaire dont la composante moyenne est facilement pilotée électriquement et dont les fluctuations périodiques sont imposées par le signal de commande de type alternatif (le plus souvent un signal sinusoïdal). Enfin, étant commandés électriquement, l'inclusion de ces dispositifs dans des algorithmes de contrôle en boucle fermée est également possible. Un certain nombre d'articles de synthèse publiés au cours des deux dernières décennies ont fait le bilan non exhaustif de différentes applications pouvant être visées par ce type d'actionneurs (Moreau, 2007; Corke, Enloe, et al., 2010; Wang et al., 2013; Kriegseis, Simon, et al., 2016; Corke & Thomas, 2018; K. Choi & Kim, 2018). Comme indiqué dans ces articles, les actionneurs plasma à base de DBD ont démontré leur capacité à réaliser un contrôle actif à plusieurs reprises mais toujours dans des conditions sub-soniques à faible vitesse.

Il est désormais établi que l'influence des actionneurs plasma se limite à des écoulements de Reynolds faibles, en tout cas non représentatifs des nombre de Reynolds tels que rencontrés dans les applications ciblées. La raison principale est la faible quantité de mouvement créée par ces actionneurs. Le principe de base de ces actionneurs est le transfert de quantité de mouvement par collisions entre particules neutres et particules chargées en mouvement sous l'influence d'un fort champs électrique. La déformation moyenne des profils de vitesse dans la zone d'interaction entre le plasma et le flux d'air est souvent insuffisante pour modifier l'écoulement dans sa globalité. De plus, dans ce type de contrôle par forçage brute, le rendement électromécanique est très faible et l'obtention d'un gain net positif est par conséquent des plus hypothétiques. Ici, les configurations d'études présentent des conditions d'écoulement favorables au développement d'instabilités hydrodynamique. En effet, dans les configurations d'écoulement présentant un profil de vitesse de type inflexionel, certaines perturbations ondulatoires de petite amplitude peuvent produire des effets importants comme l'a démontré Lord Rayleigh (1880). Ces perturbations de petite amplitude sont ce qu'on appelle des instabilités d'écoulement. Elles peuvent être des ondes dépendantes du temps/de la fréquence ou des ondes se produisant dans l'espace. Ces ondes, dans certaines bandes (de fréquence ou de longueur d'onde spatiale) dépendantes des conditions de l'écoulement, peuvent s'amplifier rapidement et perturber la stabilité hydrodynamique du fluide, provoquant ainsi un effet important. L'amplification de ces ondes d'instabilité peut être contrôlée avec un forçage de très faible amplitude qui affecte les caractéristiques de stabilité du champ d'écoulement, produisant ainsi le résultat nécessaire. Cette approche du contrôle de l'écoulement est en fait plus efficace qu'une approche de "force brute" car la dynamique de l'écoulement peut être efficacement contrôlée avec une dépense énergétique moindre. De plus, de part un contrôle uniquement basé sur des interactions pariétales ou les vitesses rencontrées restent modérées, il est attendu que cette approche de contrôle d'instabilité se montre efficace sur une large gamme de Reynolds.

1.1. Motivations et objectifs de la thèse

La motivation principale de cette thèse est d'étudier les capacités et les limites des actionneurs de plasma DBD dans la réalisation d'un contrôle actif de phénomène d'instabilité. Par nature, ces actionneurs exercent un forçage instationnaire fortement périodique. Grâce aux caractéristiques géométriques de leur construction, un forçage tridimensionnel modulé dans l'espace peut également être réalisé. Différents actionneurs à technologie DBD pouvant imposer un forçage de nature bidimensionnelle et tridimensionnelle seront conçus et fabriqués. Après une caractérisation électro-mécanique approfondie de ces actionneurs, ils seront déployés à des fins de contrôle de l'écoulement sur deux configurations distinctes. Ainsi, le cas du contrôle de l'écoulement sur une aile en flèche sera abordé avec pour objectif de manipuler la région de transition laminaire/turbulent. Ce type d'écoulement, et ses instabilités transverses intrinsèques, étant particulièrement sensibles à l'état de surface du modèle, des actionneurs spécifiques de faible intrusivité et imposant un forçage spatialement modulé ont du être développés. Dans un second temps, la technique de contrôle a été étendue au cas d'une couche de mélange bidimensionnelle dont la dynamique est principalement pilotée par des instabilités primaires mais nous verrons que des instabilités de type secondaires contribuent également au taux d'épanouissement de la couche de mélange. L'objectif ici est de commander la dynamique tourbillonnaire du sillage d'une couche de mélange à bord épais en forçant d'abord les instabilités primaires, puis en imposant un forçage modulé spatialement afin de venir cette fois modifier les instabilités de type secondaires.

Les objectifs spécifiques de cette thèse dont l'approche est purement théorique et expérimentale sont détaillés ci-dessous.

1.1.1. Actionneurs plasma

Le cœur de cette thèse étant construit sur un contrôle par actionneurs plasma DBD, certaines tâches sont nécessaires afin d'assurer une mise en œuvre efficace lors de leur déploiement sur les différentes configurations d'études qui seront réalisées par la suite. Deux types d'actionneurs vont etre nécessaires. Le premier est supposé imposer un forçage uniforme le long du bord de l'électrode au contact avec l'air. Il s'agit d'une configuration très classique, largement utilisés dans la communauté. Le second type d'actionneur est largement plus innovant puisqu'il doit permettre un forçage tridimensionnel modulé dans l'espace. Les taches principales à réaliser sont les suivantes:

Les actionneurs déployés doivent respecter des contraintes strictes d'épaisseur et de rugosité afin d'éviter que la simple présence de l'actionneur n'ait un effet déstabilisant sur l'écoulement. Il s'agit d'une préoccupation générale liée à l'utilisation de ces actionneurs pour des applications de contrôle d'instabilité se développant aux interfaces fluides/solides. Pour contourner ce problème, il est nécessaire de formuler une technique de fabrication qui produise de tels actionneurs. En outre, la technique proposée doit être capable de produire
des électrodes présentant des caractéristiques géométriques complexes.

• La deuxième tâche est liée aux actionneurs modulés spatialement. Tout d'abord, une preuve de concept est nécessaire pour démontrer la tridimensionnalité de l'écoulement induit. Deuxièmement, il est nécessaire de quantifier les performances de ces actionneurs modulés en espace.

1.1.2. Contrôle d'écoulement sur une aile en flèche

Cette thèse vise à contrôler le mécanisme d'instabilité dite transversale (CF) qui est la cause de la transition laminaire-turbulente dans les couches limites tridimensionnelles se développant sur les ailes en flèche. Une stratégie de contrôle qui a récemment attiré l'attention est l'utilisation d'éléments de rugosité discrets (DRE) positionnés au bord d'attaque afin de provoquer une déformation favorable de l'écoulement moyen le long de l'envergure de l'aile (Saric, Carrillo, et al., 1998). Ceci a ensuite été généralisé sous le nom de déformation de l'écoulement en amont (UFD) pour englober l'utilisation de tout dispositif pouvant provoquer cette déformation de l'écoulement moyen (Wassermann & Kloker, 2002). Bien que cette stratégie ait donné des résultats positifs dans le cadre de quelques efforts expérimentaux et numériques, elle n'est pas très robuste en raison de sa dépendance à l'égard d'un certain nombre de paramètres. Une autre stratégie, principalement issue de simulation numérique (Messing & Kloker, 2010; Dörr & Kloker, 2015b), est basée sur la manipulation de la cause de l'instabilité, la composante de vitesse transverse dans la couche limite, et est appelée modification de l'écoulement de base (BFM). A ce jour, cette stratégie n'a jamais été validée expérimentalement.

Dans ces travaux, les stratégies de contrôle de type UFD et BFM seront réalisées grâce à des actionneurs plasma. Le but recherché ici est la réalisation d'un retard dans la transition laminaire-turbulente. Les différentes stratégies de contrôle seront étudiées dans des conditions d'écoulement identiques afin de mieux comparer l'autorité et l'efficacité de contrôle obtenues avec ces deux stratégies. En cas de succès, cela marquera le premier effort expérimental pour obtenir un retard de la transition dominée par l'instabilité de la composante transverse de la vitesse en paroi. A ce titre, quatre objectifs sont identifiés:

- Le premier objectif est de réaliser la stratégie de contrôle de type UFD à l'aide d'actionneurs DBD qui génèrent un forçage modulé spatialement selon l'envergure du modèle. De plus, avec certaines configurations de forçage (conception des actionneurs et conditions de fonctionnement), on peut s'attendre à ce que la force exercée pour réaliser la stratégie UFD effectue simultanément et localement une action de type BFM. L'effet du forçage UFD et la façon dont les deux stratégies contribuent à contrôler la transition n'est pas bien comprise et, à ce jour, n'a pas été explorée de manière expérimentale.
- Le second objectif est de réaliser et d'étudier expérimentalement la stratégie de contrôle de type BFM en utilisant des actionneurs de plasma DBD. Avant l'investigation expérimentale, une étude théorique préliminaire sera réalisée pour étudier la stabilité de la couche limite, sous l'influence d'une force volumique reproduisant les effets de la force de type EHD produite par l'actionneur. Ce modèle numérique simplifié basé sur la théorie de la stabilité linéaire sera comparé aux résultats expérimentaux.
- Enfin, cette thèse ne fournit pas seulement la première preuve de concept expérimentale de la stratégie de contrôle de la BFM, mais constitue également la première comparaison

des deux stratégies de contrôle pour des conditions d'écoulement identiques. Ainsi, elle fournit des conditions favorables pour étudier la collaboration de ces deux stratégies. En outre, ces deux stratégies seront également comparées en termes d'efficacité de contrôle obtenue, notamment en fonction de leur effet sur la traînée de frottement totale.

1.1.3. Contrôle d'écoulement d'une couche de mélange plane

La partie suivante de cette thèse concerne l'application d'actionneurs plasma DBD sur une couche de mélange laminaire plane. Cette configuration est caractérisée par des tourbillons à grande échelle dans le sens de l'envergure qui résultent du mécanisme d'instabilité primaire de type *Kelvin-Helmholtz* (Brown & Roshko, 1974). La dynamique et l'interaction de ces tourbillons ont un impact direct sur la croissance de la couche de mélange et son taux d'épanouissement. Deuxièmement, le mécanisme d'instabilité secondaire conduit à la formation de tourbillons en épingle à cheveux qui s'enroulent autour des tourbillons transversaux. L'interaction de ces deux structures tourbillonnaires entraîne la production de turbulences à petite échelle et un mélange accru à l'échelle moléculaire. La couche de mélange est la représentation la plus simple d'une couche de cisaillement libre qui prévaut dans tout type d'écoulement de fluide. Ainsi, l'analyse de l'effet du forçage des actionneurs DBD sur la couche de mélange revêt une importance significative pour améliorer les applications de contrôle de l'écoulement de ces actionneurs. Un certain nombre d'outils analytiques seront générés et appliqués sur les données expérimentales acquises afin de discerner l'influence de la stratégie de contrôle utilisée. Actuellement, les objectifs suivants sont identifiés:

- Concevoir et mettre en place un dispositif expérimental permettant l'étude d'une couche de mélange résultante de l'interaction de deux flux présentant un gradient de vitesse transverse. Cette étape préliminaire requiert une caractérisation adéquate du champ de vitesse dans le sillage de la plaque séparatrice et une comparaison avec les caractéristiques des moyens d'essais préalablement rapportées dans la littérature.
- Générer et à valider des outils analytiques et à formuler une méthodologie pour étudier cette configuration d'écoulement. La mesure de la vitesse d'écoulement sera obtenue à partir de la vélocimétrie par image de particules résolue temporellement. Cet ensemble de données expérimentales constitue la base de la première analyse spatio-temporelle du champ de vitesse de la couche de mélange. Les outils analytiques identifiés comprennent l'analyse par stabilité linéaire, la densité spectrale de puissance, la décomposition orthogonale en modes propres appliquée au domaine spectral (SPOD) pour obtenir l'organisation spatiale des fluctuations et les techniques d'identification des tourbillons afin d'extraire la dynamique et les interactions des structures tourbillonnaires.
- Inspecter l'impact d'un forçage spatialement uniforme sur l'instabilité primaire. Il s'agit ici de tirer parti de la nature instationnaire du forçage appliqué par les actionneurs DBD. En outre, l'impact des différents paramètres de fonctionnement de l'actionneur DBD sera également étudié. En particulier, les effets de forçage aux harmoniques et sousharmoniques de l'instabilité naturellement le plus amplifiée sont déterminés.
- Réaliser une étude préliminaire de manipulation des instabilités secondaires dans une couche de mélange plane. Ceci est réalisé en appliquant un forçage modulé spatialement selon l'envergure de la plaque séparatrice. Il s'agit ici d'étudier l'influence d'un forçage tridimensionnelle et instationnaire sur le développement des structures fluidiques primaires

(transverses) et secondaires (longitudinales).

1.2. Organisation du manuscript

Cette thèse est organisée comme suit. Le chapitre 2 contient un bref descriptif des efforts de recherche réalisés dans le passé pour mieux comprendre les différents concepts étudiés dans la présente thèse. Il comprend une brève description du principe de fonctionnement des actionneurs plasma DBD. La description de l'évolution de la transition laminaire-turbulente dominée par la mécanique de l'instabilité hydrodynamique dans une couche limite sur une aile en flèche. Les différentes stratégies de contrôle conçues pour maîtriser cette instabilité afin de réaliser un retard dans la transition seront ensuite présentées. Les différents mécanismes d'instabilité qui définissent la dynamique des structures d'écoulement tourbillonnaire dans la configuration d'écoulement de la couche de mélange plane sont également abordés.

La méthodologie suivie et les outils utilisés dans cette thèse pour étudier les configurations d'écoulement correspondantes sont introduits dans le chapitre 3. Ce chapitre énumère les différentes installations d'essai expérimentales utilisées au cours de cette thèse, les systèmes de mesure déployés et enfin les différents outils utilisés afin d'analyser les données expérimentales.

Le chapitre 4 est consacré aux actionneurs plasma DBD qui seront utilisés pour le contrôle de l'écoulement et de l'instabilité dans la couche limite d'une aile en flèche et les configurations d'écoulement de la couche de mélange plane. Une nouvelle technique de fabrication pour réaliser des actionneurs DBD à géométrie complexe présentant une épaisseur d'électrode à l'échelle micrométrique est proposée et les actionneurs fabriqués sont validés. Les configurations géométriques des différents actionneurs à plasma DBD qui seront utilisés dans cette thèse sont ensuite décrites. La géométrie des actionneurs modulés spatialement destinés à appliquer un forçage tridimensionnel est ensuite caractérisée de manière approfondie à la fois électriquement et mécaniquement par des mesures de PIV tomographique.

Les efforts déployés pour étudier la couche limite sur une aile en flèche sont présentés dans le chapitre 5. Les conditions d'écoulement étudiées sont d'abord présentées, puis les résultats de l'analyse de la stabilité de la couche limite sont utilisés pour la prédiction des régions de transition. L'état de la couche limite de référence (non forcée) est ensuite présenté. Les résultats obtenus lors de la réalisation de différentes stratégies de contrôle (UFM et BFM) sont successivement présentés. Enfin, une estimation de l'efficacité du contrôle sur la base de son impact sur la traînée de frottement totale est comparée pour les deux approches de contrôle retenues dans cette thèse.

Le chapitre 6 contient tous les résultats acquis dans le cadre des recherches effectuées sur l'écoulement de type couche de mélange plane. Dans un premier temps, une caractérisation approfondie de la couche de mélange non forcée a été entreprise. Les résultats obtenus sont analysés et comparés aux efforts expérimentaux précédents. Ensuite, les deux différents actionneurs de plasma DBD sont utilisés pour manipuler les mécanismes d'instabilité sous-jacents et analyser leur impact sur le champ d'écoulement.

Enfin, une synthèse des essais réalisés et des résultats principaux est proposée en Chapitre 7. Au-delà de simplement conclure sur les réalisations expérimentales conduites lors de cette thèse, des éléments de perspective sont également proposés dans l'optique d'affiner les résultats d'ors et déjà obtenus.



Figure 2.: (a) Schematic of a typical DBD plasma actuator, reproduced from Kotsonis et al. (2011). (b) Photograph of the plasma discharge, reproduced from Forte et al. (2006).

2. REVUE DE LA LITTÉRATURE

Ce second chapitre est consacré à la présentation de quelques principes et concepts fondamentaux concernant trois aspects principaux de cette thèse, ceci en passant en revue les résultats des recherches passées sur ces sujets. Ici le chapitre est présenté en version fortement condensée et se concentre sur une présentation rapide des actionneurs plasma DBD, une description des instabilités transverses rencontrées dans la couche limite d'une aile en flèche et de leur conséquence sur le phénomène de transition laminaire/turbulent, et enfin un rappel des mécanismes d'instabilité qui régissent l'évolution et la dynamique des structures cohérentes dans une couche de mélange plane.

2.1. Actionneurs à décharge à barrière diélectrique

La faisabilité de générer un écoulement de type jet de paroi mince par l'utilisation d'une décharge de surface de type décharge à barrière diélectrique et d'utiliser ce jet local dans des applications de contrôle découlement a été démontré pour la première fois par Roth et al. (1998). La conception initialement proposée par Roth et al reste la plus souvent utilisée. L'actionneur plasma DBD est alors composé de deux électrodes métalliques qui sont montées sur les deux côtés d'un substrat diélectrique (voir figure 2a). Les diélectriques généralement utilisés sont le PMMA, le Teflon, le verre, la céramique ou les polyamides tels que le Kapton par exemple, dont l'épaisseur varie entre quelques dizaines de micromètres et quelques millimètres. La largeur des électrodes est de quelques millimètres et elles sont montées de façon asymétrique de part et d'autre du diélectrique sans ou avec chevauchement. Ce dispositif est communément appelé actionneur plasma à décharge à barrière diélectrique (DBD). Le principe de fonctionnement est basé sur l'ionisation de l'air et de la mise en mouvement des particules chargées. En appliquant un potentiel électrique suffisamment élevé à l'électrode haute-tension il est possible d'ioniser l'air (voir Figure 2b) dans la région située au-dessus de la seconde électrode qui elle est connectée à un potentiel inférieur (généralement connectée à la masse). La rupture électrique du milieu génère alors des ions et des électrons libres. Ces ions se déplacent de par le champ électrique produit et elles entrent en collision avec des particules d'air neutres à proximité, transférant ainsi une partie de leur énergie cinétique et générant une force volumique de Coulomb. Cette force



Figure 3.: (a) Temporal evolution of the different NS terms governing the development of the induced flow in the wall-parallel direction, reproduced from Kotsonis et al. (2011). (b) Wall-parallel component of the mean EFD body forced distribution (F_x in N m⁻³), reproduced from Benard, Debien, et al. (2013).

volumique reste localisée et de faible amplitude (figure 3) mais elle est suffisante pour donner lieu à un jet parallèle à la paroi diélectrique, dont la vitesse est généralement inférieure à 10 m s⁻¹. Afin d'éviter une accumulation de charges à la surface du diélectrique qui entraînerait une réduction du champ électrique local, et qui pourrait finalement conduire à l'extinction de l'ionisation, un signal de courant alternatif (AC) est imposé en entrée. Ainsi, la fréquence (f_{ac}) de la forme d'onde du courant alternatif va de 100 Hz à quelques dizaines de kilohertz et l'amplitude de tension (V_{ac}) peut être de quelques kV à environ 50 kV selon l'épaisseur et les propriétés du substrat diélectrique. Au cours des deux dernières décennies, de nombreux travaux ont été réalisés par différents groupes de recherche sur le fonctionnement de ces actionneurs et leur capacité à produire un écoulement local. Ainsi, les principes de base, le mécanisme de fonctionnement et l'utilisation de ces actionneurs pour différentes applications de contrôle d'écoulement peuvent être trouvés dans un certain nombre d'articles de synthèse publiés au cours des deux dernières décennies (Moreau, 2007; Corke, Enloe, et al., 2010; Wang et al., 2013; Benard & Moreau, 2014; Kotsonis, 2015; Kriegseis, Simon, et al., 2016; Corke & Thomas, 2018; K. Choi & Kim, 2018).

Les actionneurs plasma DBD génèrent une force volumique par ionisation de l'air, comme mentionné précédemment. Le jet résultant est souvent appelé vent ionique ou vent électrique). Les premières mesures de la vitesse du jet induit étaient acquises à l'aide d'un tube de Pitot en verre (sonde de pression totale) relié à un manomètre (voir par exemple Pons et al., 2005; Forte et al., 2007; Jolibois & Moreau, 2009). Ces sondes peuvent être fixées à un système de déplacement pour obtenir la distribution spatiale de la vitesse dans le jet. Bien que cette méthode soit simple, rapide et très précise pour mesurer la vitesse dans le jet induit, elle est intrusive et seule la composante horizontale (parallèle à la paroi) de la vitesse peut être mesurée. De plus, les faibles taux d'échantillonnage de cette mesure interdisent l'acquisition de données résolues dans le temps. Une technique de mesure qui fournit des données de vitesse dans le jet induit avec une excellente résolution temporelle est l'anémométrie à fil chaud. Cependant, cette méthode de quantification de la vitesse induite présente certaines limites inhérentes car les mesures peuvent être affectées par les interférences électromagnétiques dues à l'actionneur. De plus, la génération d'un arc électrique entre le bord de l'électrode exposé à l'air et le fil chaud est une possibilité si la sonde est très proche de l'actionneur et détruirait la sonde et l'anémomètre. En raison de ces limitations, l'utilisation de techniques de mesure optique telles que la vélocimétrie par image de particules (PIV, voir par exemple Kriegseis, Grundmann, et al.,



Figure 4.: (a) Schlieren image of a starting vortex generated by a DBD plasma actuator, reproduced from Moreau et al. (2008). (b) Time-averaged ionic wind generated by a DBD plasma actuator with a 3 mm thick dielectric substrate ($V_{ac} = 22$ kV, $f_{ac} = 1.5$ kHz), reproduced from Debien, Benard, & Moreau (2012).

2011; Kotsonis et al., 2011; Debien, Benard, & Moreau, 2012) ont été employées au cours de la dernière décennie pour étudier la topologie du jet induit. Cette méthode est non intrusive, fournit au moins deux composantes de la vitesse et la mesure du jet induit est obtenue avec des résolutions spatiales et/ou temporelles élevées. Ces mesures ont permis d'en apprendre beaucoup sur la topologie du jet induit par les actionneurs de plasma DBD. L'écoulement est accéléré dès que l'actionneur DBD est opérationnel. Cette accélération dure des dizaines de millisecondes et génère un vortex initial (voir figure 4a). Après cette étape d'accélération, un jet de paroi est produit à partir du bord de l'électrode exposée à l'air, et s'étend le long de l'électrode mise à la terre, comme le montre la figure 4b. Cet écoulement est parallèle à la surface diélectrique, avec une épaisseur de quelques millimètres. De plus, un faible effet d'aspiration est observé près du bord de l'électrode exposée à l'air (juste en amont), ce qui induit une composante de vitesse verticale orientée vers la surface diélectrique. L'amplitude de la composante de vitesse normale à la paroi dans l'écoulement induit est inférieure d'environ un ordre de grandeur à la composante parallèle à la paroi. Enfin, les acquisitions de PIV à résolution temporelle de l'écoulement induit ont révélé que l'actionneur produit un jet qui fluctue à la même fréquence que la forme d'onde du singal alternatif appliqué en entrée. Les deux demi-cycles AC (c.-à-d., positif et négatif) génèrent une vitesse horizontale. Cependant, le demi-cycle négatif induit une plus grande vitesse que le demi-cycle positif ce qui conduit à un jet fluctuant périodiquement à la fréquence d'alimentation.

2.2. Couche limite tridimensionelle sur une aile en flèche

Comme tout écoulement, les couches limites peuvent être soit laminaires, soit turbulentes. Les couches limites laminaires sont caractérisées par des couches parallèles d'écoulement de fluide qui échangent de la quantité de mouvement entre elles par des cisaillements visqueux. Aucun mouvement de particules perpendiculaire à la direction de cet écoulement parallèle ne se produit. En revanche, les couches limites turbulentes sont caractérisées par de très forts gradients de vitesse à proximité de la paroi par rapport à son homologue laminaire. Ainsi, les couches limites turbulentes entraînent des niveaux de frottement pariétal plus élevés. La transition entre les couches limites est le processus par lequel une couche limite initialement laminaire se transforme en un état turbulent. Il est donc extrêmement important de comprendre les mécanismes qui provo-



Figure 5.: The Airbus A321neo aircraft with a backward swept wing.

quent cette transition et éventuellement de les contrôler, car cela peut permettre d'améliorer l'efficacité des corps aérodynamiques en réduisant la traînée de frottement. La seconde moitié du 20e siècle a marqué un nouveau départ pour l'industrie aéronautique. L'induction des moteurs à réaction a permis aux avions de voler à des vitesses plus élevées. Lorsque les avions se rapprochent de la vitesse du son, des ondes de pression sont générées au-dessus des ailes, ce qui augmente la traînée dite d'onde, annulant ainsi dans une certaine mesure les effets bénéfiques des moteurs à réaction. Une solution pour surmonter ce problème consiste à utiliser des ailes en flèche, comme c'est le cas de la majorité des avions modernes. La géométrie de cette aile est telle qu'elle s'incline vers l'arrière (parfois vers l'avant) à partir de son ancrage au niveau du fuselage de l'avion (voir figure 5), au lieu de s'étendre latéralement de façon rectiligne. Alors que l'avion lui-même se déplace à des vitesses transsoniques, les phénomènes aérodynamiques au niveau de la surface de l'aile sont dominés par la composante de vitesse perpendiculaire au bord d'attaque. Ainsi, le nombre de Mach critique de l'avion où apparaissent les ondes de choc est augmenté, ce qui permet à l'avion de se déplacer à des vitesses plus élevées.

Sur le bord d'attaque d'une aile en flèche, où les gradients de pression sont élevés, les lignes d'écoulement sont fortement courbées. Celles-ci, avec le dérapage de l'aile, conduisent à des lignes d'écoulement qui dévient vers le bout de l'aile pour une aile balayée vers l'arrière et vers l'emplanture de l'aile pour une aile balayée vers l'avant. À proximité du corps aérodynamique, un équilibre des forces s'établit dans la direction perpendiculaire à ces lignes d'écoulement, entre la pression et les forces centrifuges. Cependant, cet équilibre des forces est rompu à l'intérieur de la couche limite en raison d'une force motrice plus faible. Il en résulte un écoulement secondaire de faible vitesse, perpendiculaire à la direction de la ligne d'écoulement locale, appelé écoulement transversal (CF). Cette topologie est présentée dans la figure 6.

2.2.1. Instabilité de type primaire

La composante de vitesse de l'écoulement transversal est nulle à la fois à la paroi et sur le bord extérieur de la couche limite, mais elle a un maximum fini à une certaine distance de la surface à l'intérieur de la couche limite (voir figure 6). Ainsi, la composante d'écoulement transversal du profil de vitesse de la couche limite comporte un point d'inflexion qui donne lieu à une instabilité non visqueuse appelée instabilité d'écoulement transversal (Saric, Reed, et al., 2003). Comme on peut le déduire d'un certain nombre de travaux impliquant la visualisation de l'écoulement de surface (voir par exemple Dagenhart et al., 1989; Deyhle & Bippes, 1996;



Figure 6.: Sketch of a generic swept wing. The inviscid streamline (ISL) and the corresponding coordinate system $([xyz]_{ISL})$ are shown. The tangential and cross-flow velocity profiles are depicted in a local inviscid-streamline-oriented coordinate system.

Downs & White, 2013; Serpieri & Kotsonis, 2016; Serpieri, Yadala, et al., 2017), la couche limite sur une ile en flèche présente un motif distinctif, caractérisé par des stries alignées à quelques degrés près avec la direction de l'écoulement (voir figure 7a). Ces stries sont l'empreinte de tourbillons co-rotatifs amplifiés par les modes d'instabilité primaires transversaux montrés dans la mesure par PIV tomographique sur la figure 7b. Ces modes d'instabilité primaires peuvent être de nature stationnaire (stable) ou instationnaire (instable), selon la réceptivité de la couche limite aux perturbations externes. Des modèles de réceptivité expliquent comment les perturbations externes interagissent et influent sur les perturbations à l'intérieur de la couche limite (Reed & Saric, 1989). Ce processus est généralement négligé dans les modèles de prédiction de la transition, mais un certain nombre d'expériences ont montré que la transition dans les couches limites tridimensionnelles montre une forte dépendance à l'égard de l'environnement des perturbations externes. Un aperçu de la réceptivité de la couche limite tridimensionnelle peut être trouvé dans Bippes (1999), Saric, Reed, et al. (2003), Tempelmann et al. (2012), Downs & White (2013), Kurz & Kloker (2014), and Rizzo et al. (2019).

2.2.2. Instabilités de type secondaire

Bien que les tourbillons résultant des instabilités primaires transversales soient relativement faibles, ils provoquent une distorsion de l'écoulement moyen et modifient la stabilité de la couche limite. Le côté ascendant déplace le fluide à faible de la sous-couche visqueuse vers les régions à forte quantité de mouvement de la partie supérieure de la couche limite. Ces actions provoquent une modulation des quantités de mouvement dans le sens de l'envergure, ce qui entraîne le développement de forts cisaillements de vitesse dans la couche limite dans le sens de l'envergure et dans le sens normal de la paroi. Ces couches de cisaillement sont inflexionnelles et instables aux mécanismes d'instabilité secondaires (White & Saric, 2005). Ces instabilités secondaires se développent rapidement dans le sens du courant, déclenchant une transition vers la turbulence. La rupture laminaire dominée par les instabilités stationnaires transverses présente un front de transition caractéristique en forme de coin ou dentelé (Reibert et al., 1996; White & Saric, 2005; Serpieri & Kotsonis, 2016). Si, en revanche, les instabilités primaires instationnaires (en espace) sont amplifiées au départ, elles ont tendance à dégrader les structures stationnaires dans la couche



Figure 7.: (a) Fluorescent oil-flow visualization on the pressure side of a 45° swept-wing, reproduced from Serpieri & Kotsonis (2015). (b) Time-average velocity magnitude along the stationary CF axis measured with tomographic PIV, reproduced from Serpieri & Kotsonis (2016).

limite et donc à conduire à un front de transition plus uniforme sur toute l'étendue (Downs & White, 2013). Malik et al. (1999) ont utilisé des équations de stabilité parabolisées (PSE) linéaires et non linéaires pour analyser l'amplification des instabilités primaires transversales. Ils ont ainsi distingué deux types de modes d'instabilité secondaire. Les premiers, appelés modes de type-I, ont été produits par le gradient de vitesse d'écoulement dans le sens du courant le long de l'envergure $(\partial U/\partial z)$. Ces modes sont composés de fluctuations de haute énergie et de basse fréquence et sont observés sur le bord des tourbillons primaires transversaux, comme le montre la figure 8a (reproduite de Bonfigli & Kloker, 2007). Les seconds, appelés modes de type-II, sont le résultat de gradients importants de la vitesse le long de la direction normale de la paroi $(\partial U/\partial y)$ et sont donc situés au sommet des tourbillons primaires (voir figure 8b). Par rapport aux modes secondaires de type-I, ces modes contiennent des fluctuations à haute fréquence mais avec des énergies plus faibles. Un troisième type d'instabilités secondaires, ou modes de type-III, a également été identifié par la suite (Högberg & Henningson, 1998; Janke & Balakumar, 2000). Ces modes sont un produit de la modulation dans le sens de l'envergure des instabilités primaires de déplacement causées par des tourbillons transversaux stationnaires. Ainsi, leur rôle dans le processus de transition nécessite l'amplification des instabilités primaires stationnaires et instationnaire mobiles, et est principalement pilotée par les fluctuations de niveaux de turbulence en écoulement libre (Downs & White, 2013).

Dans leur expérience, White & Saric (2005) ont fourni des mesures détaillées de l'instabilité secondaire. La couche limite a été forcée avec le mode d'instabilité stationnaire critique en utilisant des DRE distribués le long de l'envergure. Les modes de type-I et de type-II ont été



Figure 8.: Location of secondary instability modes along the stationary CF vortices (dotted black curves), reproduced from Bonfigli & Kloker (2007). (a) *type-I* modes, (b) *type-II* modes and (c) *type-III* modes.

clairement identifiés et leur croissance soudaine a provoqué une transition du régime laminaire. De plus, l'emplacement de ces modes secondaires sur les tourbillons primaires transversaux était en accord avec les modélisations de Malik et al. (1999).

Cette étude (White & Saric, 2005) a surtout révélé que l'apparition des mécanismes d'instabilité secondaires de type-I ou de type-II dépend de la longueur d'onde des instabilités primaires transverses présentes dans la couche limite par rapport à celle de l'instabilité critique. Le mode de type-I est le plus courant et il est observé lorsque les tourbillons primaires transversaux ont une longueur d'onde similaire à celle de l'instabilité primaire critique ou inférieure (sous-critique), en raison du fort cisaillement généré dans le sens de l'envergure. Lorsque des modes stationnaires super-critiques sont forcés (c'est-à-dire des modes avec une longueur d'onde supérieure au mode critique), il en résulte une réduction du cisaillement dans le sens de l'envergure et donc une amplification du mode de type-I. Ainsi, les modes de type-II déclenchent la transition vers la turbulence.

2.3. Phénomènes d'instabilités dans une couche de mélange plane

L'interface de deux flux parallèles l'un à l'autre et ayant des vitesses ou des densités différentes est couramment observée dans de nombreux phénomènes naturels et applications pratiques. Un mélange intense des deux flux se produit dans la région présentant un gradient du paramètre spécifique et, par conséquent, ces types d'écoulements systèmes de fluides sont généralement appelés couches de mélange. L'interaction des deux courants se traduit par une forte action de cisaillement et, par conséquent, la couche de mélange est fondamentalement une représentation d'une classe plus large d'écoulement non visqueux à cisaillement libre. Cette configuration d'écoulement domine un certain nombre de phénomènes, notamment l'écoulement dans les jets et les sillages, la dispersion des polluants dans l'atmosphère, la météorologie en relation avec la formation de cyclones, entre autres.



Figure 9.: (a) Vorticity distribution in a free-shear layer theoretically computed by Michalke (1965b). (b) Smoke visualisation showing vortex formation in a jet flow (flow from left), reproduced from Freymuth (1966).

Les couches de mélange provenant de conditions initiales laminaires sont caractérisées par un mouvement ordonné de structures d'écoulement à grande échelle et atteignent leur autosimilarité beaucoup plus tard que celles qui résultent de conditions initiales turbulentes. Les couches de mélange réellement turbulentes sont quant à elle caractérisées par des mouvements à petite échelle qui permettent un mélange efficace au niveau moléculaire. En fait, les mécanismes qui conduisent à la transition laminaire-turbulente dans les couches de mélange ont été largement étudiés. Si l'état de la couche limite entrante est un facteur primordial dans le développement ultérieur d'une couche de mélange auto-similaire, un certain nombre d'autres paramètres (tels que l'épaisseur du bord de fuite, le niveau de turbulence du courant libre, etc.) peuvent également avoir un effet (Dziomba & Fiedler, 1985).

2.3.1. Instabilité de type primaire

Les écoulements à cisaillement libre, y compris les couches de mélange et les jets, présentent un profil de vitesse avec un point d'inflexion. Il est bien connu, grâce aux travaux menés par Lord Rayleigh (1880), qu'un fluide non-visqueux dont le profil de vitesse comporte un point d'inflexion est instable à certaines perturbations ondulatoires de faible amplitude. Les champs d'écoulement à cisaillement libre sont instables aux perturbations résultant du mécanisme d'instabilité de *Kelvin-Helmholtz* (KH) qui est susceptible de provoquer des instabilités secondaires tridimensionnelles pouvant faciliter l'apparition rapide de turbulences. Taylor (1935) postule que le mouvement au sein de la couche de mélange était homogène et isotrope.

Cependant, les efforts théoriques de Michalke Michalke (1965a,b) ont confirmés l'existence de tourbillons discrets dans les couches de cisaillement libre. En fait, Michalke (1965b) a utilisé l'équation non linéaire de *Helmholtz* sur un profil de vitesse hyperbolique-tangente pour comprendre la formation des tourbillons dans les écoulements à cisaillement libre. Ainsi il a était observé que la couche de cisaillement libre devenait plus mince dans certaines régions et plus épaisse dans d'autres, l'épaisseur étant directement liée à la présence de structures cohérentes large-échelles (figure 9a). Ceci a été très bien corroboré par les expériences de visualisation par injection de fumée, comme celle de Freymuth (1966) dans l'écoulement de jet illustré en figure 9b. Plus tard, Winant & Browand (1974) ont découvert que ces structures fluidiques à large-échelles se développent plus en aval et forment un train périodique de tourbillons bidimensionnels. En raison de petites variations de la force de ces tourbillons ou de l'espacement entre eux, les tourbillons voisins peuvent s'enrouler les uns autour des autres et, par diffusion visqueuse, former un tourbillon unique de plus grande dimension. Il a été constaté que cette association ou fusion de tourbillons était responsable de la croissance de la couche de mélange. De plus, les résultats suggèrent que l'entraînement du fluide irrotationnel ne se produit pas

seulement en raison de la diffusion du vortex de la couche de mélange vers l'écoulement potentiel extérieur. Mais un entraînement plus important se produit lors de l'appariement des tourbillons, où le fluide irrotationnel entre les tourbillons de l'appariement est englouti dans le tourbillon résultant. Dans une couche de mélange non forcée, ces interactions tourbillonnaires se produisent de manière stochastique dans l'espace et le temps, entraînant une croissance linéaire. Ho & Huang (1982) ont étudié l'influence des instabilités aux sub-harmoniques sur l'évolution de la couche de mélange, en focalisant particulièrement l'analyse sur le processus d'appariement des structures cohérentes à large-échelle.

2.3.2. Instabilité de type secondaire

Outre les tourbillons à grande échelle dans le sens de l'envergure issue des instabilités primaires décrites précédemment, Brown & Roshko (1974) ont également observé des stries distinctives dans les visualisations en plan de la couche de mélange. Ces traînées dans le sens de l'écoulement sont causées par des instabilités secondaires prévalant dans le champ d'écoulement de la couche de mélange.

Les tourbillons primaires dans le sens de l'envergure créent une contrainte positive dans la région de la tresse (région entre deux tourbillons primaires) qui étire la vorticité résiduelle dans cette région. Ce mécanisme d'étirement finalement concentre la vorticité, ce qui entraîne la formation de paires de boucles de tourbillons contrarotatives orientées dans le sens de l'écoulement (Lasheras & H. Choi, 1988; Nygaard & Glezer, 1991; Moser & Rogers, 1991; Rogers & Moser, 1992). Ainsi, près du bord de fuite de la plaque séparatrice, la couche de mélange contient des structures très organisées dans le sens de l'envergure. Plus loin, les structures longitudinales secondaires se forment et l'interaction de ces deux types de structures va jouer un rôle important dans le déclenchement de la transition du régime de mélange. En effet, les tourbillons secondaires orientés longitudinalement superposent des turbulences à petite échelle aux structures d'écoulement à grande échelle. Ce changement dans l'état de la couche de mélange favorise une transition de mélange à petite échelle, car il a un effet important sur le mélange au niveau moléculaire.

Cette structuration de l'écoulement de couche de mélange a été confirmée par Lasheras & H. Choi (1988) lors d'une étude expérimentale sur le développement tridimensionnel d'une couche de mélange plane soumise à des perturbations sinusoïdales de petite amplitude. Comme l'instabilité primaire du KH n'a pas été forcée, la couche de cisaillement s'est développée naturellement, formant des tubes tourbillonnaires bidimensionnels dans le sens de l'envergure. Grâce à la visualisation de l'écoulement, les auteurs observent que le développement de ces deux structures d'écoulement vortex est indépendant au départ. Cependant, plus en aval, les tourbillons dans le sens de l'écoulement induisent une ondulation des cœurs tourbillonnaires des structures longitudinales primaires. Cette ondulation des tourbillons dans le sens de l'envergure n'est rien d'autre que l'ondulation observée précédemment par Breidenthal (1981). Finalement, les tourbillons dans le sens de l'écoulement s'enroulent autour du tourbillon primaire, comme le montre la figure 10.



Figure 10.: Plane view of the mixing layer after the secondary vortices wrap around the large-scale spanwise vortices. (b) Topology of the vorticity field between two consecutive spanwise vortices. Reproduced from Lasheras & H. Choi (1988).

3. MOYENS ET MÉTHODES

Les recherches présentées dans cette thèse sont principalement issues d'investigations expérimentales qui s'appuient sur quelques calculs théoriques. Cette thèse comprend l'étude de l'effet d'une décharge plasma sur les phénomènes d'instabilités hydrodynamiques de deux configurations d'écoulement distinctes. Les recherches relatives à la couche limite des ailes en flèche et au contrôle des instabilités transversales associées ont été menées au *département AWEP de la* faculté d'ingénierie aérospatiale de TU Delft, aux Pays-Bas. Les investigations relatives au cas de la couche de mélange plane ont, quant à elles, été réalisées à l'Institut PPRIME qui dépend de l'Université de Poitiers en France.

Ce chapitre présente les différents outils utilisés pour mener à bien les recherches de ce doctorat. Les installations expérimentales qui comprennent les différentes souffleries utilisées et les modèles employés, sont décrites. Les différents métrologies utilisées telles que la thermographie infrarouge (IR) pour analyser l'évolution de la transition laminaire-turbulente sur l'aile en flèche, la vélocimétrie par image de particules (PIV) pour obtenir des données quantitatives sur les vitesses ainsi que les différentes méthodes utilisées pour caractériser la performance des actionneurs de plasma DBD. Enfin, divers outils analytiques utilisés pour étudier les données expérimentales acquises sont examinés.



Figure 11.: The low-turbulence wind tunnel at at Delft University of Technology, The Netherlands.

3.1. Banc expérimentaux principaux

3.1.1. Couche limite tridimensionnelle : Soufflerie, modèle et conditions d'essai

Les expériences sur la couche limite tridimensionnelle qui se développe sur une aile en flèche ont été menées dans la soufflerie à faible turbulence (LTT, figure 11) en boucle fermée de TU Delft, aux Pays-Bas. Les dimensions de la chambre d'essai sont de 1.25 m × 1.80 m × 2.6 m en hauteur, largeur et longueur, respectivement. La soufflerie est équipée de sept grilles successives suivis d'un convergent dont le rapport de contraction est de 17 : 1. Une caractérisation récente de l'écoulement dans cette soufflerie a révélé que la turbulence était de $Tu/U_{\infty} = 0.024\%$ à une vitesse d'écoulement libre $U_{\infty} = 24$ m s⁻¹.

Le modèle (appelé M3J) utilisé dans cette thèse est une aile en flèche de 45°. L'aile utilisée est une version modifiée de la forme NACA 66018 nommée 66018M3J. Le modèle d'aile est tel que le bout de l'aile inférieure est en amont par rapport au bout de l'aile supérieure. Le modèle a une envergure de 1,25 m et une corde dans le sens de l'écoulement de c = 1.27 m. La longueur de la corde perpendiculaire au bord d'attaque est de $c_x = 0.898$ m. Le modèle est fait de fibre de verre avec une surface polie (écart type de rugosité de surface de $0.2 \ \mu m$). Le modèle d'aile est équipé de deux rangées de 46 prises de pression pariétale orientées dans le sens de l'écoulement à l'infini. Ces prises ont été connectées à un transducteur pour mesurer la distribution de la pression à la surface du modèle. Ainsi, aucune différence n'est perceptible dans la distribution de la pression du côté de l'aile, mesurée par les deux réseaux de prises de pression de surface. La présence d'instabilité transverse exige que le flux d'air soit invariant le long du modèle (Deyhle & Bippes, 1996; Reibert et al., 1996; Radeztsky et al., 1999). Pour un angle d'incidence de l'aile de $\alpha = 3^{\circ}$ et une vitesse d'écoulement libre de $U_{\infty} = 25 \text{ m s}^{-1}$ ($Re_c = 2.1 \cdot 10^6$, ce qui correspond aux conditions d'écoulement principalement étudiées dans cette thèse), l'absence de différence dans les mesures de pressions indique un degré élevé d'uniformité de l'écoulement le long de l'envergure comme souhaité.

3.1.2. Couche de mélange plane : Soufflerie, modèle et conditions d'essai

Les expériences sur la couche de mélange plane ont été réalisées dans une soufflerie en boucle fermée, spécialement modifiée pour permettre des études sur cette configuration à l'Institut PPrime de l'Université de Poitiers, en France. Les dimensions de la section d'essai de cette souf-



Figure 12.: (a) Wind-tunnel set-up. Location of the foams and perforated plate (dark rectangle), metallic grids (four vertical lines) and the splitter plate (red) are depicted. (b) Schematic (not to scale) of PMMA attachment (light grey) fixed to the aluminium plate (black) to complete the splitter-plate. Spanwise incision (dark grey) to house an electrode of the DBD plasma actuator is shown. Coordinate system used in this study is presented in both figures.

flerie sont de 30 cm \times 30 cm \times 1.2 m en hauteur, largeur et longueur. Une plaque d'aluminium (plaque séparatrice) de 15 mm d'épaisseur s'étend de l'entrée de la chambre de tranquillisation, juste après les grilles en nid d'abeille, à travers le convergent et jusqu'à l'entrée de la section d'essai, comme l'illustre la figure 12a (ligne rouge épaisse). Cette plaque d'aluminium divise la soufflerie en deux moitiés égales dans le sens vertical, générant ainsi deux flux d'écoulement parallèles. L'installation a été conçue de telle sorte que la vitesse du courant supérieur (U_2) soit différente de celle du courant inférieur (U_1), afin de réaliser une configuration plane de la couche de mélange entre deux flux parallèles.

Une grille métallique (avec un diamètre de fil de 0.18 mm et une solidité de 35%) a été installée juste en amont du convergent dans le flux d'air supérieur afin de casser les tourbillons entrants et réduire la turbulence de l'écoulement. L'intensité de la turbulence a été calculée à partir de mesures de vitesse acquises avec un capteur de type fil chaud (bande passante filtrée entre 10 et 1000 Hz). L'intensité de la turbulence du flux d'air à grande vitesse à $U_2 = 10 \text{ m s}^{-1}$ est de $Tu_2/U_2 = 0.22\%$. Afin de réaliser une différence de vitesse entre les deux flux, des pertes de charges ont été installés dans le flux inférieur, juste en amont du convergent (rectangle noir avec d'épaisses lignes verticales sur la figure 12a). Une combinaison de mousses et de plaques perforées a été installée ici pour réduire la vitesse du flux inférieur. Un ensemble de quatre grilles métalliques (les mêmes que dans le courant supérieur), séparées les unes des autres de 8 mm, a été installé en aval de la plaque perforée pour réduire le niveau de turbulence également dans ce flux. Grâce à ces modifications, la vitesse du flux inférieur a été réduite de moitié par rapport à celle du flux supérieur. Ainsi, le rapport de vitesse entre les deux courants $r (= U_1/U_2) = 0.5$ (R = (1 - r)/(1 + r) = 0,33). Le niveau de turbulence dans le flux inférieur a été mesuré avec le même fil chaud. Lorsque la vitesse du flux inférieur était $U_1 = 10 \text{ m s}^{-1} (U_2 = 20 \text{ m s}^{-1})$, l'intensité de la turbulence (bande passante filtrée entre 10 et 1000 Hz) était $Tu_1/U_1 = 1.5\%$. À $U_1 = 5 \text{ m s}^{-1}$ ($U_2 = 10 \text{ m s}^{-1}$), l'intensité de turbulence a été mesurée à $Tu_1/U_1 = 1.3\%$.

3.2. Métrologies optiques

Comme tout écoulement, les couches limites peuvent être soit laminaires, soit turbulentes. Les couches limites laminaires sont caractérisées par des couches parallèles d'écoulement de fluide

qui échangent de la quantité de mouvement entre elles par des cisaillements visqueux. Aucun mouvement de particules perpendiculaire à la direction de cet écoulement parallèle ne se produit. En revanche, les couches limites turbulentes sont composées d'un mouvement chaotique des particules. Ces couches limites sont caractérisées par de très forts gradients de vitesse à proximité de la paroi par rapport à son homologue laminaire. Ainsi, les couches limites turbulentes entraînent des niveaux de frottement pariétal plus élevés. La transition entre les couches limites est le processus par lequel une couche limite initialement laminaire se transforme en un état turbulent. Il est donc extrêmement important de comprendre les mécanismes qui provoquent cette transition et éventuellement de les contrôler, car cela peut permettre d'améliorer l'efficacité des corps aérodynamiques en réduisant la traînée de frottement. La seconde moitié du 20e siècle a marqué un nouveau départ pour l'industrie aéronautique. L'induction des moteurs à réaction a permis aux avions de voler à des vitesses plus élevées. Lorsque les avions se rapprochent de la vitesse du son, des ondes de pression sont générées au-dessus des ailes, ce qui augmente la traînée dite d'onde, annulant ainsi dans une certaine mesure les effets bénéfiques des moteurs à réaction. Une solution pour surmonter ce problème consiste à utiliser des ailes en flèche, comme c'est le cas de la majorité des avions modernes. La géométrie de cette aile est telle qu'elle s'incline vers l'arrière (parfois vers l'avant) à partir de son ancrage au niveau du fuselage de l'avion (voir figure 5), au lieu de s'étendre latéralement de façon rectiligne. Alors que l'avion lui-même se déplace à des vitesses transsoniques, les phénomènes aérodynamiques au niveau de la surface de l'aile sont dominés par la composante de vitesse perpendiculaire au bord d'attaque. Ainsi, le nombre de Mach critique de l'avion où apparaissent les ondes de choc est augmenté, ce qui permet à l'avion de se déplacer à des vitesses plus élevées.

3.2.1. PIV tomographique

La PIV tomographique à résolution temporelle a été utilisée dans cette thèse pour acquérir des données de vitesse à trois composantes du champ d'écoulement induit par un actionneur de plasma DBD. La figure 13 illustre la mise en place de cette expérience. Cette série d'acquisition a été réalisée pour des conditions de repos, c'est-à-dire sans écoulement extérieur afin de caractériser les performances de l'actionneur DBD. L'actionneur a été placé à l'intérieur d'une enceinte équipée d'accès optiques. L'air dans l'enceinte été ensemencé avec des gouttelettes d'huile de paraffine, d'un diamètre moyen de 1 μ m, générées par une buse de pulvérisation TSI. L'illumination des particules a été assurée par un laser PIV Nd:YAG (18 mJ par impulsion). Le laser est placé à l'extérieur de l'enceinte. Le faisceau laser a été transformé en une nappe de 4 mm d'épaisseur en utilisant une optique spécifique (combinaison de lentilles cylindriques et sphériques) et il a été orienté parallèlement à la surface diélectrique. L'objectif ici est de vérifier la tridimensionnalité de l'écoulement lorsque qu'un actionneur plasma modulé spatialement est utilisé. Dans ces conditions, bien que longue et couteuse en équipement et temps de calcul, la technique de tomographie PIV semble la plus adaptée.

Le système d'imagerie était composé de quatre caméras numériques Photron FastCAM SA1 dotées chacune d'un capteur de 1024 px² (pixels de 20 μ m et résolution numérique de 12 bits). Le capteur des caméras a été réduit à 960 × 440 px². Les caméras étaient équipés d'objectif Nikon-Nikkor de 200 mm avec une ouverture numérique de $f_{\#} = 8$. Le volume imagé était de $(x_p \times y_p \times z_p =)$ 33 × 4 × 16 mm³. Un mécanisme de décalage d'objectif était fixé entre les appareils et les objectifs afin de respecter la condition Scheimflug et d'ajuster l'alignement du plan focal sur les capteurs des appareils. Le laser et les quatre caméras ont été synchronisés à l'aide d'un contrôleur haute vitesse Lavision et de la suite Lavision Davis 8.2. L'acquisition des images, la reconstruction du volume et la corrélation croisée ont été effectuées à l'aide de la



Figure 13.: Tomographic PIV set-up to acquire three-component velocity data of the DBD plasma actuator's induced flow field. The x_p direction is represented by the red arrow.

même suite logicielle. Le taux d'acquisition était de $f_s = 10$ kHz en mode image par image, ce qui donne une séparation d'impulsion de 100 µs entre des images consécutives.

Une cible à double couche a d'abord été utilisée pour l'étalonnage de ce système de mesure et la fonction de cartographie obtenue a été corrigée à l'aide de la technique d'autoétalonnage du volume (Wieneke, 2008). Comme la mesure actuelle est résolue dans le temps, l'opérateur MART (Lynch & Scarano, 2015) à suivi de mouvement simultané (SMTE-) a été utilisé pour la reconstruction volumétrique des champs acquis. La corrélation croisée pour obtenir les champs vectoriels de vitesse a été effectuée avec un volume d'interrogation final de $48 \times 32 \times 24$ voxels³, avec un chevauchement de 75%. Cela a permis d'obtenir une résolution vectorielle finale de 0.22 mm dans les trois directions.

3.2.2. Thermographie Infrarouge

Une nécessité essentielle dans cette thèse est de visualiser et d'étudier la topologie de la couche limite qui s'étend sur l'aile en flèche et de détecter la position de la transition laminaireturbulente si elle existe. Une technique utilisée pour visualiser la topologie de la transition d'un écoulement est la thermographie infrarouge (IR) qui capte le rayonnement IR émis par un corps. Le corps aérodynamique testé est chauffé en continu par une source externe (comme des circuits imprimés chauffant le modèle par effet Joule, des radiations de lampes halogènes, etc.) et refroidi par l'écoulement par convection. Il s'agit d'une technique non intrusive, ce qui est une caractéristique importante pour l'utilisation visée, car même les particules minuscules introduites dans les installations d'essai peuvent entraîner une évolution et une transition de la couche limite sensiblement différentes en raison du mécanisme de réceptivité. Ainsi, dans la présente thèse, la visualisation du flux IR est largement utilisée pour inspecter l'évolution de la transition de la couche limite de l'aile en flèche. Il convient toutefois de noter que la conduction thermique à travers la surface du modèle doit être négligeable, car cela permet de mieux déterminer l'effet de l'écoulement, en augmentant la sensibilité thermique du système de mesure par rapport à l'écoulement. Le modèle d'aile balayée utilisé dans cette thèse est fait de fibre de verre qui n'est pas un bon conducteur thermique. Il facilite donc l'utilisation de la thermographie infrarouge pour la visualisation de l'écoulement. Ici, le modèle a été éclairé à l'aide de six lampes halogènes de 1 kW à travers les vitres de la section d'essai pour maintenir



Figure 14.: (a) Captured calibration image corresponding to the *global IR-FOV*. (b) Raw image of the thermal signature of the boundary layer over the swept-wing model's surface. (c) Full IR field, after the application of dewarping/deskewing. (d) Detected transition front in the region enclosed by the dashed-white rectangle in (c).

un flux de chaleur constant et améliorer le contraste. La caméra IR utilisée pendant l'expérience était une *Optris PI64*0. Elle comporte une matrice de plan focal non refroidie de $640 \times 480 \text{ px}^2$ qui répond à la plage spectrale électromagnétique comprise entre 7.5 et 13 μ m. Sa sensibilité thermique (NETD) est de 75 mK. La caméra IR a également été placée à l'extérieur de la soufflerie, à environ 1 m du modèle. Deux champs de vision différents (FOV) ont été captés pour inspecter les différentes caractéristiques de la couche limite. Pour le premier champ de vision, la caméra IR était équipée d'un objectif grand angle de $15^{\circ} \times 11^{\circ}$ avec une longueur focale de 41.5 mm afin de capturer l'écoulement sur l'ensemble du modèle et de détecter la transition laminaire-turbulente. Cependant, la résolution spatiale obtenue lors de cette série de mesure ne permet pas d'inspecter la longueur d'onde des modes CFI stationnaires considérés ici. À cet effet, un deuxième FOV zoomé a été capturé en équipant la caméra IR d'un objectif $60^{\circ} \times 45^{\circ}$ d'une longueur focale de 10.5 mm.

La caméra enregistre 100 champs IR à un taux de 5 Hz pour chaque cas d'écoulement testé. Le signal moyen est utilisé afin d'augmenter le rapport signal/bruit. Le traitement des images IR et la détection du lieu de transition ont été réalisés à l'aide d'un logiciel *MATLAB* interne (département AWEP de l'université technologique de Delft). La procédure est illustrée à la figure 14 (voir également Rius Vidales et al., 2018). Tout d'abord, une feuille d'étalonnage a été conçue et installée sur l'aile. Le modèle d'aile a ensuite été placé à l'angle d'attaque étudié et la feuille de calibrage a été imagée avec la caméra IR. La figure 14a présente une image de calibrage correspondant à l'IR-FOV global. La feuille de calibrage a ensuite été retirée et la surface de l'aile nettoyée, avant d'acquérir les filaments IR nécessaires, dont un exemple est présenté à la figure 3.6b. Pour tracer ces images IR dans le plan xz, on applique un traitement

des images (dewarping/deskewing) en utilisant les images de calibrage prédéfinies, ce qui permet de supprimer les erreurs de courbure et de parallaxe (angle d'observation) de la surface de l'aile.

La différence de température entre les régions à écoulement laminaire et turbulente peut être observée dans la figure 14c (régions claires et sombres respectivement). Cette différence de température est exploitée dans cette thèse et le point de transition est identifié comme l'emplacement moyen du gradient de température maximum dans le sens de l'écoulement (voir les points bleus dans figure 14d). Pour augmenter la robustesse de cette méthode d'identification, les valeurs aberrantes sont éliminées sur la base de l'écart-type des positions détectées dans le sens de l'écoulement. Le motif dentelé typique du front de transition nécessite l'utilisation d'une ligne d'interpolation linéaire passant par les points de transition le long de l'envergure, qui est représentée par la ligne rouge continue en figure 14d. L'erreur standard de cette interpolation linéaire est utilisée comme une estimation de la largeur du front de transition (ligne rouge en pointillés en figure 14d).

3.3. Outils d'analyse

Une fois l'acquisition des données terminée, un traitement supplémentaire de ces ensembles de données est nécessaire pour faciliter l'analyse des configurations de flux correspondantes. Un certain nombre d'outils analytiques sont déployés dans le cadre de cette thèse, soit pour utiliser les données expérimentales acquises afin de faire des prévisions basées sur des bases théoriques, soit pour aider à mieux comprendre le phénomène physique sous-jacent.

Ces outils sont présentés dans cette section.

3.3.1. Analyse par stabilité linéaire

La théorie de la stabilité linéaire (LST) a été traditionnellement utilisée dans les études sur les instabilités hydrodynamiques à la fois dans les couches de cisaillement libre et les couches limites. C'est un outil polyvalent qui permet de prédire l'apparition d'instabilités, les longueurs d'onde et/ou fréquences associées et les taux de croissance spatiale. Il repose sur l'hypothèse qu'un écoulement de base parallèle (c'est-à-dire que l'écoulement ne varie pas dans le sens de l'amplification des perturbations) est instable aux perturbations ondulatoires infinitésimales. La formulation du solveur LST est brièvement décrite ci-dessous.

Le champ d'écoulement instantané peut être étendu comme suit: $U = \overline{U} + \widetilde{U}$ où $\overline{U}(=[\overline{u}, \overline{v}, \overline{w}])$ est l'écoulement de base qui satisfait à l'équation de Navier-Stokes constante et incompressible et $\widetilde{U}(=[\widetilde{u}, \widetilde{v}, \widetilde{w}])$ représente les perturbations ou les fluctuations de vitesse. Si le flux de base est supposé être parallèle et que les perturbations sont considérées comme infinitésimales et semblables à des vagues, ces perturbations peuvent être représentées par la fonction suivante,

$$\widetilde{U}(x, y, z, t) = \psi(y)e^{i(\alpha x + \beta z - \omega t)}$$
(1)

où $\psi(y)$ est la fonction propre de la perturbation, α et β sont les nombres d'ondes complexes dans le sens de l'écoulement et de l'envergure respectivement, et ω est la fréquence angulaire. Comme les instabilités augmentent dans le sens de l'écoulement, l'évolution du nombre d'ondes dans cette direction α (= $\alpha_r - i\alpha_i$) est particulièrement intéressante, car -i représente le taux de croissance des perturbations correspondant aux différents modes d'instabilité. Le nombre d'ondes dans le sens de l'envergure () est également complexe. Cependant, l'invariance du champ d'écoulement dans le sens de l'envergure donne $_i = 0$. La fréquence angulaire (ω) est réelle car la formulation spatiale de la LST est appliquée ici comme déjà mentionné.

L'évolution de l'écoulement et donc des perturbations qu'il contient est régie par les équations incompressibles de Navier-Stokes (voir Schlichting & Gersten, 2016, pour la dérivation). L'expansion du flux instantané est appliquée à l'équation de Navier-Stokes et tous les termes ne comportant que l'écoulement de base sont négligés. De plus, en partant de l'hypothèse de perturbations infinitésimales, tous les termes quadratiques de ces composantes fluctuantes sont également négligés. Il en résulte les équations de Navier-Stokes linéarisées qui régissent l'évolution des perturbations et qui sont données dans l'équation 2.

$$\nabla \cdot \tilde{U} = 0 \tag{2a}$$

$$\frac{\partial \widetilde{U}}{\partial t} + (\overline{U} \cdot \nabla) \widetilde{U} + (\widetilde{U} \cdot \nabla) \overline{U} = -\frac{\Delta \widetilde{p}}{\rho} + \nu \nabla^2 \widetilde{U}$$
(2b)

Ici, ρ est la densité du fluide et ν est la viscosité cinématique. Les équations 2a et 2b représentent respectivement la conservation de la masse et du moment. En remplaçant la fonction de flux de la perturbation donnée dans l'équation 3.3 par l'équation 2 et en négligeant \tilde{u} , \tilde{w} et \tilde{p} , on obtient l'équation d'Orr-Sommerfeld.

$$\left[\left(\frac{\partial^2}{\partial y^2} - \alpha^2 - \beta^2 \right)^2 - iRe(\alpha \overline{u} + \beta \overline{w} - \omega) \left(\frac{\partial^2}{\partial y^2} - \alpha^2 - \beta^2 \right) + iRe\left(\alpha \frac{\partial^2 \overline{u}}{\partial y^2} + \beta \frac{\partial^2 \overline{w}}{\partial y^2} \right) \right] \psi(y) = 0$$
(3)

La fonction propre $\psi(y)$ avec les conditions limites défiées est ensuite mappée sur le domaine de Tchebychev et, en conséquence, les dérivés de deuxième et quatrième ordre sont exprimés sous forme de matrices de dérivés de Chebyshev. Ces calculs sont effectués avec l'algorithme (et la suite MatLab) fourni par Weideman & Reddy (2000). Comme indiqué précédemment, les perturbations liées aux différentes instabilités ne croissent que dans le sens du courant (x)et, par conséquent, β et ω sont tous deux fournis en entrée de l'équation. En conséquence, l'équation 3 peut être réorganisée comme un problème aux valeurs propres avec uniquement le nombre d'onde dans le sens du courant . Ce problème est ensuite résolu selon la technique de matrice complémentaire proposée par Bridges & Morris (1984). Le facteur N d'amplification (Van Ingen, 1956, 2008) de chaque mode d'instabilité avec une fréquence et une longueur d'onde données $\lambda_z (= 2\pi/\beta_r)$ est évalué par l'intégrale explicitée dans l'équation 4,

$$N(x,\lambda_z,\omega) = \int_{x_0(\lambda_z,\omega)}^x -\alpha_i(x,\lambda_z,\omega)dx$$
(4)

où, $x_0(\lambda_z, \omega)$ est la première position où le mode d'instabilité considéré est instable. Dans les deux configurations d'écoulement étudiées, cette formulation LST est utilisée pour estimer les

caractéristiques de stabilité du champ d'écoulement respectif.

3.3.2. Decomposition orthogonal aux modes propres dans l'espace spectral

Afin de mieux comprendre les différents modes d'instabilité de la couche de mélange plane et, éventuellement, analyser la croissance de la couche, on cherche à déterminer l'organisation spatiale de l'écoulement pour une valeur fréquentielle donnée.. À cette fin, une décomposition orthogonale aux modes propres (POD) appropriée a été appliquée sur l'ensemble de la série temporelle obtenue par PIV. La POD, proposée pour la première fois par Lumley (1967), est une technique permettant d'identifier objectivement les structures cohérentes dans un écoulement turbulent. Une forme spécifique de POD appelée décomposition orthogonale propre spectrale (SPOD) (Delville et al., 1999; Towne et al., 2018) est utilisée dans cette thèse. Les modes SPOD décrivent des structures dans un flux statistiquement stationnaire qui sont cohérentes à la fois dans l'espace et dans le temps. Ici, la décomposition du tenseur trans-spectral (transformée de Fourier directe du tenseur de corrélation obtenue à partir de données expérimentales) est effectuée. Ainsi, cette méthode aboutit à des modes propres qui oscillent chacun à une seule fréquence, ce qui est avantageux pour étudier la configuration de cette thèse.

La formulation de la SPOD se trouve dans Delville et al. (1999) qui l'ont également appliquée sur une couche de mélange plane et ont observé l'organisation des structures primaires et secondaires dans le premier mode. Cependant, ils ont appliqué le SPOD sur des données de vitesse acquises à partir de mesures ponctuelles par fil chaud alors qu'ici, il est appliqué pour la première fois sur des mesures de PIV corrélées dans l'espace et dans le temps. Dans un article récent, Towne et al. (2018) montrent également la formulation et la procédure de la SPOD et discutent de sa relation avec la décomposition en mode dynamique et l'analyse de résolution. En fait, une implémentation MatLab fournie par les auteurs est utilisée ici pour effectuer des mesures SPOD sur nos données expérimentales.

La méthode est maintenant brièvement décrite. Considérons une variable stochastique $Q(\Omega, t)$ qui, dans les travaux actuels, est constituée par les fluctuations des séries temporelles acquises au moyen de la PIV. Ici, Ω représente le domaine spatial des mesures tandis que t est le temps. Dans un premier temps, les champs vectoriels de vitesse mesurés sont convertis du domaine temporel au domaine fréquentiel en utilisant une transformée de Fourier discrète (TFD) comme suit:

$$\widetilde{Q}(\Omega, f) = \int_{-\infty}^{\infty} \widetilde{Q}(\Omega, t) \cdot e^{-i2\pi f t} dt$$
(5)

Le tenseur trans-spectral à chaque fréquence est alors calculé par l'équation,

$$S(\Omega, \Omega', f) = \widetilde{Q}^*(\Omega, f)\widetilde{Q}(\Omega, f)$$
(6)

où * représente le complexe conjugué. Le problème de valeur propre intégrale à résoudre pour décomposer le champ de flux est alors donné par l'équation 7a,

$$\int_{\Omega} S(\Omega, \Omega', f) W(\Omega') \Psi(\Omega', f) d\Omega = \Lambda(f) \Psi(\Omega, f)$$
(7a)

$$\Theta(f) = \int \widetilde{Q}(\Omega, f) \Psi^*(\Omega, f) \sqrt{\Lambda(f)} d\Omega$$
(7b)

où, $\Psi(\Omega, f)$ est la fonction propre dépendant de la fréquence et $\Lambda(f)$ est la valeur propre correspondante qui est représentative de l'énergie cinétique turbulente intégrée (TKE) des fluctuations de ce mode particulier. La variable $W(\Omega)$ est un Hermitien à définition positive. Les vecteurs propres correspondants (Θ) sont calculés par l'équation

La méthode de Welch (1967) est utilisée ici. La série temporelle est décomposée en un certain nombre de sous-ensembles ou blocs qui se chevauchent et qui comportent chacun un nombre égal d'images (N_{fr}) , et la TFD est calculée sur chacun de ces blocs. Les composantes de Fourier d'une fréquence particulière dans tous les blocs sont empilées et le SPOD est effectué. On obtient ainsi un nombre de modes SPOD égal au nombre de blocs à chaque fréquence. Ci-après, les notations $\Psi(f,k)$ et $\Lambda(f,k)$ seront utilisées pour désigner respectivement la fonction propre et les valeurs propres de chaque mode (k) à une fréquence particulière (f).

La détection des structures cohérentes en rotation et l'étude de leur dynamique et de leurs interactions dues au forçage appliqué sont souhaitées. À cette fin, il est nécessaire de reconstruire les champs de vitesse à partir des modes SPOD. Ceci est fait selon l'équation 8.

$$\widetilde{Q}_{R}(\Omega, t) = \sum_{f} \left[\sum_{k} \Theta(f, k) \Psi(f, k) \sqrt{\Lambda(f, k)} \right] \cdot e^{i2\pi f t}$$
(8)

4. DÉVELOPPEMENT ET CARACTÉRISATION D'ACTIONNEURS DBD

Les actionneurs plasma DBD font partie du thème central de cette thèse car ils fournissent à priori le forçage suffisant pour espérer pouvoir réaliser un contrôle des instabilités et de l'écoulement résultant sur les deux configurations d'essais. Afin d'atteindre les objectifs de cette thèse quelques innovations concernant les actionneurs plasma DBD sont impératives. La première question qui se pose est celle de l'épaisseur de l'électrode exposée à l'air. Si ce paramètre géométrique ne respecte pas des contraintes strictes, l'introduction d'actionneurs à plasma DBD à la surface d'un modèle aérodynamique peut avoir un effet néfaste et provoquer la déstabilisation de l'écoulement qui est à contrôler. C'est notamment le cas de la couche limite sur une aile en flèche qui est très réceptive à la rugosité de surface. Ainsi, un des objectifs de cette thèse est de définir une nouvelle technique de fabrication pour générer des actionneurs plasma avec des électrodes respectant des contraintes strictes en termes d'épaisseur et de rugosité.

La technique de fabrication doit également pouvoir répondre à une seconde contrainte. En effet, l'instabilité transverse stationnaire apparaissant dans la couche limite tridimensionnelle étudiée et les instabilités secondaires dans la couche de mélange plane sont des exemples d'instabilités se produisant dans l'espace. Les instabilités transverses instationnaires et les instabilités secondaires dans la couche limite d'une aile en flèche ainsi les instabilités primaires de type *Kelvin*-



Figure 15.: Design of employed DBD plasma actuators (not for scale). (a) To apply spanwise-uniform EFD forcing. The dielectric substrate, air-exposed electrode (E.E., width w_E), covered or encapsulated electrode (C.E., width w_C) are shown. (b) To apply spanwise-modulated EFD forcing. Width and length of the stems (w_S and l_S respectively), spanwise wavelength (λ_z), and relative overlap between the electrodes (o) are represented. In both schematics, the spanwise extent on which the actuator applies flow control (z_{PA}) and distance between plasma-generating edge of the air-exposed electrode and the edge of the dielectric surface (x_{PA}) are shown as well.

Helmholtz dans les couches de mélange planes sont des exemples d'instabilités dépendantes du temps et de la fréquence. Les instabilités temporelles peuvent être contrôlées par le signal AC fourni à l'actionneur. Les instabilités spatiales, en revanche, nécessitent l'application d'un forçage modulé qui peut être obtenu grâce à la géométrie des électrodes. Ainsi, le second objectif est d'intégrer la contrainte de pouvoir produire des électrodes à géométrie complexe avec la technique développée.

Les actionneurs plasmas développés dans ce chapitre sont similaires à ceux qui seront utilisés par la suite dans des scénarii de contrôle d'écoulement via une action ciblée sur les phénomènes d'instabilités. Il est donc nécessaire d'en vérifier le fonctionnement. Pour cela, des essais de caractérisation électromécanique ont été réalisés à l'aide de différentes métrologies électriques et mécaniques. Ici, seuls quelques-uns des résultats principaux sont présentés.

4.1. Géométries étudiées

Les deux différentes configurations d'actionneurs plasma DBD qui sont les plus utilisées dans cette thèse sont présentées ici. Les schémas de ces deux types d'actionneurs sont présentés sur la figure 15. Ils comprennent des actionneurs qui exercent une force volumique EFD bidimensionnelle, homogène selon l'envergure, et ceux qui créent une force volumique modulée en trois dimensions.

Les actionneurs plasma DBD fabriqués pour exercer une force uniforme sont ceux indiqués sur la figure 15a. Ils sont composés de bandes uniformes de métal conducteur comme électrodes, w_E et w_C étant respectivement les largeurs des électrodes exposées à l'air et encapsulées. Ces électrodes sont fabriquées de manière asymétrique sur les deux faces du substrat diélectrique. Ici, la formation de plasma se produit à l'interface des deux électrodes (z_{PA} en figure 15a), ce qui entraîne un effort volumique bidimensionnel. Il s'agit de la conception classique des actionneurs à plasma DBD qui a été largement étudiée dans le passé (voir Moreau, 2007; Corke, Enloe, et al., 2010; Benard & Moreau, 2014). Dans le cas des actionneurs à plasma DBD qui appliquent un forçage modulé dans le sens de l'envergure, la conception de ces actionneurs est présentée dans la figure 15b. L'électrode exposée à l'air est une bande métallique uniforme comme auparavant, tandis que l'électrode encapsulée est composée d'une série de tiges équidistantes. La longueur et la largeur de chaque tige sont respectivement de l_S et de w_S . Ici, de par la géométrie singulière de l'électrode encapsulée (électrode de masse), la zone d'ionisation ne devrait correspondre qu'aux zones d'interface entre les électrodes (c'est-à-dire là où les électrodes se chevauchent). Cette approche assure vraisemblablement une modulation de la force volumique selon l'envergure. L'espacement entre les tiges adjacentes définit la longueur d'onde (λ_z) du forçage du plasma dans le sens de l'envergure.

4.2. Actionneurs DBD à faible épaisseur

Un problème commun à toutes les applications aérodynamiques des décharges plasma utilisées en tant qu'actionneur est l'épaisseur des électrodes. Si l'épaisseur, en particulier de l'électrode exposée à l'air, est comparable à celle de la couche de cisaillement contrôlée, la simple présence de l'actionneur peut potentiellement modifier le contenu en fréquence et les longueurs d'onde d'instabilité dans la couche de cisaillement, déstabilisant ainsi l'écoulement et provoquant une transition précoce vers la turbulence. La situation n'est pas très différente dans le cas de la couche limite sur une aile en flèche. Afin de contrôler l'instabilité primaire et la transition éventuelle qu'elle provoque, le système de forçage doit être appliqué près du bord d'attaque (Reibert et al., 1996; Radeztsky et al., 1999; Tempelmann et al., 2012; Serpieri, Yadala, et al., 2017). Toutefois, cette configuration d'écoulement est très réceptive, en particulier près du bord d'attaque, à la rugosité de surface (Bippes, 1999; Saric, Reed, et al., 2003; Downs & White, 2013; Rizzo et al., 2019). Ainsi, si les électrodes de l'actionneur n'adhèrent pas à des contraintes strictes d'épaisseur et de rugosité, elles peuvent potentiellement induire des perturbations transverses stationnaires supplémentaires dans la couche limite, favorisant de façon non souhaitée la transition.

4.2.1. Méthode de fabrication

Avec ce jeu de contrainte technique, on envisage une technique de fabrication qui consiste à pulvériser des particules d'argent conductrices d'une taille inférieure au micron dans l'espace négatif laissé par un masque physique qui porte les caractéristiques géométriques de l'électrode à réaliser. Pour garantir la répétabilité et un dépôt uniforme du métal sur toute la surface qui servir d'électrode, l'ensemble du processus a été réalisé à l'aide d'un traceur (CNC)-XY *Tooli* à commande numérique par ordinateur. Les différentes étapes de ce processus de fabrication sont brièvement expliquées ici.

Étape 1: Conception de l'électrode

Un dessin vectorisé de l'électrode nécessaire est généré à l'aide de n'importe quelle suite de dessin assisté par ordinateur (CAO) disponible dans le commerce. Cette tâche est essentielle pour articuler les caractéristiques géométriques requises des électrodes et pour maintenir des tolérances strictes. Ces dessins vectorisés sont utilisés pour générer les séquences nécessaires de commandes de machine afin de régir le mouvement des différentes têtes d'outils du traceur CNC-XY.

Étape 2: Génération du masque

La surface diélectrique est d'abord préparée pour la fabrication. Une feuille noire, non réfléchissante

et autocollante (*DC Fix Decorative Foil 346-8091*) est déposée sur le substrat diélectrique. Elle est placée dans la zone de travail du traceur CNC-XY de manière à ce que l'électrode soit fabriquée à l'endroit voulu. La forme géométrique de l'électrode est découpée dans la feuille autocollante à l'aide d'une tête laser de 2.5 W (445 nm, diode laser de 9 mm avec une lentille de verre à trois éléments). La puissance du laser et le nombre de passages ont été choisis de telle sorte qu'il ne coupe que la feuille autocollante et n'a aucun effet sur la surface diélectrique. Une fois la découpe laser terminée, l'autocollant ayant la forme de l'électrode est retiré du substrat diélectrique, générant ainsi un masque approprié à utiliser lors de l'étape 3. La génération de masques de la manière indiquée est avantageuse car elle améliore la précision globale du processus de fabrication, et elle réduit les risques de mauvais alignement des électrodes. Toutefois, il peut y avoir d'autres moyens de générer les masques nécessaires. Il faut souligner que le placement et l'alignement précis des deux électrodes de part et d'autre du substrat diélectrique est une priorité afin d'éviter la discontinuité du champs électrique et de garantir ainsi une force volumique homogène.

Étape 3 : Pulvérisation contrôlée par ordinateur

Une fois les masques générés, une tête d'outil avec une buse de pulvérisation est fixée au traceur CNC-XY. Un revêtement conducteur en argent (3830, Holland Shielding Systems BV) est dilué avec un solvant industriel pour produire l'encre métallique qui est introduite dans la buse de pulvérisation. Le revêtement utilisé est composé de particules d'argent de taille micrométrique (25-50% du volume) mélangées à une résine thermoplastique (4-Methylpentan-2-one, moins de 60% du volume). Dans les travaux actuels, une partie du revêtement est diluée avec deux parties du solvant industriel, la Méthyl-éthyl-cétone (MEK). Cette encre métallique est ensuite pulvérisée dans l'espace négatif généré par le masque. Lorsqu'il est pulvérisé, le solvant s'évapore, déposant les particules d'argent sur le substrat diélectrique. Le revêtement est laissé à sécher pendant quelques minutes avant que le masque ne soit retiré. Le substrat diélectrique est immédiatement nettoyé avec de l'alcool éthylique pour éliminer les particules d'argent et les taches d'adhésif du masque. L'épaisseur des électrodes est contrôlée par le débit d'encre et la vitesse d'avance de la tête de l'outil de pulvérisation, ainsi que par le nombre de passages effectués.

Les caractéristiques de surface des électrodes en argent fabriquées avec cette nouvelle technique ont été étudiées par interférométrie. Deux actionneurs sont préférentiellement étudiés. Le premier actionneur possède une électrode de surface fabriquée en un seul passage de la buse de pulvérisation. Son épaisseur attendue d'environ 6 μ m (actionneur Ag_{6 μ m}). Le second actionneur présente une électrode réalisée par cinq passages du spray et l'épaisseur finale est d'environ 30 µm (actionneur Ag_{30µm}). Ces actionneurs offrent une circonstance favorable pour étudier l'influence de l'épaisseur de l'électrode exposée à l'air sur la vitesse dans le champ d'écoulement induit. Un interféromètre à lumière blanche Talysurf CCI 6000 a été utilisé afin de qualifier la qualité de la réalisation. Un champ de visualisation de $3 \times 3 \text{ mm}^2$ dans le plan $x_p - z_p$ centré sur le bord des électrodes exposées à l'air a été capturé avec ce dispositif pour deux actionneurs à électrodes d'argent conductrices. Les profils de surface acquis avaient une résolution spatiale finale de 1.8 μ m dans les deux directions x_p et z_p . Ces résultats sont présentés sur la figure 16, les points 16a et c correspondant à l'électrode fabriquée avec un passage de la buse de pulvérisation et les points 16b et d correspondant à l'électrode fabriquée après cinq passages de la buse de pulvérisation. Il convient de noter que ces actionneurs ont fonctionnés avant l'obtention de ces caractéristiques de surface.

L'échelle de couleur des cartes de contour des figures 16a et b représente la hauteur de la surface la plus haute à chaque endroit du champ de mesure par rapport à la surface diélectrique.



Figure 16.: Surface features of the electrodes fabricated with the novel technique obtained through interferometry. Coloured contour levels (in μ m) in (a) and (b) represent height of the top-most surface from the dielectric substrate of the electrodes produced with one and five passes of the spraying nozzle. Surface features of the same in three-dimensions (colour scales in μ m) are presented in (c) and (d).

Bien que les électrodes fabriquées semblent continues et lisses à l'œil, ces courbes de niveau indiquent qu'elles présentent des irrégularités à l'échelle micrométrique. Cela apparaît plus clairement dans les graphiques de surface tridimensionnels présentés dans les figures 16c et d. Cependant, un inconvénient de l'utilisation de l'interférométrie en lumière blanche pour les besoins actuels doit être souligné avant d'aller plus loin. L'argent est un matériau réfléchissant, surtout en lumière blanche. En conséquence, un certain nombre de pics de haute amplitude sont observés sur les deux graphiques de surface. Ces pics ne sont donc pas censés être représentatifs d'une caractéristique géométrique particulière, mais plutôt dus à la lumière de haute intensité provenant de la surface réfléchissante de l'électrode.



Figure 17.: Time-averaged velocity magnitude fields (in m s⁻¹). $V_{ac} = 16$ kV, $f_{ac} = 2$ kHz. The two grey bars below $y_p = 0$ represents the electrodes of the DBD actuator. The white solid (Ag_{6µm}) and dash-dotted (Ag_{30µm}) curves represent $u_{p,max}$ at each x_p location.

4.2.2. Caractérisation des performances mécaniques

Bien que les bords des électrodes ne soient pas parfaits, ces actionneurs offrent une belle opportunité d'étudier l'influence de l'épaisseur de l'électrode exposée à l'air sur la vitesse de l'écoulement pouvant être créé. Cet exercice nécessite des mesures du champ de vitesse avec une haute résolution spatiale pour saisir les potentielles différences dans la distribution de la vitesse. Pour cela, une campagne d'essai par PIV haute-résolution a été menée. Quelques-uns des résultats obtenus sont présentés ici.

Les actionneurs produisent un jet à paroi mince d'environ 0.75 mm d'épaisseur (voir figure 17). Un faible effet d'aspiration est détecté en amont du bord de l'électrode exposée à l'air $(x_p = 0)$. On observe que l'actuateur $Ag_{6\mu m}$ produit un jet comparativement plus fin. L'emplacement de la vitesse u maximale $(u_{p,max})$ est également plus proche de la paroi pour cet actionneur. Pour mieux comprendre cela, l'évolution de $u_{p,max}$ dans le vent ionique le long de x_p pour les deux actionneurs est présentée sur la figure 18a (courbes bleue et rouge sur les figures 17a et b). On voit que le $u_{p,max}$ se déplace vers la paroi, jusqu'à $x_p \approx 2$ mm. Ensuite, il commence à s'en éloigner. Il est clairement visible que le maximum de vitesse $u_{p,max}$ est plus proche de la paroi pour le vent ionique généré par l'actionneur $Ag_{6\mu m}$ comparé à l'actionneur $Ag_{30\mu m}$, ceci pour l'ensemble du champ moyen. Pour l'actionneur $Ag_{6\mu m}$, à $x_p = 2$ mm, le $u_{p,max} = 3,25$ m s⁻¹ et à $y_p \approx 0.13$ mm. À cet endroit, il est inférieur d'environ 125 μ m à la valeur $u_{p,max}$ de l'actionneur $Ag_{6\mu m}$.

La figure 18b présente le changement de position du $u_{p,max}$ avec l'épaisseur de l'électrode exposée à l'air, à $x_p = 2$ mm pour cinq amplitudes de tension. Dans tous les cas testés, une tendance similaire est observée. Cela confirme que l'épaisseur de l'électrode exposée à l'air a une influence sur la topologie du vent ionique produit. On peut s'attendre à ce qu'une électrode plus fine génère un champ électrique local plus fort à proximité du bord de l'électrode. Cela entraîne la génération et l'accélération d'espèces chargées plus près de la paroi. Ainsi, on observe que l'électrode plus fine produit un jet plus proche de la paroi mais la vitesse de ce jet est réduite par la présence d'une forte dissipation visqueuse dans cette région.

La poussée générée par les actionneurs de la DBD est calculée à partir des champs vectoriels moyens de la PIV en utilisant l'approche de bilan intégrale de quantité de mouvement. Le volume



Figure 18.: (a) Evolution of maximum u_p -velocity along x_p direction (from figure 17). (b) y_p -location of maximum u_p -velocity at $x_p = 2$ mm (dashed black line in (a)) versus voltage.



Figure 19.: Thrust (T_x) generated versus voltage amplitude (a) and frequency (b) for the three DBD actuators.

de contrôle choisi est l'ensemble du champ vectoriel moyen mesuré de manière à minimiser les effets de la pression. La poussée générée (T_x) en fonction de l'amplitude et de la fréquence de la tension appliquée est présentée sur la figure 19. La poussée générée par l'actionneur Cu_{40µm} (actionneur de référence construit à partir d'adhésif de cuivre d'épaisseur 40 µm) augmente avec la tension appliquée, comme indiqué précédemment (Kotsonis et al., 2011; Debien, Benard, & Moreau, 2012). Pour les actionneurs avec électrode argentée exposée à l'air, l'augmentation de la poussée est quasi-linéaire dans la plage de tension testée jusqu'à $V_{ac} = 18$ kV. Cependant, la courbe diminue légèrement à $V_{ac} = 20$ kV, ce qui suggère le début d'une décharge atteignant un régime filamentaire. Ceci se confirme avec les résultats de up,max observés précédemment. On observe que la poussée augmente linéairement avec la fréquence. La courbe montre une petite réduction pour l'actionneur Ag_{6µm} après $f_{ac} = 2$ kHz qui est à nouveau attribuée à une décharge devenant filamentaire.

4.3. Actionneurs DBD à modulation spatiale

Dans cette thèse, une comparaison est faite entre l'effet du forçage d'actionneur plasma DBD modulé spatialement et d'actionneur uniforme, ceci pour les deux configurations d'écoulement considérées. Alors que l'influence des paramètres du signal d'entrée (en particulier l'amplitude de la tension) sur les performances électriques et mécaniques est bien comprise pour les actionneurs uniformes, on sait peu de choses sur les dispositifs modulés spatialement. Des travaux de recherche antérieurs (voir par exemple Dörr & Kloker, 2017; Serpieri, Yadala, et al., 2017; Guo & Kloker, 2020) ont considéré que les performances des actionneurs modulés sont similaires à celles d'un actionneur correspondant à longueur uniforme. Dans l'étude actuelle où une comparaison des deux types de forçage est effectuée, la performance des actionneurs modulés en portée est indispensable pour mieux quantifier l'efficacité du contrôle. La première nécessité à cet égard est de démontrer la tridimensionnalité de la décharge plasma, puis il s'agira de mesurer l'écoulement produit afin de vérifier que la condition de forçage est belle et bien tridimensionnelle. Enfin, les performances de ces actionneurs modulés en espace seront comparées à un actionneur traditionnel à bord linéaire.

4.3.1. Caractérisation optique des régions ionisées

Tout d'abord, les caractéristiques de la décharge ont été étudiées en mesurant le signal de courant et en capturant la décharge elle-même avec une caméra iCCD. Les décharges pendant les moitiés positives et négatives du cycle du courant alternatif ont été capturées séparément. Pour chaque demi-cycle, deux champs avec des largeurs de porte d'ouverture différentes de la caméra iCCD ont été acquis. Les deux premiers champs (A et B) ont été capturés avec un temps d'ouverture de 14 μ s, tandis que les deux autres (C et D) ont été acquis avec une d'ouverture de 140 μ s. La figure 20 présente le signal de courant mesuré et les images de la décharge de plasma lorsque l'actionneur fonctionne pour $V_{ac} = 3$ kV et $f_{ac} = 2$ kHz. On voit que le courant de décharge a deux régimes, l'un pendant la moitié positive du signal d'entrée en courant alternatif et l'autre pendant la moitié négative. Comme observé dans les travaux de recherche précédents, le signal est composé de pics positifs de courant élevé pendant la moitié positive du cycle AC. Pendant la moitié négative, alors que les pics ont une polarité négative.

Dans ces deux régimes, la décharge est limitée aux régions de l'envergure entre les lignes vertes en pointillés $w_S = 3$ mm espacées les unes des autres, qui représentent l'interface entre l'air exposé et les tiges de l'électrode mise à la terre. Cependant, quelques canaux conducteurs sont également observés juste à l'extérieur de ces régions. Les champs C et D, où la largeur d'ouverture de la caméra iCCD est portée à 140 μ s, confirment l'apparition de cette décharge "secondaire". Cette décharge est principalement attribuée à deux phénomènes. Le premier est un faible champ électrique généré entre le bord générateur de plasma de l'électrode exposée à l'air et les bords longitudinaux des tiges (c'est-à-dire l_S dans la figure 15b). Le second est dû au dépôt de charge sur la surface diélectrique au cours des cycles de courant alternatif suivants. La longueur (moyenne) de cette décharge secondaire dans le sens de l'envergure, comparée à la largeur de chaque tige, est d'environ $0.4 \cdot w_S$. Il faut étudier l'effet que cela aura sur les champs de flux induits. Néanmoins, cette configuration géométrique permet d'obtenir une bonne modulation de la couche de plasma dans le sens de l'envergure. En tenant compte de cette décharge secondaire, seul 46% de l'envergure est ionisé. Enfin, dans le sens du courant, la décharge s'étend jusqu'à environ $x_p = 2$ mm.

4.3.2. Caractérisation mécanique par PIV tomographique

La section précédente a montré que la démarche adoptée permet de moduler spatialement la décharge plasma. L'objectif est désormais de vérifier que l'écoulement produit est lui aussi localisé en certaines régions le long du bord de l'électrode, les régions ionisées par le champ électrique. Il s'agit donc de mettre en place une métrologie permettant d'accéder simultanément



Figure 20.: Discharge current and fast imaging of the developing plasma layer for the spanwise-modulated DBD actuator configuration. The employed actuator is operated at $V_{ac} = 3$ kV and $f_{ac} = 2kHz$. Dashedgreen lines in A, B, C and D represent the stems of the grounded electrode. A & B: 14 μ s opening gate width; C & D: 140 μ s opening gate width.

aux trois composantes de la vitesse. À cet effet, le dispositif de PIV tomographique présenté précédemment a été utilisé pour mesurer les champs de vitesse induits. Le champs de mesure était telle que l'organisation de l'écoulement le long de deux tiges adjacentes de l'électrode mise à la terre (distance de $\lambda_z = 9$ mm) peut être mesuré. Les champs de vitesse moyens dans le temps ($U_p = \sqrt{u_p^2 + v_p^2 + w_p^2}$) de l'écoulement induit par l'actionneur plasma à modulation d'amplitude sont présentés sur la figure 21.

Le champ de vitesse induit moyenné dans le temps lorsque l'actionneur de plasma DBD fonctionne à $V_{ac} = 3$ kV et $f_{ac} = 2$ kHz est présenté dans la figure 21a. Ce champ indique clairement la nature tridimensionnelle du champ d'écoulement produit. Deux jets sont observés, chacun prenant naissance à l'interface de l'électrode exposée à l'air et d'une tige de l'électrode mise à la terre. Une faible vitesse est observée juste en amont de cette interface. C'est la faible aspiration qui alimente le fluide dans les jets créés. On observe que l'actionneur n'accélère le flux qu'aux endroits où les électrodes se chevauchent et conduit à la formation de jets pariétaux sur la longueur des tiges. Ainsi, il en résulte un champ d'écoulement qui est modulé dans le sens de l'envergure en fonction de la distance entre les tiges adjacentes, qui est de $\lambda_z = 9$ mm dans ce cas. Cela indique en outre que la force volumique générée par ces actionneurs sera également modulée dans le sens de l'envergure (tridimensionnelle). En outre, la décharge secondaire à proximité des tiges, observée sur les images iCCD de la couche de plasma de la figure 20, ne





Figure 21.: Time-averaged velocity magnitude field (m s^{-1}) of the flow induced by the spanwise-modulated plasma actuator captured with tomographic PIV.

semble pas avoir d'impact significatif sur le champ d'écoulement induit, car les jets générés sont principalement limités à la largeur des tiges.

Le flux induit par l'actionneur lorsque la fréquence du signal AC d'entrée est augmentée à $f_{ac} = 6$ kHz est montré dans figure 21b. Cette augmentation de la fréquence du courant alternatif entraînera une augmentation de la puissance fournie à l'actionneur, car ces paramètres varient de manière linéaire les uns par rapport aux autres. L'écoulement produit est toujours tridimensionnel et est modulé en fonction de l'espacement entre les tiges ($\lambda_z = 9$ mm). Dans cette condition de fonctionnement avec une puissance d'entrée plus élevée, un deuxième noyau de jet plus petit est observé dans le jet induit centré autour de $z_p = 9$ mm, au bord de la tige ($z_p \approx 11$ mm). La raison de cette différence entre les deux jets n'est pas connue, mais elle pourrait être attribuée à une décharge secondaire plus forte au voisinage de la tige à $z_p = 9$ mm. En outre, le jet induit est principalement limité par la largeur de la tige. Le faible écoulement induit à l'extérieur de la tige à $z_p = 0$ a la même magnitude de vitesse ($U_p \approx 0.1$ m s⁻¹) que dans le champ induit précédent montré sur la figure 21a.

4.3.3. Performances électromécaniques

Les sections précédentes ont montré que la décharge plasma était modulée spatialement par le design proposé et que l'écoulement produit était lui aussi modulé comme souhaité. Pour une caractérisation approfondie des performances de ces actionneurs, en particulier pour mesurer la force générée qui sera utilisée pour les applications de contrôle d'écoulement, l'utilisation de mesures de PIV tomographique n'est pas très pratique. Le coût de calcul pour obtenir les champs vectoriels de vitesse, la quantité de données générées et les limites de la VFO rendent l'utilisation de ces techniques compliquée dès lors qu'une grande matrice d'essai est nécessaire pour caractériser en détail l'actionneur. Bien qu'un certain nombre de paramètres puissent être mesurés pour déterminer la performance de l'actionneur, deux paramètres, à savoir la puissance consommée par l'actionneur et la poussée correspondante générée, sont choisis pour quantifier la performance des actionneurs modulés spatialement. Cinq actionneurs à plasma DBD différents ont été fabriqués à cet effet. Ces actionneurs ont tous été construits sur des plaques de PMMA de 2 mm d'épaisseur qui ont servi de diélectrique. Le choix de ce diélectrique est motivé par le fait que les actionneurs à portée uniforme construits sur ce diélectrique sont très robustes et que leurs performances sont bien documentées dans la littérature (voir par exemple Jolibois & Moreau, 2009). Les électrodes ont été fabriquées avec des bandes d'aluminium recouvertes d'un revêtement noir non réfléchissant (épaisseur totale de 85 μ m).

Tout d'abord, les mesures concernant l'actionneur de référence (bord d'électrode rectiligne et plasma uniforme), les résultats acquis montrent que la variation de la consommation d'énergie avec l'amplitude de la tension d'entrée en courant alternatif est bien approchée par une courbe $P_{avg} = A_P (V_{ac} - V_0)^2$ (courbe noire continue dans la figure 22a). La tension d'allumage de la décharge plasma est de $V_0 = 4$ kV et le coefficient constant est $A_P = 6.82 \cdot 10^{-8}$ W V⁻² (ajustement effectué sur les données mesurées, avant normalisation). Il est à noter que cet ajustement n'est valable qu'après la tension d'allumage du plasma (V_0). Ainsi, le coefficient constant obtenu ici est comparable à celui rapporté par Pons et al. (2005).

La puissance consommée par les actionneurs modulés spatialement, normalisée avec la longueur de décharge correspondante (l_p) est également présentée dans la figure 22a. On peut voir que la puissance normalisée de ces actionneurs est très similaire à celle de l'actionneur de référence. Ici aussi, la puissance consommée augmente selon la même loi de puissance $P_{avg} = A_P (V_{ac} - V_0)^2$.



Figure 22.: Quantification of the spanwise-modulated-actuator performance. (a) and (b) show the average power consumed (P_{avg}) and thrust generated (T_x) along the wall-parallel direction respectively, as a function of voltage amplitude V_{ac} ($f_{ac} = 2$ kHz). Wall-parallel thrust generated as a function of average power consumed are presented in (c). Solid black curves in (a) and (b) are the power-law fit on the corresponding data of the spanwise-uniform (2D) actuator.

La poussée parallèle à la surface du diélectrique (T_x) , moyennée dans le temps, produite par l'actionneur DBD de référence est présentée dans la figure 22b. Il convient de noter que les valeurs mesurées ne sont pas la force totale générée par l'actionneur DBD, mais la force EFD soustraite des forces de frottement encourues en raison de la présence du support diélectrique. Comme pour la consommation d'énergie, la force générée par l'actionneur de référence augmente avec l'amplitude de la tension appliquée selon une loi de puissance (Thomas, Corke, et al., 2009; Kotsonis et al., 2011). Après la tension d'allumage du plasma, elle est bien approchée par la relation $T_X \propto (V_{ac} - V_0)^2$, comme observé par Debien, Benard, & Moreau (2012). Dans le cas des actionneurs modulés spatialement, la force générée est comparable à celle de l'actionneur de référence lorsque la mesure est normalisée avec la longueur de décharge, comme observé précédemment avec la consommation électrique.

On peut conclure que la modulation spatiale de la décharge n'a que peu d'influence sur les régimes du plasma et les caractéristiques de l'ionisation. En fait, les performances de ces actionneurs modulés spatialement s'échelonnent en fonction de la longueur de la décharge. Cela indique que, ici, la modification de la géométrie n'a pas d'effet significatif sur le plasma lui-même (en termes de physique du plasma). Ainsi, les résultats démontrent que la performance des actionneurs modulés spatialement peut être estimée avec succès en connaissant la performance de l'actionneur bidimensionnel uniforme correspondant. Ceci est particulièrement utile car la fabrication et la mesure des paramètres de performance des actionneurs bidimensionnels sont à la fois simples et rapides par rapport aux actionneurs modulés en espace.

5. COUCHE LIMITE TRIDIMENSIONELLE SUR UNE AILE EN FLÈCHE

Les résultats des mesures de contrôle sur la couche limite d'une aile en flèche sont résumé dans ce chapitre. Les avions à réaction modernes ont des ailes dont les bords d'attaque sont en flèche afin d'augmenter le nombre de Mach critique. Cette géométrie donne naissance à une couche limite laminaire tridimensionnelle dont la composante de vitesse transverse forme intrinsèquement un point d'inflexion, ce qui entraîne l'instabilité de l'écoulement transversal non visqueux, qui se caractérise par alors par des tourbillons en rotation de nature stationnaire (stable) ou mobile (instable), en fonction des perturbations extérieures (Bippes, 1999; Downs & White, 2013). Ici, les actionneurs DBD sont déployés pour réaliser deux stratégies distinctes afin de contrôler les instabilités et la transition éventuelle de la couche limite vers la turbulence. Ces méthodes sont la déformation en amont du flux (UFD) et la modification du flux de base (BFM). La stratégie UFD sera réalisée avec un actionneur modulé en espace, tandis que la stratégie BFM sera implémentée en utilisant un forçage uniforme selon l'envergure.

L'objectif est également de déterminer l'efficacité de cette méthode contrôle par plasma pour retarder la transition laminaire-turbulente. En plus d'être la première étude expérimentale de l'efficacité de la stratégie BFM, la première comparaison de cette méthode de contrôle avec la stratégie UFD dans des conditions d'écoulement expérimentales identiques est également effectuée ici. Si cela facilite la comparaison des deux stratégies en termes d'efficacité et d'efficience, cela permet également d'examiner la collaboration (le cas échéant) entre les stratégies de contrôle UFD et BFM.

5.1. Conditions expérimentales des essais

5.1.1. Régimes étudiés

Les études expérimentales sur la couche limite tridimensionnelle ont été menées dans la soufflerie à faible turbulence de l'Université Technique de Delft, aux Pays-Bas. Le modèle utilisé est une aile en flèche à 45° appelée *M3J*. La corde de ce modèle c = 1.27 m est utilisée dans ce chapitre pour adimensionner les coordonnées spatiales. Pour l'ensemble des essais l'angle d'incidence de l'aile a été fixé à $\alpha = 3^{\circ}$. Pour la plupart des cas testés, la vitesse de l'écoulement à l'infini était de $U_{\infty} = 25$ m s⁻¹ ($M_{\infty} = 0.07$), ce qui correspond à un nombre de Reynold de $Re_c = 2.1 \cdot 10^6$. Un certain nombre de cas ont également été acquis à $U_{\infty} = 27.5$ m s⁻¹ ($Re_c = 2.3 \cdot 10^6$) et $U_{\infty} = 30$ m s⁻¹ ($Re_c = 2.5 \cdot 10^6$).

5.1.2. Conditionnement de l'écoulement par éléments rugueux

Une analyse par stabilité linéaire est utilisée pour déterminer les caractéristiques de stabilité de la couche limite étudiée, en particulier la longueur d'onde du mode d'instabilité transverse stationnaire critique dans les conditions d'écoulement étudiées de $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$. Les variables de stabilité ont été calculées selon la procédure introduite précédemment. La solution de la couche limite laminaire du côté pression de l'aile balayée a été calculée à partir de la distribution de pression correspondante acquise expérimentalement. La solution de la couche limite été utilisée pour calculer et analyser les caractéristiques de



Figure 23.: N-factors of different stationary CF instability modes (f = 0 Hz) of the laminar boundary layer generated on the pressure side of the swept-wing at flow conditions $U_{\infty} = 25$ m s⁻¹ ($Re_c = 2.5 \cdot 10^6$) and $\alpha = 3^\circ$.



Figure 24.: IR thermography mean fields. The flow comes from left. (a) Boundary-layer transition topology of the clean wing. (b) Boundary-layer transition topology with an array of DREs installed at x/c = 0.017, forcing the $\lambda_{z,crit} = 8$ mm stationary CF instability mode. In both figures, PET foils are mounted on the swept-wing model but no DBD plasma actuator is fabricated on the top most foil.

stabilité de la couche limite. Les facteurs d'amplification N dans le sens de la corde de certains modes d'instabilité transverse stationnaires dans la couche limite laminaire ont été calculés et sont présentés sur la figure 23. À l'emplacement le plus en aval (proche des minima de pression), l'instabilité la plus dominante a une longueur d'onde de z = 8 mm. Ainsi, dans la couche limite actuelle, l'instabilité transverse stationnaire critique a une longueur d'onde de $\lambda_{z,crit} = 8$ mm.

Dans les conditions d'écoulement étudiées ($Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$), la transition naturelle sur le modèle utilisé se situe en aval du point minimum de pression de x/c 0.63. Ainsi, le processus de transition est non seulement dominé par le mécanisme d'instabilité transverse mais aussi par la séparation laminaire comme le montre la figure 24a. Ainsi, le mode d'instabilité transverse stationnaire critique a été forcé en utilisant un ensemble d'éléments de rugosité discrets (DRE). Le mode d'instabilité transverse stationnaire le plus instable ou le plus critique dans la couche limite étudiée a été identifié, par analyse LST, comme ayant une longueur d'onde de $\lambda_{z,crit} = 8 \text{ mm}$. En outre, la LST prédit que le début (point neutre) de ce mode se situe à x/c = 0.0068 et montre une croissance linéaire jusqu'à la station la plus en aval dans le sens de la corde. Ce mode d'instabilité critique a été forcé avec le réseau DRE. Le dispositif était composé d'éléments cylindriques d'un diamètre de $d_{DRE} = 2 \text{ mm}$ et d'une hauteur de $k_{DRE} = 60 \mu$ m, de sorte que la distance entre les DRE adjacents était de ($\lambda_{z,crit} =) 8 \text{ mm}$.

Le réseau de DRE a été installé à x/c = 0.017, juste en aval du point neutre du mode d'instabilité critique. Le rapport entre la hauteur du DRE et l'épaisseur du déplacement de la couche limite locale est $k_{DRE}/\delta^* = 0.47$. Le résultat de ce forçage sur la couche limite et sa transition éventuelle est visible sur la figure 24b. Le forçage passif par DRE permet de ramener la transition en amont à environ $x_{tr}/c 0.32$ et d'obtenir une transition laminaire-turbulente dominée par l'instabilité transverse. Un autre facteur de la plus haute importance est que l'amplification du mode critique présente un "scénario du pire" pour tout schéma de contrôle qui serait appliqué dans ces conditions. Ces conditions de couche limite conditionnée par un réseau de DRE à une longueur d'onde de $\lambda_z = 8$ mm est considérée comme le cas de référence des études présentées par la suite.

5.1.3. Conditions de forçage

L'installation des actionneurs plasma le modèle d'aile en flèche est brièvement décrite ici. Il est absolument nécessaire de s'assurer qu'en cas de panne de l'un des actionneurs DBD utilisés, la surface du modèle d'aile ne sera pas endommagée. Ainsi, deux feuilles de PET de 500 µm d'épaisseur ont été installées sur la surface du modèle. La feuille supérieure était fixée sur la feuille inférieure (qui était plus grande) de telle sorte que leurs bords inférieurs étaient alignés l'un avec l'autre. La feuille supérieure était géométriquement et dimensionnellement similaire à la feuille PET sur laquelle les actionneurs plasma du DBD ont été fabriqués par la méthode présentée précédemment.

Afin de comparer l'autorité de contrôle des actionneurs pour différentes configurations et conditions de forçage, il est nécessaire de normaliser la force appliquée par l'actionneur et d'obtenir une estimation de son amplitude par rapport à la couche limite expérimentale. À cet effet, les valeurs de poussée intégrées T_x sont utilisées pour calculer le coefficient de moment cinétique c_{μ} selon l'équation 9,

$$c_{\mu} = \frac{T_x}{\frac{1}{2}\rho u_e^2 \theta_e} \cdot \frac{w_S}{\lambda_z} \tag{9}$$

où, u_e et θ_e sont respectivement la vitesse au bord de la couche limite et l'épaisseur du moment cinétique, à l'endroit où le forçage est appliqué par l'actionneur (x/c = 0.035). Ces données sont extraites de la solution de la couche limite laminaire dans les conditions d'écoulement étudiées. La poussée générée, calculée à partir des mesures de vitesse par PIV et l'application d'un bilan intégrale des quantités de mouvement, est utilisée pour calculer c_{μ} . Un rapport supplémentaire w_S/λ_z , qui représente le rapport-cyclique des actionneurs lorsqu'ils sont modulés spatialement, fait ici partie de l'équation. Ceci est fait pour tenir compte de la tridimensionnalité du forçage appliqué. Ainsi pour un forçage uniforme, ce rapport est égal à 1.

La figure 25a présente la variation de cµ pour le forçage uniforme selon l'amplitude de la tension d'entrée, ceci pour trois conditions d'écoulement ($Re_c = 2.1 \cdot 10^6$, $2.3 \cdot 10^6$ et $2.5 \cdot 10^6$). Le coefficient c_{μ} augmente avec l'amplitude de la tension d'entrée. Cependant, on observe qu'il diminue avec l'augmentation de la vitesse de l'écoulement. Ceci est prévisible car c_{μ} est inversement proportionnel au carré de la vitesse de bord de la couche limite (u_e) qui augmente avec la vitesse de l'écoulement à l'infini. La figure 25b présente la variation de cµ pour les configurations de forçage modulées dans l'espace à $Re_c = 2.1 \cdot 10^6$. Les coefficients c_{μ} de ces actionneurs sont nettement inférieurs à ceux de l'actionneur uniforme en raison du rapport supplémentaire w_S/λ_z introduit dans l'équation 9. Dans cette thèse, la multitude de cas étudiés fournit une large gamme de valeurs c_{μ} . En considérant uniquement les conditions d'écoulement à $Re_c = 2.1 \cdot 10^6$, la valeur testée la plus basse est $c_{\mu} = 0.01$ (actionneur de $\lambda_z = 5$ mm à


Figure 25.: Momentum coefficient versus applied AC voltage amplitude, the AC frequency being $f_{ac} = 10$ kHz. (a) Spanwise-uniform EFD forcing at three different Reynold's number. (b) Spanwise-modulated EFD forcing at $Re_c = 2.1 \cdot 10^6$.

 $V_{ac} = 2.5 \text{ kV}$) et la valeur testée la plus élevée est $c_{\mu} = 0.71$ (actionneur uniforme dans l'espace $V_{ac} = 4.5 \text{ kV}$). Désormais, cµ sera utilisé pour discuter des différentes conditions de forçage.

5.2. Contrôle des instabilités transverses par actionneur DBD

Dans cette étude, l'utilisation d'un réseau de DRE ne conditionne pas seulement la couche limite en verrouillant le mode d'instabilité transverse stationnaire critique $\lambda_{z,crit} = 8$ mm, mais le réseau fixe également la position des tourbillons transverse dans le sens de l'envergure le long de la ligne d'écoulement non visqueuse. Ce faisant, il fournit une condition expérimentale idéale pour étudier comment les stratégies UFD et BFM réalisées avec le forçage des actionneurs plasma DBD collaborent ou s'opposent l'une à l'autre.

5.2.1. Contrôle par approche UFD

Les actionneurs DBD à décharge modulé dans l'espace sont maintenant utilisés pour forcer des modes d'instabilité transverse stationnaires spécifiques, afin de réaliser la stratégie de contrôle de l'UFD. Il convient de noter que tous les cas d'essais présentés dans cette section correspondent aux conditions d'écoulement à nombre de Reynold de $Re_c = 2.1 \cdot 10^6$ ($U_{\infty} = 25 \text{ m s}^{-1}$). Seule cette condition est étudiée car le mode critique varie avec le nombre de Reynold, et donc si les conditions d'écoulement sont modifiées, le mode stationnaire de 8 mm forcé par le réseau de DRE ne restera pas le mode critique. Deux actionneurs DBD qui effectuent le contrôle de l'UFD en générant des jets modulés dans le sens de l'envergure pour forcer un seul mode d'instabilité transverse stationnaire monochromatique ont été conçus et fabriqués. Le premier actionneur a été conçu pour forcer le mode d'instabilité sous-critique de λ_z ($= \lambda_{z,crit}/2$) = 4 mm. Le second actionneur force le mode d'instabilité de λ_z ($\lambda_{z,crit} \cdot 2/3$) = 5 mm.

Avec l'actionneur DBD forçant le mode d'instabilité sous-critique avec la longueur d'onde $\lambda_z = 4$ mm, en manipulant la position relative dans le sens de l'envergure de l'actionneur DBD par rapport aux DRE, deux configurations de forçage différentes ont été réalisées. Dans la première, l'actionneur DBD a été installé de telle sorte que la ligne d'écoulement non-visqueuse entrante passe au-dessus de la position x/c = 0.035, à l'interface des électrodes où la force EFD est exercée (voir figure 26a). Comme les tourbillons transverses stationnaires sont approxima-



Figure 26.: Schematic of the forcing configurations to realize the UFD control strategy. The DRE arrangement (•) and the DBD actuator with its encapsulated electrode (orange, with stems) and airexposed electrode (grey bar) are depicted. Consecutive inviscid streamlines originating from the DRE are also shown. (a) UFD4_{in} configuration; (b) UFD4_{out} configuration; (c) UFD5 configuration.

tivement alignés avec l'écoulement produit par l'actionneur, les jets de plasma sont "en phase" avec les tourbillons CF entrants (ci-après dénommés UFD4_{in}). Dans la seconde configuration, l'actionneur DBD est déplacé de 2 mm selon l'envergure (direction z), ce qui fait que les jets produit par la décharge sont maintenant "déphasés" par rapport aux tourbillons entrants (configuration UFD4_{out}). Dans ce cas, les tourbillons transverses stationnaires passent entre deux tiges adjacentes de l'électrode encapsulée (lieu de création du ven électrique) comme le montre la figure 26b. Le deuxième actionneur DBD forçant les modes d'instabilité transverse stationnaire sous-critique de $\lambda_z = 5$ mm fournit une troisième configuration de forçage (figure 26c, ci-après dénommée configuration UFD5). Cette configuration est telle qu'un tourbillon transverse sur cinq passe directement sur l'interface des électrodes à x/c = 0.035 (comme dans la configuration UFD4_{in}). Alors que les quatre autres tourbillons passent la station x/c = 0.035 dans le sens de la corde à des distances variables des zones de décharge ce qui entraîne une relation de phase variable dans l'espace entre les deux.

La figure 27 présente la topologie d'écoulement et de transition pour $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$, visualisée à l'aide de la thermographie IR. Pour les cas de référence (pas de forçage EFD) présentés aux figures 27a, d et g, les actionneurs DBD correspondant respectivement aux configurations de forçage UFD4_{in}, UFD4_{out} et UFD5, sont montés sur le modèle mais ne sont pas en fonctionnement. Les zones claires de ces figures correspondent aux régions de l'aile où l'écoulement est en régime laminaire, tandis que les zones plus sombres correspondent à un écoulement turbulent. Dans tous les champs IR de base, la rupture laminaire se produit selon le modèle dentelé caractéristique (Saric, Reed, et al., 2003; Saric, Carpenter, et al., 2011; Downs & White, 2013; Serpieri & Kotsonis, 2016). En outre, la transition se produit approximativement à la même position, dans le sens de la corde, que celle où seul le réseau de DRE conditionne la couche limite. Cela indique que les électrodes exposées à l'air des actionneurs fabriqués selon la nouvelle technique n'ont aucune influence sur la couche limite et sa transition naturelle. Cela indique gue les de fabrication proposées ont réussi à atteindre l'un de leurs principaux objectifs: produire des électrodes d'une épaisseur de quelques micromètres qui n'influencent en rien la couche limite.

La topologie de l'écoulement résultant du forçage EFD appliqué par les actionneurs plasma DBD ($V_{ac} = 4 \text{ kV}$ et $f_{ac} = 10 \text{ kHz}$) est illustrée dans les figures 27b, e et h. Dans tous les cas, une



Figure 27.: IR thermography mean fields. DBD plasma actuators were operated at $V_{ac} = 4 \text{ kV}$, $f_{ac} = 10 \text{ kHz}$. (a) UFD4_{in} unforced flow; (b) $c_{\mu} = -0.068$; (c) Subtraction of (a) from (b); (d) UFD4_{out} unforced flow; (e) $c_{\mu} = -0.068$; (f) Subtraction of (d) from (e); (g) UFD5 unforced flow; (h) $c_{\mu} = -0.055$; (i) Subtraction of (g) from (h). The solid white lines represent constant chord positions.

bande brillante est observée dans le coin supérieur gauche. Il s'agit de la signature thermique de la décharge à barrière diélectrique. En outre, la topologie du front de transition permet d'évaluer plus en détail les caractéristiques des instabilités sous-jacentes. Plus précisément, un front de transition laminaire-turbulent uniforme sur toute l'étendue de la zone suggère que la transition est dominée par les modes d'instabilité mobiles. Lorsque ces modes instables se propagent dans une direction perpendiculaire à la ligne de flottaison non-visqueuse locale, ils distribuent ou "lissent" la forte modulation de vitesse dans le sens de l'envergure causée par les modes d'instabilité transverses stationnaires. In fine, cela érode le front de transition traditionnellement très 'déchiqueté' (Downs & White, 2013). Dans les visualisations IR présentées ici, le front de transition présente toujours un motif dentelé persistant lorsque le forçage EFD est appliqué, ce qui suggère que la transition vers la turbulence est toujours dominée par les modes d'instabilité transverse stationnaires.

Malgré la confirmation visuelle du retard de transition, évidente dans la figure 27, une estimation quantitative du lieu de transition est maintenant recherchée. Les déplacements du front de transition déterminés expérimentalement ($\Delta x_{tr}/c$) sont extraits des champs IR pour les configurations de forçage du plasma testé et différents coefficients de moment. Les résultats sont présentés dans la figure 28. On observe que la position moyenne du front de transition sur toute l'envergure se déplace vers l'aval avec cµ pour les trois configurations étudiées. Pour la configuration UFD4_{in}, le plus grand retard de transition obtenu est d'environ $\Delta x_{tr}/c = 2.9\%$ de la corde, se produisant pour le forçage le plus intense à $c_{\mu} = -0.068$. Pour la configuration de forçage UFD4_{out}, le plus grand retard de transition atteint est $\Delta x_{tr}/c = 1.7\%$, ce qui est nettement inférieur à celui obtenu avec la configuration précédente. Le plus grand déplacement du front de transition avec la configuration de forçage UFD5 est encore plus faible puisqu'il n'est



Figure 28.: Experimentally determined laminar-to-turbulent transition-front displacement $(\Delta x_{tr}/c)$ versus momentum coefficient c_{μ} for the three forcing configurations investigated here.

que d'environ $\Delta x_{tr}/c = 0.8\%$ ($c_{\mu} = -0.055$).

5.2.2. Contrôle par approche BFM

La discussion de la section précédente montre un retard dans la transition laminaire-turbulente grâce à l'application d'un forçage modulé dans le sens de l'envergure. Cependant, les résultats indiquent également qu'avec la configuration de l'actionneur modulé dans le sens de l'envergure, il n'y a aucune certitude concernant la réalisation du contrôle de l'instabilité. En outre, les résultats indiquent également que le retard de transition obtenu n'est pas uniquement dû à l'action UFD imposé par l'actionneur, mais une modification locale de l'écoulement de base (BFM) est aussi soupçonnée. Cela suggère qu'une action UFD pure ne peut pas être réalisée avec la géométrie actuelle de l'actionneur plasma. D'autre part, une stratégie de contrôle purement BFM peut être réalisée avec un actionneur qui applique un forçage EFD uniforme sur toute l'envergure. Les résultats y afférents, obtenus à la fois dans les domaines théorique et expérimental, sont présentés dans cette section.

Au voisinage de la ligne d'attache au bord d'attaque d'une aile en flèche, les lignes d'écoulement non-visqueux sont presque alignées avec la direction z. Ainsi, la composante transverse (w_{ISL}) de l'écoulement de base dans cette région est approximativement alignée avec la direction x, comme le montre la figure 29. L'application d'un forçage dans la direction x aura alors un impact important sur la composante transverse du flux de base et contribuera à la réalisation de la stratégie de contrôle BFM. Ceci est réalisé ici par l'application d'une force de corps EFD bidimensionnelle et uniforme dans l'envergure dans la direction x, c'est-à-dire orientée perpendiculairement au bord d'attaque. Cette conception est adoptée pour éviter le forçage indésirable des modes d'instabilité non désirés et assurer une configuration de forçage simple et robuste. À cet effet, deux actionneurs plasma DBD ont été fabriqués, l'un pour appliquer un forçage $-F_x$ (figure 29a) et l'autre pour appliquer un forçage F_x (figure 29b) sur la couche limite.

La topologie de l'écoulement, visualisée par thermographie IR, est présentée à la figure 5.12 pour $Re_c = 2.1 \cdot 10^6 \ (U_{\infty} = 25 \text{ m s}^{-1})$. Les images IR des cas non forcés sont présentées dans les figures 30a et d. Comme précédemment, les actionneurs sont installés sur l'aile mais ne fonctionnent pas ici non plus. La position de la transition reste approximativement similaires ce qui confirme à nouveau l'absence d'influence des électrodes fines. On observe également un



Figure 29.: Schematic of the forcing configurations to realize the BFM control strategy (not to scale). The swept-wing leading edge (thick black line), inviscid streamline (thin black line) and local velocity components (u_{ISL}, w_{ISL}) are shown. The DRE arrangement (•) and the DBD actuator's air-exposed electrode (grey bar) are also depicted. (a) $-F_x$ EFD plasma forcing (red arrows); (b) F_x EFD plasma forcing (blue arrows).

motif strié dans ces champs. L'analyse spectrale de ces images confirme que le mode dominant dans la couche limite est le mode $\lambda_{z,crit} = 8$ mm forcé par les DRE. Cela suggère que, comme prévu, le forçage EFD uniforme appliqué dans le sens de l'envergure n'a aucun effet sur les modes d'instabilité transverse dominants dans la couche limite. En outre, pour faciliter la visualisation du mouvement du front de transition, les figures 30c et f montrent la soustraction du champ IR avec le forçage EFD du cas non forcé correspondant. Le forçage F_x déplace évidemment le front de transition vers l'amont, par rapport au cas non forcé. Au contraire, le forçage $-F_x$, dans les mêmes conditions de forçage et d'écoulement, déplace le front de transition vers l'aval.

Les points de transition mesurés expérimentalement sont présentés à la figure 31, pour les cas d'un forçage selon $-F_x$ et F_x . Le retard ou l'avancement de la transition est proportionnel à l'amplitude de c_{μ} . Dans les cas testés, le plus grand retard de transition est d'environ 4.5%. Il est observé au plus petit nombre de Reynolds et au plus grand coefficient cµ testé ($Re_c =$ $2.1 \cdot 10^6$, $c_{\mu} = -0.71$). Pour les nombres de Reynolds plus élevés, le déplacement des fronts de transition avec cµ montre des tendances similaires. Pour $Re_c = 2.3 \cdot 10^6 (U_{\infty} = 27.5 \text{ m s}^{-1})$ \triangle), $c_{\mu} = -0.62$ et $Re_c = 2.5 \cdot 10^6 (U_{\infty} = 30 \text{ m s}^{-1}, \Box)$, $c_{\mu} = -0.55$ le délai de transition est respectivement de 3.7% et 1.6%. Ici, la technique de contrôle de l'écoulement est une approche directe visant à affecter directement la composante de vitesse transverse dont l'amplitude est faible. Cette composante augmente avec la vitesse de l'écoulement à l'infini et donc, l'autorité de la force du corps de l'EFD appliquée sur la couche limite diminue à mesure que le nombre de Reynolds augmente. En outre, le mode critique dans la couche limite change avec le nombre de Reynolds, mais on observe que le forçage appliqué a toujours un effet. Cela démontre une fois de plus l'insensibilité de la méthode UFD au nombre d'ondes des modes d'instabilité stationnaires transverses. Les barres d'incertitude de la figure 30 représentent la largeur du front de transition. La largeur du front de transition observée ici semble plutôt insensible au forçage. La présence d'un front de transition qui conserve une forme dentelée suggère qu'aucune fluctuation dans la bande de fréquence des instabilités transverses instationnaire n'a été introduite. Ceci, est une indication préliminaire que l'application du forçage EFD à des fréquences élevées n'entraînera pas l'amplification des fluctuations instables dans ces couches limites.

L'évolution des facteurs N pour le mode d'instabilité critique à $\lambda_{z,crit} = 8$ mm, calculée par le modèle simplifié, est présentée dans la figure 32. Si l'existence d'autres modes stationnaires ne peut être exclue, le mode 8 mm est le plus dominant et le plus observé pour entraîner la transition, car il est conditionné par les DRE. Par conséquent, le facteur N de ce seul mode est présenté et analysé ici. En utilisant les emplacements de transition obtenus à partir des mesures de thermographie IR, on observe que le facteur N critique auquel la transition se



Figure 30.: IR thermography mean fields $(Re_c = 2.1 \cdot 10^6, U_{\infty} = 25 \text{ m s}^{-1})$. (a) Unforced flow. (b) $-F_x$ $(c_{\mu} = -0.71)$. (c) Subtraction of (a) from (b). (d) Unforced flow. (e) F_x $(c_{\mu} = 0.71)$. (f) Subtraction of (d) from (e). The solid white lines represent constant chord positions.



Figure 31.: Experimentally measured transition locations (x_{tr}/c) versus momentum coefficient c_{μ} . \bigcirc : $Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$; \triangle : $Re_c = 2.3 \cdot 10^6$, $U_{\infty} = 27.5 \text{ m s}^{-1}$; \Box : $Re_c = 2.5 \cdot 10^6$, $U_{\infty} = 30 \text{ m}$ s⁻¹; \bullet : transition location predicted by simplified numerical model for $Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$.

produit pour le cas d'écoulement non forcé est $N_{crit} \approx 5.6$ (ligne en pointillé sur la figure 32). Conformément aux diagrammes de stabilité, bien que le forçage appliqué affecte l'amplification



Figure 32.: N-factors of the $\lambda_{z,crit} = 8$ mm stationary CF instability mode in unforced and forced numerical boundary layers. Chordwise locations of the DRE array (dotted line) and DBD plasma actuator (dashed line) are depicted. Experimentally obtained transition locations (\Box) are also shown. Unforced boundary layer $N_{crit} = 5.6$ (dash-dotted line).

des instabilités localement, son effet intégral se traduit par une réponse globale sur les facteurs N. La réduction des taux de croissance du mode d'instabilité avec le forçage $-F_x$ entraîne une diminution globale des facteurs N par rapport à un écoulement non forcé. Ainsi, la valeur critique Ncrit à laquelle la transition se produit naturellement dans le cas non forcé est atteinte en aval, retardant ainsi la transition. Au contraire, le forçage F_x augmente l'amplification des instabilités et entraîne une augmentation des amplitudes du facteur N. Cette augmentation des facteurs N a pour conséquence d'atteindre N_{crit} en amont du cas non forcé, favorisant ainsi la transition la finaire-turbulente. Les emplacements de transition détectés expérimentalement des cas d'écoulement forcé sont également indiqués dans la figure 32, ce qui indique la capacité de l'approche de modélisation simplifiée à saisir les tendances générales de la manipulation de la transition avec forçage. Il est en effet évident que le modèle est capable de saisir la tendance expérimentale, mais qu'il ne parvient pas à prédire avec précision le lieu de transition. Cela est prévisible, compte tenu de l'éventail des hypothèses et des simplifications qui entrent dans la chaîne des étapes de modélisation. En dépit des réserves émises, l'approche de modélisation simplifiée présente un outil robuste, rapide et efficace, idéalement adapté pour compléter le processus de conception expérimentale et fournir une mise à l'échelle de base des performances de contrôle potentiellement atteignables.

5.2.3. Comparaison des rendements et gains énergétiques

Le gain net réalisé grâce à la manipulation de l'écoulement de couche limite appliquée est maintenant évalué. Il s'agit du rapport entre la puissance économisée grâce à la réduction de la traînée de frottement et la puissance fournie à l'actionneur DBD normalisée avec l'étendue de l'actionneur (P_{avg}/z_{PA}). La puissance économisée par unité d'envergure est égale au produit de la variation du coefficient de frottement de paroi ΔC_f , de la pression dynamique $\rho U_{\infty}^2/2$, de la corde c dans le sens de l'écoulement et de la vitesse à l'infini U_{∞} , comme le montre l'équation 10. Il convient de noter qu'une négation de Cf est effectuée ici puisqu'une réduction de la traînée de frottement pariétal représente un gain de puissance.

$$\eta = \frac{(-\Delta C_f) \cdot \frac{\rho U_{\infty}^2}{2} \cdot c \cdot U_{\infty}}{(P_{avg}/z_{PA})}$$
(10)



Figure 33.: Net-gain estimation due to the achieved transition delay. The dotted (black) line represents a gain of $\eta = 1$. The investigated flow conditions correspond to $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$.

L'estimation du gain net réalisé grâce aux différentes configurations de forçage étudiées ici est présentée dans la figure 33. La configuration de forçage uniforme utilisée pour la méthode BFM, tout en étant robuste, n'est pas très efficace. Le gain maximal pour cette configuration est seulement de $\eta = 0.4$ ($P_{avg}/z_{PA} = 1.28$ W m⁻¹, $c_{\mu} = -0.07$). En fait, pour le cas avec une puissance d'entrée maximale de $P_{avg}/z_{PA} = 10.5$ W m⁻¹ ($c_{\mu} = -0.27$), qui produit le retard de transition le plus élevé parmi tous les cas considérés ($\Delta x_{tr}/c = 4.5\%$), le gain net est de seulement $\eta = 0.16$. Ainsi, la puissance consommée pour réaliser le contrôle nécessaire est supérieure à celle obtenue par la réduction de la traînée de frottement résultant de la nouvelle position de transition.

Les configurations de forçage modulées spatialement montrent une amélioration significative de l'efficacité globale du système. En particulier avec les configurations UFD4_{in}, un gain maximum de $\eta = 6.9$ est réalisé grâce à un fonctionnement avec une faible puissance d'entrée. Les cas de faible puissance d'entrée de la configuration UFD5 entraînent également un gain supérieur à $\eta = 1$. Cependant, dans ces deux configurations, le gain diminue avec l'augmentation de la puissance d'entrée et devient inférieur à $\eta = 1$, même si le délai de transition obtenu continue d'augmenter. La configuration UFD4_{out}, bien qu'elle ne soit pas aussi efficace que les autres méthodes, reste plus efficace que le cas du forçage uniforme de la méthode BFM. Ce résultat suggère que même un petit retard de transition peut entraîner une amélioration considérable de l'efficacité du système global. Cela est particulièrement utile pour les nombres de Reynold plus élevés, où l'autorité du système de contrôle, bien que marginale, pourrait entraîner des améliorations significatives de l'efficacité du système.

6. CONTRÔLE DES INSTABILITÉS PRIMAIRES ET SECONDAIRES D'UNE COUCHE DE MÉLANGE PLANE

Ce chapitre s'intéresse au contrôle d'une couche de mélange formée par deux écoulements parallèles ayant des vitesses d'écoulement libres différentes et qui interagissent l'un avec l'autre en aval du bord de fuite d'une plaque séparatrice. Ce type d'écoulement est dominé par deux mécanismes d'instabilité. L'instabilité primaire de type *Kelvin-Helmholtz* (KH) va régir l'évolution des structures d'écoulement tourbillonnaire à grande échelle et le contrôle de cette instabilité est fortement dépendant de la fréquence. Ces tourbillons à grande échelle présente un intérêt primordial puisqu'ils sont responsables de la croissance de la couche de mélange lors de son épanouissement. En raison de la présence d'instabilités sub-harmoniques et de leur phase relative par rapport à l'instabilité fondamentale KH, ces tourbillons peuvent fusionner entrainant ainsi une augmentation brusque de la largeur de la couche de mélange. Nous avons vu précédemment que l'instabilité secondaire se manifeste sous la forme de tourbillons orientés dans le sens de l'écoulement distribués le long de l'envergure. Ces tourbillons s'enroulent autour des noyaux des tourbillons primaires et lorsque ces derniers sont appariés, ces tourbillons sont engloutis dans le nouveau tourbillon primaire. Cela finit par provoquer l'apparition de turbulences à petite échelle, ce qui contribue à améliorer le mélange. Étant donné que ces tourbillons sont répartis le long de la couche de mélange, ces instabilités peuvent être manipulées par un forçage spatial appliqué le long du bord de fuite de la plaque séparatrice. Dans cette thèse, les instabilités primaires sont contrôlées par un actionneur plasma DBD exerçant une force EFD uniforme sur toute l'envergure. La réponse en fréquence de ces actionneurs est exploitée pour exciter ces instabilités et manipuler les caractéristiques de stabilité du champ d'écoulement de la couche de mélange. Le caractère unique de ce forçage doit être souligné ici. Le forçage EFD exercé n'a pas de forte composante verticale et la perturbation impartie est confinée près de la surface diélectrique. Ainsi, le contrôle nécessaire de l'instabilité est réalisé ici de manière indirecte, en manipulant le cisaillement et le contenu spectral dans la couche limite entrante. Le forçage largement horizontal reproduit l'effet des instabilités de la couche limite (c'est-à-dire les ondes de Tollmien-Schlichting) qui se développent en amont du bord de fuite de la plaque de séparation. Ainsi, la technique de forçage choisie fournit un environnement de forçage plus "réaliste" que les autres techniques et reproduit les caractéristiques des perturbations dans le cas naturel, bien qu'à une fréquence et une amplitude bien contrôlées. Dans un second temps, les modes d'instabilité secondaire sont excités à l'aide d'actionneur modulée spatialement. Le déploiement d'actionneurs DBD à cette fin offre un autre scénario unique. Alors que le forçage modulé en espace est obtenu grâce à la géométrie de l'actionneur, le signal électrique d'entrée sera utilisé pour imposer des fluctuations correspondant aux instabilités KH primaires.

6.1. Conditions expérimentales des essais

6.1.1. Modèle

Les expériences sur les couches de mélange ont été menées dans une soufflerie en boucle fermée, spécialement modifiée pour permettre des études sur les couches de mélange à l'Institut PPrime de l'Université de Poitiers, en France. La plaque séparatirce constituée d'aluminium se termine par un insert en PMMA. Cet insert constitue la pièce finale de la plaque séparatrice (voir figure 12b). Il s'agit d'un cône de 120 mm de long de 3.1° sur les côtés supérieur et inférieur. Ainsi la plaque séparatrice se termine par un bord de fuite de 3 mm d'épaisseur. L'utilisation du PMMA comme matériau pour cette partie du modèle est motivée par son utilisation comme diélectrique pour l'actionneur plasma. À cet effet, une incision de 40 mm de large (dans le sens de l'écoulement) a été pratiquée sur la face inférieure, de sorte que 2 mm de PMMA compose la face supérieure (épaisseur du substrat diélectrique).

6.1.2. Régimes d'étude

La configuration de la couche de mélange plane a été générée avec deux flux d'air à des vitesses différentes. La vitesse du flux inférieur était de $U_1 = 5 \text{ m s}^{-1}$ et celle du flux supérieur était

maintenue constante à $U_2 = 10 \text{ m s}^{-1}$. Il en résulte un rapport de vitesse de $r = U_1/U_2 = 0.5$. Deux autres paramètres de vitesse sont utilisés dans cet article et sont présentés ici. Le premier est la différence de vitesse entre les deux écoulements $\Delta U = U_2 - U_1 = 5 \text{ m s}^{-1}$. Le second paramètre est la moyenne des deux vitesses $\overline{U} = (U_2 + U_1)/2 = 7.5 \text{ m s}^{-1}$, qui devrait être la vitesse de convection des tourbillons dans la couche de mélange qui s'ensuit (Oster, Wygnanski, & Fiedler, 1977). Avec ces paramètres de vitesse, un second rapport de vitesse est généralement définit comme $R = \Delta U/(2U)$ (ici R = 0.33). Dans ces conditions d'écoulement, l'intensité de la turbulence mesurée par anémométrie à fil chaud (sonde à fil unique, bande passante filtrée entre 10 et 1000 Hz) dans les écoulements à basse et haute vitesse sont respectivement Tu_1/\overline{U} 0.9% et Tu_2/\overline{U} 0.3%. Enfin, il convient de rappeler que les deux couches limites de la plaque séparatrice ne sont pas déclenchées. Il en résulte que les profils de vitesse sur la plaque séparatrice correspondent à des profils de Blasius, indiquant que les deux couches limites sont laminaires. Ceci est confirmé par les facteurs de forme calculés avec les données de vitesse à la surface de la plaque de séparation jusqu'à $y = 2 \cdot \delta_{99}$ qui sont $H_1 \approx 2.8$ et $H_2 \approx 2.9$. L'épaisseur de la couche limite de la plaque de séparation du côté haute vitesse à x = -1 mm a été calculée comme étant égale à: $\theta_i = 0.55$ mm. Cette valeur est désormais utilisée pour adimensionaliser les coordonnées spatiales et autres paramètres de longueur lorsque cela est nécessaire. Le nombre de Reynold basé sur l'épaisseur de quantité de mouvement et la viscosité cinématique est calculé par $Re_{\theta_i} = \overline{U}\theta_i/\nu$. Pour les conditions d'écoulement étudiées où la vitesse moyenne entre les deux courants est $\overline{U} = 7.5 \text{ m s}^{-1}$ et en supposant que la viscosité cinématique soit $\nu = 1.47 \cdot 10^{-5}$ $m^2 s^{-1}$ (de l'air à 15° C), le nombre de Reynold est $Re_{\theta_i} = 280$.

6.1.3. Description des actionneurs

La construction est telle que l'actionneur exerce une force volumique uniforme sur la partie supérieure de la plaque séparatrice. Le bord de l'électrode active est positionnée à $x/\theta_i = -22$ (x = -12 mm). L'actionneur fonctionne à l'aide d'une forme d'onde sinusoïdale modulée par une fonction porte à basse fréquence. La fréquence d'alimentation de la décharge est de $f_{ac} = 2$ kHz, qui est largement supérieur à la fréquence de l'instabilité la plus amplifiée. Cette fréquence AC est responsable du processus d'ionisation et donc de la production de la force EFD. Une modulation par porte à basse fréquence (f_b) a été appliquée à la forme d'onde sinusoïdale d'entrée, ce qui a donné un signal d'entrée final composée d'une succession de phase on et off. Le paramètre de fonctionnement que l'on fait varier pour étudier l'effet de l'excitation de différents modes d'instabilité KH dans la couche de mélange est la fréquence de la modulation par porte f_b . Cette fréquence sera représentée par le nombre de Strouhal correspondant, calculé comme suit: $St_f = f_b \theta_i/\overline{U}$.

Afin d'obtenir une estimation de l'amplitude du forçage par rapport au champ d'écoulement expérimental, le coefficient de moment (c_{μ}) est également utilisé ici (voir chapitre précédent). Ici, ρ (= 1.225 kg m⁻³) est la densité de l'air dans les conditions ambiantes, θ_i (= 0.55 mm) est l'épaisseur du moment de la couche limite latérale à haute vitesse près du bord de fuite de la plaque de séparation et U_2 est la vitesse de l'écoulement à l'infini. Le rapport w_S/λ_z doit tenir compte du rapport-cyclique lorsqu'un forçage tridimensionnel modulé dans le sens de l'envergure est appliqué à l'écoulement. Dans le cas d'un forçage uniforme, ce rapport est naturellement égal à 1. La variation du coefficient de moment calculée de cette manière avec la puissance moyenne consommée est présentée dans la figure 34. Correspondant à la poussée parallèle à la paroi générée présentée sur la figure 22c, le coefficient de moment augmente lui aussi de manière linéaire avec la puissance consommée par l'actionneur. Pour la plupart des cas présentés dans cette section, l'actionneur plasma fonctionne à une amplitude de tension $V_{ac} = 14$ kV et une



Figure 34.: Variation of momentum coefficient c_{μ} with average input power P_{avg} . Solid red marker is an estimate of c_{μ} when the spanwise-uniform DBD plasma actuator is operated at $V_{ac} = 14$ kV, $f_{ac} = 2$ kHz with a burst-modulation duty of 50%.

fréquence $f_{ac} = 2$ kHz. Le rapport cyclique du cycle de modulation est également maintenu constant à 50%. La puissance moyenne consommée par l'actionneur DBD lorsqu'il fonctionne à $V_{ac} = 14$ kV, $f_{ac} = 2$ kHz est de $P_{avg} = 52.2$ W m⁻¹. Si l'on considère l'application d'une modulation avec un rapport cyclique de 50% sur le signal d'entrée en courant alternatif pour manipuler les instabilités primaires dans la couche de mélange, cette valeur de puissance peut être réduite de moitié (voir Benard & Moreau, 2009). Ainsi, la puissance moyenne consommée par l'actionneur lors de l'application de la force EFD nécessaire sur la couche de mélange est $P_{avg,f} = 26.1$ W m⁻¹. Pour cette valeur de puissance, le coefficient c_{μ} est de 0.22 si l'on se base sur l'estimation de la poussée à partir des données de la littérature. De meme, la vitesse maximale moyenne dans le temps du jet induit est estimée à $U_f \approx 2.5$ m s⁻¹ (à partir des données expérimentales de Jolibois & Moreau, 2009). Le rapport de vitesse entre cette vitesse du vent électrique produit et le flux à grande vitesse est alors $U_f/U_2 \approx 0.25$. Le coefficient de $c_{\mu} = 0.22$ pour cette puissance d'entrée est comparable au rapport de vitesse estimé. Comme la quantité de mouvement du jet libre ne change pas, le coefficient de quantité de mouvement calculé est représentatif de ce rapport de vitesse.

6.2. Caractérisation de l'écoulement naturel

6.2.1. Champs d'écoulement moyens

La plaque de séparation utilisée pour cette étude a un bord de fuite de 3 mm d'épaisseur. Il en résulte un effet de sillage qui se manifeste par la formation d'une zone de recirculation juste en aval du bord de fuite de la plaque de séparation. Dans la couche de mélange actuelle non forcée, on observe que la zone de recirculation dure jusqu'à environ $x/\theta_i = 44$ ($x \approx 24$ mm). Au-delà de cet emplacement dans le sens du courant, on observe l'apparition de la couche de mélange. L'épaisseur de vorticité de la couche de mélange non forcée est présentée sur la figure 35. Jusqu'à $x/\theta_i \approx 30$, aucune croissance n'est observée en raison de la région de recirculation. Dans la zone de recirculation, l'épaisseur de vorticité mesure $\delta_{\omega}/\theta_i \approx 3$, ce qui, en termes de dimensions ($\delta_{\omega} = 1.65$ mm), correspond approximativement à la moitié de l'épaisseur du bord de fuite de la plaque séparatrice. Une augmentation de l'épaisseur de vorticité est observée juste en aval de cet emplacement, puis cette augmentation devient linéaire après x/θ_i 40. Cela correspond la région laminaire où l'instabilité linéaire provoque l'enroulement de la couche de cisaillement à la



Figure 35.: Thickness of the unforced mixing layer along the streamwise direction.

fréquence fondamentale (ou naturellement préférée), ce qui entraîne la formation de tourbillons orientés dans le sens de l'envergure (Michalke, 1965a; Winant & Browand, 1974). Il s'ensuit une croissance linéaire à partir de $x/\theta_i \approx 100$ qui peut être attribuée à l'appariement des tourbillons se produisant de manière stochastique dans l'espace et le temps (Winant & Browand, 1974; Oster & Wygnanski, 1982; Ho & Huang, 1982). Il convient également de noter que le schéma de croissance de la couche de mélange actuelle ressemble à celui de la couche de mélange résultant d'un bord de fuite mince dans les travaux numériques de Laizet et al. (2010).

Le taux de croissance de la couche de mélange peut être approché par la méthode définie par Abramovich (1963)-Sabin (1965). Dans la région de croissance linéaire de la couche de mélange actuelle, le taux de croissance global est bien approximé par $d\delta_{\omega}/dx = 0.25 \cdot R$ (ligne rouge en pointillé dans la figure 35). La valeur du paramètre d'ajustement c = 0.25 est supérieure à la plage généralement acceptée de c = 0.16 to 0.18 (Brown & Roshko, 1974; Browand & Latigo, 1979; Mehta, 1991). Le taux de croissance élevé de cette étude pourrait être attribué au niveau plus élevé de turbulence du flux à faible vitesse, qui entraîne une augmentation de l'entraînement et donc du taux de croissance (Pui & Gartshore, 1979). Cependant, l'interaction complexe du sillage de la plaque de séparation avec la couche de mélange ne peut pas non plus être négligée. La valeur du paramètre d'ajustement c est inférieure à celle rapportée par Laizet et al. (2010) (c = 0.34) mais similaire à celle observée dans les expériences de Huang & Ho (1990) (approximativement 0.27). La figure 35 montre également l'épaisseur intégrale de la couche de mélange (souvent appelée épaisseur de la quantité de mouvement θ). Dans la région turbulente de la couche de mélange non forcée actuelle, l'épaisseur intégrale augmente également de façon linéaire avec la position longitudinale. Le taux de croissance mesuré correspondant à l'épaisseur intégrale est $d\theta/dx \approx 0.017 \ (= 0.05 \cdot R)$. Ce taux est comparable à celui observé par Ho & Huang (1982) $(d\theta/dx \approx 0.013)$ dans leur couche de mélange avec un rapport de vitesse (R = 0.31) similaire à celui de la présente étude.

6.2.2. Identification des instabilités par analyse LST et périodogrammes

La figure 36a montre les taux d'amplification spatiale non dimensionnelle $(-\alpha_i \cdot \delta_\omega(x))$ des ondes d'instabilité KH de différentes fréquences calculées à l'aide de la LST. Ici, l'axe vertical indique la fréquence de l'instabilité de KH représentée sous forme non dimensionnelle par le nombre de Strouhal $St_{\theta_i} = f\theta_i/\overline{U}$, où θ_i et \overline{U} sont respectivement l'épaisseur du moment de la couche limite à haute vitesse et la vitesse moyenne des deux flux. Ces paramètres sont choisis afin de calculer le nombre de Strouhal à l'image des travaux de Monkewitz & Huerre (1982) and Ho & Huang (1982).

On constate que la fréquence des fluctuations les plus instables augmente jusqu'à $x/\theta_i \approx 14$ puis commence à diminuer, comme le montre la courbe noire de la figure 36a (et c) qui représente le mode d'instabilité le plus amplifié à chaque position. Jusqu'à $x/\theta_i \approx 40$, les fluctuations supérieures à $St_{\theta_i} = 0.044$ sont les plus amplifiées. Cependant, en raison de l'apparition du sillage de la plaque de séparatrice, les fluctuations dans cette région sont escomptées. Au-delà de cette position, la fréquence des fluctuations les plus instables diminue comme prévu avec la position longitudinale. Cette caractéristique est d'autant plus claire lorsque l'on s'intéresse à l'amplification des facteurs N des différents modes d'instabilité (figure 36b). Cela montre que les fluctuations à basse fréquence sont plus amplifiées en aval. Cependant, il est important de noter qu'une fois que l'instabilité fondamentale devient neutre, la croissance de l'instabilité sousharmonique est influencée par les interactions non linéaire avec la fondamentale (Monkewitz, 1988). Cela rend alors le calcul par LST inexact. Ainsi, seules des tendances peuvent être observées et aucun commentaire ne peut être fait concernant les valeurs calculées. L'épaisseur de vorticité de cette couche de mélange dans le sens du courant présentée sur la figure 35 montre une augmentation soudaine entre $x/\theta_i = 40$ et jusqu'à 75, ce qui signifie l'enroulement des tourbillons dans le sens de l'envergure en raison des ondes de perturbation générées par l'instabilité KH fondamentale. Dans cette région, la LST prévoit que les modes les plus instables se situent entre $St_{\theta_i} = 0.034$ et 0.013, comme le montre la vue agrandie du taux d'amplification de la figure 36c. Ainsi, l'un de ces modes doit être l'instabilité KH fondamentale de la configuration actuelle de la couche de mélange.

Une autre méthode pour déterminer l'instabilité fondamentale de KH consiste à mesurer la fréquence de passage des tourbillons primaires dans la région initiale, la fréquence dominante correspondra alors à la fréquence de l'instabilité fondamentale (Ho & Huang, 1982). À cet effet, des séries temporelles des fluctuations de la composante de vitesse transversale (v') échantillonnées le long de la ligne $y_{0.5}$ dans la couche de mélange non forcée sont analysées dans le domaine de Fourier. Les spectres de puissance sont calculés selon la méthode de Welch (1967). À cet effet, les séries temporelles acquises ont été décomposées en segments qui se chevauchent, chacun contenant 1900 instantanés et se chevauchant de 50%. Il en est résulté une résolution de fréquence finale de $\Delta S t_{\theta_i} = 3.67 \cdot 10^{-4}$ ($\Delta f = 5$ Hz).

Les spectres de puissance calculés sont présentés à la figure 37. Les fluctuations avec la plus grande énergie correspondent à $St_{\theta_i} = 0.022$ (bande de fréquence de 300 ± 10 Hz). On observe que ces fluctuations sont les plus fortes à proximité de la plaque de séparation et qu'elles diminuent en aval, devenant comparativement faibles pour $x/\theta_i > 350$. Ainsi, l'instabilité fondamentale du KH dans la configuration étudiée correspond à un nombre de Strouhal de $St_0 = 0.022$ $(f_0 = 300$ Hz). La longueur d'onde correspondante de l'onde d'instabilité fondamentale est λ_0/θ_i $(= \overline{U}/f_0\theta_i = 1/St_0) = 45.5$ ($\lambda_0 = 25$ mm). Cette instabilité se situe dans la bande de fréquences que le calcul de la LST a prévue comme pouvant être l'instabilité fondamentale. Ceci est une démonstration du mécanisme de rétroaction. Comme le soulignent Ho & Huang (1982), les calculs de stabilité fournissent une bande de fréquences possibles qui peut être l'instabilité KH fondamentale. Mais la rétroaction des événements d'appariement des tourbillons se produisant en aval accorde le développement de la couche de mélange à une fréquence spécifique.

Le nombre de Strouhal calculé avec l'épaisseur du moment intégral local à $x/\theta_i = 91$ est $f_0\theta(x)/\overline{U} = 0.03$, ce qui corrobore bien les travaux antérieurs impliquant la théorie de la stabilité linéaire (voir Ho & Huerre, 1984). Le choix de cet emplacement est motivé par les suggestions de Ho & Huerre (1984), qui recommandent que la comparaison entre les valeurs mesurées et calculées du nombre de Strouhal se fasse à une longueur d'onde d'instabilité en aval du bord de fuite et surtout avant l'apparition de toute croissance non linéaire causée par interactions des



Figure 36.: (a) Non-dimensional amplification rate $(-\alpha_i \cdot \delta_{\omega}(x))$ of primary instabilities in the unforced mixing layer at different frequencies computed using LST. The black curve represents the most amplified instability mode at each streamwise station. (b) Amplification N-factors computed from $-\alpha_i$ according to equation 3.6. (c) Zoomed-in view of non-dimensional amplification rate $(-\alpha_i \cdot \delta_{\omega}(x))$ from the dashed-green rectangle in (a).

instabilités sous-harmoniques.

6.2.3. Analyse de l'écoulement par SPOD

Les premiers résultats de la décomposition par méthode SPOD montre que les fluctuations sur $St_{\theta_i} = 0.022$ contiennent la plus grande énergie (données non montrées ici). Ceci est strictement conforme à la discussion précédente où la fréquence de l'instabilité fondamentale s'est avérée être $St_0 = 0.022$. Le premier mode propre à cette fréquence contient 89% de la TKE de tous les modes à cette fréquence. La composante verticale de la fonction propre correspondante



Figure 37.: Non-dimensional, normalised power spectra $(\Phi_N^{v'} \cdot \Delta f/\overline{U}^2 \text{ (dB)}, \Delta St_{\theta_i} = 3.67 \cdot 10^{-4}, \Delta f = 5 \text{ Hz})$ of the v'-velocity component sampled along the $y_{0.5}$ -trace.

 $(\Psi_v(0.022, 1))$ est présentée dans la figure 38a. L'échelle de couleurs de la figure montre la distribution de la contrainte de cisaillement de Reynold du champ fluctuant correspondant. À cette fin, le champ d'écoulement a été reconstruit avec ce mode SPOD particulier.

Les structures cohérentes sont clairement visibles jusqu'à $x/\theta_i \approx 250$ dans ce mode. Il est important de noter que ces structures ne représentent que l'organisation spatiale des fluctuations dans le champ d'écoulement et ne représentent aucune structure physique de l'écoulement. Au départ, ces structures cohérentes forment un angle obtus avec la direction du courant. Cela correspond au domaine spatial où la contrainte de Reynold est positive. Comme l'ont montré Rajaee et al. (1994) qui ont réalisé une POD "instantanée" (Sirovich, 1987) sur leurs mesures au fil chaud de la couche de mélange plane, cette contrainte de Reynold positive qui agit contre le taux de déformation de cisaillement positif du champ d'écoulement moyen représente le transfert d'énergie de l'écoulement moyen vers le champ fluctuant. Les structures atteignent finalement un alignement vertical à $x/\theta_i \approx 80$ qui représente le mode atteignant son amplitude maximale. Au-delà de ce point, les structures ont un angle aigu avec la direction du flux et la contrainte de Reynold dans cette région est observée comme étant négative. L'action de cette contrainte de Reynold négative contre le taux de déformation positif du flux moyen entraîne le transfert d'énergie du champ fluctuant correspondant vers le flux moyen.

La fonction propre $\Psi_v(0.011, 1)$ contenant les fluctuations au niveau de la première sousharmonique de l'instabilité fondamentale est présentée dans la figure 6.10b. On observe que ces fluctuations augmentent jusqu'à $x/\theta_i \approx 340$ où elles atteignent leur amplitude maximale et diminuent plus en aval. Les structures relatives aux fluctuations de la deuxième sous-harmonique de l'instabilité fondamentale sont représentées par la fonction propre $\Psi_v(0.0055, 1)$ dans la figure 38c. Par rapport aux fluctuations précédentes, les structures initiales ne sont observées ici que plus en aval à $x/\theta_i \approx 210$ et donc, ces fluctuations ne s'amplifient qu'au-delà de cette position aval. Ces fluctuations s'amplifient jusqu'à $x/\theta_i \approx 700$, au-delà duquel elles commencent à s'amortir. Enfin, la fonction propre des fluctuations à $St_{\theta_i} = 0.0117$ qui est proche de la première instabilité sous-harmonique est présentée dans la figure 38d. L'organisation de ces fluctuations est présentée ici car l'instabilité correspondant à ces fluctuations sera étudiée ultérieurement dans le manuscrit original, afin d'étudier l'effet d'un forçage par désaccord de la première sousharmonique. Les structures de cette fonction propre ressemblent à celles de la première sousharmonique, mais ces fluctuations atteignent leur amplitude maximale légèrement en amont à $x/\theta_i \approx 320$, au-delà de laquelle elles commencent à se désaccorder.

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Figure 38.: Vertical component of the eigenfunction of the first mode $(\Psi_v(St_{\theta_i}, 1))$ at different frequencies (grey - negative; black - positive). The colour-scale in (a) represents the normalised Reynold's stress component of the fluctuations corresponding to this mode $(R_{xy}(St_{\theta_i}, 1)/(\Delta U)^2)$.

6.3. Contrôle de l'instabilité primaire

Dans cette partie, les résultats des essais destinés à manipuler l'instabilité primaire et les structures cohérentes de KH sont présentés. Ici le forçage est linéaire selon l'envergure quel que soit la condition d'essai. Dans un premier temps, le champ de vitesse et l'écoulement non contrôlé sont détaillés. Dans un second temps, une description des effets du forçage sur les quantités moyennes et la dynamique tourbillonnaire est proposée.



Figure 39.: Growth of the mixing layers when different primary KH instabilities are forced using the DBD plasma actuator ($c_{\mu} = 0.22$). Dashed (black) lines represent linear fits to different segments of the growth curves. Black (thick) curve represents the growth of the unforced mixing layer and the dashed (red) line represents the corresponding $0.25 \cdot R$ linear fit (see figure 35).

6.3.1. Effets sur les caractéristiques de l'écoulement moyen

L'évolution de la couche de mélange due au forçage EFD uniforme à c_{μ} = 0.22 pour exciter différentes instabilités KH dans la couche de mélange, est discutée ici. Le forçage EFD entraîne un développement initial de la couche de mélange en amont par rapport au cas non forcé. En effet, le forçage réduit l'effet de sillage du bord de la plaque de séparatrice ce qui permet un déclenchement plus rapide du phénomène d'enroulement tourbillonnaire. Indépendamment du mode d'instabilité forcée, le développement initial de la couche de mélange forcée se trouve à $x/\theta_i \approx 15$ qui est en effet en amont où le développement initial de la couche de mélange non forcée est observé. La croissance des différentes couches de mélange forcé est illustrée à la figure 39. Chaque condition de forçage est présentée séparément et comparée à la couche de mélange non forcée. La croissance initiale de toutes les couches de mélange forcées, indiquant l'enroulement des tourbillons bidimensionnels dans le sens de l'envergure, dure jusqu'à $x/\theta_i \approx 60$. Plus en aval, le fait de forcer les différentes instabilités entraîne des modes de croissance différents. Le fait de forcer l'instabilité fondamentale $(St_f = St_0 = 0.022)$ (figure 39a) réduit considérablement la croissance des couches de mélange ($\approx 0.07 \cdot R$) jusqu'à $x/\theta_i = 200$ par rapport à la couche de mélange non forcée. Cela devrait être dû à l'atténuation des instabilités sous-harmoniques (voir le cas mode I de Ho & Huang, 1982) qui jouent un rôle essentiel dans la croissance de la couche de mélange. La couche de mélange croît alors rapidement ($\approx 0.7 \cdot R$) jusqu'à ce que $x/\theta_i = 500$. Plus en aval, la courbe de croissance s'effondre sur celle de la couche de mélange non forcée.

Le forçage à la première sous-harmonique de l'instabilité fondamentale $(St_f = St_0/2 = 0.011)$ conduit à une courbe de croissance différente (figure 39b). Après l'enroulement initial, la couche de mélange commence à croître à partir de $x/\theta_i \approx 100$, pour finalement atteindre une croissance linéaire à un taux de $\approx 0.43 \cdot R$. Après cela, la croissance s'arrête temporairement de $x/\theta_i \approx 300$ à environ 600 ($\approx 0.089 \cdot R$). La couche croît ensuite de manière linéaire, jusqu'à la fin du domaine spatial mesuré, à un taux similaire ($\approx 0.26 \cdot R$) à celui de la couche de mélange non forcée. Cette courbe de croissance présente des similitudes qualitatives avec celles observées dans les travaux expérimentaux de Oster & Wygnanski (1982), Ho & Huang (1982) et Huang & Ho (1990), lorsqu'un forçage sous-harmonique a été appliqué sur la couche de mélange.

6.3.2. Réorganisation de l'écoulement – Forçage harmonique

Les résultats rapportés jusqu'à présent ont démontré que le forçage exercé par l'actionneur plasma DBD peut influencer la croissance de la couche de mélange et ses caractéristiques spectrales. Afin de mieux comprendre l'effet de ce forçage basse fréquence, l'organisation spatiale des champs fluctuants est maintenant analysée. À cette fin, la méthode SPOD a été appliquée sur les séries temporelles obtenues par PIV.

Les résultats de la méthode SPOD appliquée à la couche de mélange forcée à l'instabilité fondamentale sont présentés dans la figure 40. La composante de TKE à différentes fréquences (en noir sur la figure 40a) montre que les modes ayant la plus grande énergie se rapportent à l'instabilité fondamentale forcée $St_f = St_0 = 0.022$. La concentration de TKE est également observée aux harmoniques du mode forcé comme prévu; cependant, les instabilités sous-harmoniques ont été atténuées. Celles-ci sont également évidentes dans l'amplitude des fluctuations correspondantes $(A_v(St_{\theta_i}, 1))$ présenté dans la figure 40b. La composante verticale de la fonction propre du premier mode SPOD contenant les fluctuations de l'instabilité forcée ($\Psi_v(0.022,1)$) est présentée sur la figure 40c. En comparaison avec le même mode dans la couche de mélange non forcée, des structures cohérentes se forment plus en amont. C'est une conséquence directe du déplacement vers l'amont du développement de la couche de mélange conséquence du forçage EFD appliqué. Les fluctuations liées à ce mode s'amplifient à un taux de $-\alpha_i \cdot \theta_i \approx 0.0071$ et atteignent leur amplitude maximale à environ $x/\theta_i \approx 80$ (comme dans le cas d'un écoulement non forcé) ce qui correspond à l'alignement vertical des structures. Toutefois, l'amplitude maximale de ces fluctuations est nettement plus élevée $(A_{v,max}(0.022,1)/U = 0.326)$ par rapport à la couche de mélange non forcée. Ici, les structures restent alignées verticalement et sont présentes jusqu'à environ $x/\theta_i \approx 320$. La composante verticale de la fonction $\Psi_v(0.011, 1)$ du premier mode SPOD contenant les fluctuations relatives à la première sous-harmonique de l'instabilité forcée (figure 40d). Les structures initiales sont observées plus en aval (à environ $x/\theta_i \approx 250$) de celles de la fonction propre correspondante dans la couche de mélange non forcée. Par conséquent, ces fluctuations ne s'amplifient qu'après cette position (figure 40b). Ces fluctuations atteignent leur amplitude maximale à $x/\theta_i \approx 550$, où l'amplitude de l'instabilité fondamentale forcée est maintenant très faible. De même, des structures cohérentes relatives à la deuxième sous-harmonique de l'instabilité fondamentale ($St_{\theta_i} = 0.0055$, figure 40e) sont observées plus en aval à $x/\theta_i \approx 550$ également. Par rapport au cas non forcé, l'amplitude de ces deux instabilités sous-harmoniques est nettement inférieure. Cela indique clairement l'atténuation des modes d'instabilité sousharmoniques lorsque l'instabilité fondamentale est forcée.

L'effet du forçage sur la dynamique et les interactions des structures fluidiques cohérentes est maintenant discuté. Grâce à une synchronisation des acquisitions PIV et du signal de commande de l'actionneur il est possible e reconstruire l'écoulement à l'aide de moyenne de phase. Le critère Q proposé par Hunt et al. (1988), est ensuite appliqué à chacun des champs moyennés en phase pour détecter les tourbillons cohérents dans le champ. Ici les champs de vitesse sont reconstruits



Figure 40.: SPOD decomposition of the mixing layer forced at $St_f = St_0 = 0.022$ ($c_{\mu} = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of selected frequencies along the $y_{0.5}(x)$ trace. (c), (d) and (e) show the vertical component of the eigenfunction of the first mode $(\Psi_v(St_{\theta_i}, 1))$ of these different frequencies (grey - negative; black - positive).

à partir d'un nombre limité de mode SPOD. Une telle reconstruction permet de sélectionner des instabilités et d'analyser leur influence sur la dynamique globale de l'écoulement. En outre, cet exercice agit également comme un filtre pour éliminer tout bruit à haute fréquence des séries temporelles de PIV acquises.

Le calcul du critère Q sur les différents champs moyennés par phase est présenté dans la figure 41. Le forçage conduit à des tourbillons cohérents qui se forment juste en aval du bord de fuite de la plaque séparatrice. Tous les tourbillons sont séparés de leurs voisins par $\Delta x/\theta_i = 45.5$

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Figure 41.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0 = 0.022$ ($c_\mu = 0.22$).

qui est égal à la longueur d'onde de l'instabilité fondamentale forcée $(\lambda_0/\theta_i = \overline{U}/f_0\theta_i = 1/St_0)$. Ainsi, la fréquence de passage du tourbillon est effectivement égale à l'instabilité forcée. Aucun cas d'appariement de tourbillons n'est observé ici en raison de l'atténuation des modes d'instabilité aux sous-harmoniques par le forçage de l'instabilité au fondamental. Cette inhibition de l'interaction des tourbillons se traduit par une croissance négligeable de cette couche de mélange jusqu'à $x/\theta_i = 200$, comme discuté précédemment. On observe que les structures cohérentes se renforcent jusqu'à la position $x/\theta_i \approx 80$. Cette position est proche de l'endroit où les fluctuations liées à l'instabilité fondamentale forcée ont atteint leur plus grande amplitude. En aval, les structures tourbillonnaires semblent s'atténuer et aucun tourbillon clair n'est perceptible après $x/\theta_i \approx 450$. Cela suggère que la couche de mélange actuelle n'est plus en phase avec la fréquence de modulation du forçage appliqué.

6.3.3. Réorganisation de l'écoulement – Forçage sous-harmonique

Les résultats de la SPOD appliquée sur la couche de mélange forcée à la première sous-harmonique de l'instabilité fondamentale ($St_f = St_0/2 = 0.011$) sont maintenant discutés (figure 42). Comme précédemment, un pic important est observé dans la distribution TKE à l'instabilité forcée, comme le montre la figure 42a. Ces fluctuations contiennent la plus grande énergie dans la couche de mélange actuelle ($\approx 7\%$), le premier mode SPOD représentant la majeure partie de cette énergie. On observe également une concentration de l'énergie dans les modes d'instabilité fondamentale. La composante verticale de la fonction $\Psi_v(0.011, 1)$ contenant les fluctuations de l'instabilité sous-harmonique forcée est représentée sur la figure 42d et son amplitude selon la position longitudinale est présentée sur la figure 42b. Contrairement au flux non forcé où l'amplitude initiale de ces fluctuations était négligeable, ici, l'amplitude de ce mode à $x/\theta_i = 0$ est $A_v(0.011,1)/\overline{U} = 0.018$. C'est une conséquence directe de l'ajout de quantité de mouvement imposé par l'actionneur plasma. Ces fluctuations augmentent ensuite rapidement jusqu'à $x/\theta_i \approx 50$, elles s'amplifient ensuite progressivement et atteignent leur amplitude maximale à $x/\theta_i \approx 320$, près de l'endroit où la croissance de la couche de mélange s'arrête brusquement. Comme dans le cas précédent où l'instabilité fondamentale est forcée, ici aussi, l'instabilité sousharmonique forcée atteint son amplitude maximale au même endroit que celui observé dans la couche de mélange non forcée. Cependant, l'amplitude maximale est plus élevée que dans le cas non forcé et consécutivement l'instabilité est amortie beaucoup plus tardivement.

La composante verticale de la fonction $\Psi_v(0.022, 1)$ contenant les fluctuations de l'instabilité fondamentale est illustrée à la figure 42c. Cette fonction présente quelques similitudes avec celle de la couche de mélange où l'instabilité fondamentale a été forcée. L'amplitude de cette instabilité est négligeable au départ, comme le montre la figure 42b, ce qui suggère que ces fluctuations ne sont pas influencées par le forçage appliqué. Elles croient ensuite rapidement jusqu'à $x/\theta_i \approx 50$ où elles atteignent leur amplitude maximale $(A_{v,max}(0.022,1)/\overline{U} = 0.053)$, puis diminue progressivement par la suite. L'amplitude maximale de ce mode est ici plus faible que non seulement la couche de mélange forcée à l'instabilité fondamentale (ce qui est prévisible), mais aussi le cas d'écoulement non forcé. Lorsque l'amplitude des fluctuations correspondant à l'instabilité fondamentale commence à diminuer, celle de l'instabilité sous-harmonique forcée augmente. Ceci est représentatif du fait que la sous-harmonique est renforcée par le transfert d'énergie du champ d'écoulement moyen qui oscille au niveau de l'instabilité fondamentale. Cela conduit à une croissance en résonnance de l'instabilité sous-harmonique entre $x/i \approx 80$ et 300 à un taux de $\alpha_i \cdot \theta_i \approx 0.0003$. Ce mode d'instabilité forcée atteint une amplitude maximale de $A_{v,max}(0.011, 1)/\overline{U} = 0.099$.

La distribution du critère Q calculés sur les champs de vitesse en moyenne de phase est présentée dans la figure 43. On observe que des tourbillons cohérents se forment juste en aval de la plaque de séparation. L'enroulement initial des tourbillons dure jusqu'à $x/\theta_i \approx 50$, qui est proche de la position où l'instabilité fondamentale du KH a atteint son amplitude maximale. Un exemple d'appariement des tourbillons impliquant deux tourbillons successifs (rectangle rouge pointillé dans les champs) est observé dans la couche de mélange. L'interaction de ces deux tourbillons est illustrée dans les instantanés présentés à la figure 43. Parmi ces deux tourbillons, le tourbillon le plus faible est généré par l'instabilité KH fondamentale naturelle au cours du cycle de modulation de la salve précédent. Alors que le tourbillon le plus fort, qui se trouve en amont, est influencé par le forçage appliqué. Ceci est clairement évident lorsque l'on observe les deux premiers tourbillons à la moyenne de phase $\phi = 1\pi$. Les deux tourbillons subissent ensuite un appariement dû à l'induction mutuelle, comme l'expliquent Husain & Hussain (1995). Le processus d'appariement des tourbillons commence à $x/\theta_i \approx 100$, qui correspond à l'endroit

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Figure 42.: SPOD decomposition of the mixing layer forced at $St_f = St_0/2 = 0.011$ ($c_{\mu} = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the nondimensional amplitude ($A_v(St_{\theta_i}, 1)/\overline{U}$) of selected frequencies along the $y_{0.5}(x)$ trace. (c)-(e) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) of these different frequencies (grey - negative; black - positive).

où l'amplitude de l'instabilité sous-harmonique forcée devient supérieure à celle de l'instabilité fondamentale, comme l'observe la figure 42b. Le processus d'appariement se produit lorsque l'amplitude de l'instabilité sous-harmonique forcée augmente de manière linéaire et se termine près de l'endroit où l'amplitude atteint sa valeur maximale, à environ $x/\theta_i \approx 280$. Cela correspond aux observations de Ho & Huang (1982) qui ont montré que le phénomène d'appariement des tourbillons se produit alors que l'instabilité sous-harmonique augmente de manière linéaire. Le processus d'appariement est achevé au bout de deux cycles de forçage et il se traduit par un taux de croissance plus élevé dans cette région. A la suite de l'appariement, les structures



Figure 43.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0/2 = 0.011$ ($c_{\mu} = 0.22$). Snapshots of the vortex pairing process (vortices in dashed red rectangle in the fields) along the low-frequency burst modulation cycle is also shown.

fluidiques ne présentent plus d'interactions.

6.4. Contrôle de l'instabilité secondaire

Il s'agit désormais d'étudier l'effet obtenu lorsque le forçage n'est plus homogène mais discontinu le long du bord de fuite de la plaque séparatrice. L'objectif est de tirer avantage de la nature spatiale du forçage appliqué afin d'exciter des instabilités secondaires. Comme présenté précédemment, les structures longitudinales de l'écoulement résultent de ce mécanisme d'instabilité secondaire et, combinées au phénomène d'appariement des tourbillons qui peut être piloté par l'instabilité primaire, peuvent entraîner une amélioration importante du mélange.

Dans cette partie de la thèse, la géométrie de l'actionneur est constituée de tiges équidistantes qui permettent de localiser le forçage. Chacune de ces tiges a une largeur de $w_S = 3$ mm et sont espacées de 37 mm les unes des autres. Il en résulte une force modulée dans le sens de l'envergure avec une longueur d'onde de $\lambda_z = 37$ mm et un rapport-cyclique de $w_S/\lambda_z = 8\%$. Les conditions d'écoulement et le système de coordonnées/références pour ces investigations sont les mêmes que dans le reste de ce chapitre.

Un système PIV résolue en temps est utilisé comme métrologie de base. Le forçage étant tridimensionnel, les mesures préliminaires proposées se limitent à deux plans de mesure. Le premier plan est localisé au centre d'une région sous l'influence de la décharge plasma (z = 0, $z/\theta_i = 0$), tandis que le second plan est positionné entre deux régions de décharge (z = 14 mm, $z/\theta_i = 25.5$).

6.4.1. Effets sur l'écoulement moyen

La croissance des couches de mélange forcées est présentée sur la figure 44 pour les deux positions le long de l'envergure. Le forçage du mode fondamental conduit à un fort épaississement de la couche de mélange dans le plan de mesure $z/\theta_i = 0$. L'enroulement initial des tourbillons semble donner lieu à une couche de mélange plus épaisse dès les premières positions de mesure (jusqu'à $x/\theta_i \approx 50$). La couche de mélange s'épanouie ensuite à un taux plus élevé ($\approx 0.47 \cdot R$, à partir de $x/\theta_i \approx 150$) que le taux observé dans le cas non forcé jusqu'à $x/\theta_i \approx 325$, après quoi la croissance est temporairement arrêtée jusqu'à $x/\theta_i \approx 430$. Au-delà, la croissance de la couche, bien que d'une ampleur légèrement supérieure, se fait au même rythme que le flux non forcé jusqu'à la fin du domaine spatial. Pour le second plan de mesure à $z/\theta_i = 25.5$, la croissance de la couche de mélange est significativement différente et elle est très similaire aux observations réalisées pour un forçage homogène selon l'envergure. Ces résultats démontrent une fois de plus que l'influence du forçage, même lorsqu'il est modulé en fonction de l'envergure, est remarquable. Ce qui est le plus intéressant, c'est la différence flagrante observée dans la croissance des couches de mélange correspondantes aux deux positions des plans PIV dans le sens de l'envergure. C'est la première indication d'une modulation dans le sens de l'envergure de l'épaisseur de la couche de mélange forcée.

6.4.2. Réorganisation de l'écoulement

Les résultats obtenus après application de la SPOD sur les séries temporelles lorsque l'on impose des fluctuations à l'instabilité fondamentale $St_f = St_0 = 0.022$ sont présentés dans la figure 45. La distribution de l'énergie dans les différents modes SPOD par rapport à la TKE du flux est présentée dans les figures 45a et b. Comme dans le cas correspondant au forçage uniforme, les fluctuations maximales aux deux positions de mesure sont liées à la fréquence de l'instabilité fondamentale forcée et de ses harmoniques. De plus, aucune amplification des instabilités sousharmoniques n'est observée, ceci quelle que soit la position étudiée.

Les fluctuations les plus fortes à $z/\theta_i = 25.5$, la position entre deux zones de décharge, sont



Figure 44.: Growth of the mixing layer when fluctuations pertaining to different primary KH instabilities are imparted using the spanwise-modulated DBD plasma actuator ($c_{\mu} = 0.009$). The unforced mixing layer growth presented here corresponds to that at $z/\theta_i = 0$.

à la fréquence de l'instabilité fondamentale forcée. La fonction propre $\Psi_v(0.022, 1)$ à cette position est représentée sur la figure 45d. Les structures atteignent un alignement vertical avec la direction du courant à $x/\theta_i \approx 80$. Elles restent ainsi orientées jusqu'à $x/\theta_i \approx 300$, au-delà duquel les structures claires ne sont pas perceptibles. L'organisation spatiale de ces fluctuations est très similaire à celle qui se produit lorsque cette instabilité est forcée avec l'actionneur à portée uniforme (voir figure 40c). Cela indique que, bien que l'instabilité fondamentale de KH est forcée bien que l'actionneur présente un forçage localisé et discret. La fonction propre contenant les fluctuations de la première sous-harmonique de l'instabilité fondamentale de KH illustrée à la figure 45f est également similaire à celle du cas de forçage uniforme. Des structures claires sont présentes en aval de $x/\theta_i = 200$ et ces structures atteignent un alignement vertical à $x/\theta_i \approx 450$. Cette région dans le sens de l'écoulement est celle où le taux de croissance de la couche de mélange était significativement plus élevé.

Les structures cohérentes dans la fonction propre ($\Psi_v(0.022, 1)$) à $z/\theta_i = 0$ sont illustrées sur la figure 45c. Ici, les structures cohérentes ne sont observées que jusqu'à $x/\theta_i \approx 170$. En outre, ces structures ont une forme différente de celles dans le champ à $z/\theta_i = 25.5$. Elles semblent être allongées dans la direction positive de l'axe des y, les maxima restant toujours proches de $y/\theta_i = 0$. En aval de $x/\theta_i = 50$, ces allongements s'alignent progressivement dans le sens de l'écoulement, pour finalement atteindre en aval la partie haute de la structure cohérente. Les structures relatives aux fluctuations de la première sous-harmonique présentent également des différences marquées par rapport aux cas précédents. Alors que les structures ne sont observées qu'en aval de $x/\theta_i = 200$, ces structures ont initialement deux maxima. Les maxima en haut semblent disparaître à terme au-delà de $x/\theta_i = 400$. Cependant, il convient de noter que le TKE de ces fluctuations sous-harmoniques est significativement faible (deux ordres de grandeur) par rapport aux fluctuations imparties à l'instabilité fondamentale. Ainsi, l'épaisseur

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Figure 45.: SPOD of the mixing layer forced at $St_f = St_0 = 0.022$ ($c_\mu = 0.009$). Figures on the left correspond to $z/\theta_i = 0$ (a, c, and e) and figures on the right correspond to $z/\theta_i = 25.5$ (b, d, and f). (a) and (b) show component of TKE contained in each mode normalized with the total TKE (in %). (c)-(f) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) at different frequencies (grey - negative; black - positive).

plus importante de la couche de mélange à cette position ne devrait pas être due au mécanisme d'appariement des tourbillons.

6.4.3. Effets sur les structures vortiques de la couche de mélange

6.4.3.1. Effets sur les structures primaires de la couche de mélange L'étude de la dynamique des tourbillons qui composent la couche de mélange est maintenant réalisée pour les deux positions de mesure le long de l'envergure. Comme précédemment, les champs de vitesse sont reconstruits à partir d'un nombre limité de modes SPOD. Le critère Q à partir des champs de vitesse en moyenne de phase pour un forçage de l'instabilité fondamentale $St_f = St_0 = 0.022$ est présenté à la figure 46.

À la position $z/\theta_i = 25.5$ le forçage produit de forts tourbillons juste en aval du bord de fuite de la plaque de séparation, comme le montre la figure 46b (contours colorés). Ce sont les tourbillons dans le sens de l'envergure qui résultent du mécanisme d'instabilité KH primaire. Comme observé avec le forçage uniforme, ces tourbillons s'atténuent ensuite en aval sans interagir les uns avec les autres, ce qui entraîne une croissance faiblee de la couche. La séparation entre deux tourbillons successifs est $\Delta x/\theta_i \approx 45.5$ qui est la longueur d'onde de l'instabilité fondamentale forcée. Cela confirme les conclusions tirées du calcul SPOD selon lesquelles, bien que les fluctuations soient



Figure 46.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0 = 0.022$ with the spanwise-modulated DBD plasma actuator ($c_{\mu} = 0.009$). Phase-averaged fields at both spanwise stations (a) $z/\theta_i = 0$ and (b) $z/\theta_i = 25.5$ are presented. Contours in (b) represent $Q/Q_{max} = 0.01$ (green) and 0.1 (black) from the corresponding field in (a). The dash-dotted (black) lines represent (from bottom to top) the $y_{0.05}$, $y_{0.5}$ and $y_{0.95}$ traces at the corresponding spanwise stations.

communiquées via un forçage tridimensionnel il est pourtant possible de réorganiser l'écoulement bidimensionnel. À $z/\theta_i = 0$ (figure 46a), le forçage EFD appliqué génère aussi de forts tourbillons (primaires) dans le sens de l'envergure juste en aval du bord de fuite de la plaque de séparation. Comme observé à la position $z/\theta_i = 25.5$, ces tourbillons s'atténuent progressivement sans interaction notable. La figure 46b correspondant à $z/\theta_i = 25.5$, montre également les contours de $Q/Q_{max} = 0.01$ et 0.1 (respectivement en vert et en noir) des champs correspondants en moyenne de phase à $z/\theta_i = 0$. Cette superposition des résultats indique que l'emplacement des tourbillons primaires dans le sens de l'écoulement n'est pas identique en z à un instant donné. Le noyau du tourbillon à $z/\theta_i = 0$ est décalé par rapport à $z/\theta_i = 25.5$ et ce décalage augmente progressivement avec les positions aval. Ceci montre la présence d'une ondulation des structures fluidiques primaires due à la présence des tourbillons secondaires qui a été observée

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Figure 46.: Caption in previous page

par Breidenthal (1981, appelée "wiggle").

6.4.3.2. Effets sur les structures secondaires de la couche de mélange La campagne d'essai a été limitée à des mesures PIV bidimensionnelle. Cependant comme les structures secondaires bouclent autour des noyaux des structures primaires et que leur vorticité est perpendiculaire au plan mesuré, il est alors possible de les observer dans les plans de mesures selon l'envergure. Ainsi, un noyau de la structure secondaire (qui reste faible) est observé au-dessus de chaque structure primaire à la station dans le sens de l'envergure $z/\theta_i = 0$, approximativement le long de la trace $y_{0.95}$. Ce noyau n'est pas visible sur les plans de mesure à $z/\theta_i = 25.5$. Ce sont les signatures des vortex secondaires qui se superposent aux noyaux des vortex primaires (Bernal & Roshko, 1986; Lasheras & H. Choi, 1988; Nygaard & Glezer, 1991). Le renforcement de ces structures secondaires est bien corrélée avec la région de l'écoulement où le taux de croissance de la couche de mélange était significativement plus élevée. Cette croissance plus élevée serait due à un entraînement plus important causé par les tourbillons secondaires dans cette région. Aucune structure tourbillonnaire n'est observée ici passé $x/\theta_i = 300$, ce qui correspond à l'endroit où

la croissance de la couche de mélange s'est temporairement arrêtée. L'absence de structures d'écoulement secondaires à $z/\theta_i = 25.5$ entraîne clairement un schéma de croissance similaire à celui qui consiste à forcer l'instabilité fondamentale en utilisant un actionneur à forçage uniforme.

7. CONCLUSIONS

Dans cette thèse, la capacité des actionneurs plasma à décharge à barrière diélectrique (DBD) pour manipuler différentes configurations d'écoulement dont l'évolution est régie par des mécanismes d'instabilité hydrodynamique est étudiée. Le contrôle de l'instabilité de l'écoulement pivote autour de la transmission de perturbations de faible amplitude afin de manipuler les caractéristiques de stabilité de la configuration d'écoulement, de manière à la contrôler et à réaliser un effet bénéfique le plus large possible. Les instabilités d'écoulement peuvent être classées en deux grandes catégories. Le premier sous-ensemble comprend les instabilités qui s'échelonnent en fonction des longueurs d'onde temporelles et qui dépendent donc de la fréquence. L'excitation de ces instabilités nécessite donc un forçage instationnaire/fluctuant. La deuxième catégorie comprend les ondes d'instabilité qui se produisent dans l'espace et qui sont définies en fonction de leur longueur d'onde/nombre d'onde spatial. La manipulation de ces instabilités nécessite un forçage tridimensionnel modulé dans l'espace. Étant donné l'instabilité inhérente au forçage appliqué par les actionneurs de plasma DBD, ces actionneurs peuvent, sans effort particulier, manipuler la première catégorie d'instabilités. Grâce à la variation de la géométrie de leur construction (en particulier celle des électrodes), un scénario de forçage tridimensionnel (bien que toujours instationnaire) peut également être réalisé avec les actionneurs DBD. Ainsi, l'un des objectif de cette thèse était de vérifier que les actionneurs plasma sont capable de manipuler la deuxième catégorie d'instabilités.

Afin d'étudier les capacités des actionneurs DBD à cet égard, deux actionneurs de configurations géométriques différentes ont été conçus et utilisés. Le premier est l'actionneur couramment utilisé qui exerce une force électro-fluide-dynamique (EFD) instationnaire, bidimensionnelle et uniforme dans le sens de l'envergure. Le second est un actionneur spécialement conçu qui exerce une force EFD tridimensionnelle modulée spatialement dans le sens de l'envergure. Ces actionneurs ont été déployés dans deux configurations d'écoulement qui sont régies par des mécanismes de stabilité hydrodynamique. La première étant la couche limite sur une aile en flèche, dominée par le mécanisme d'instabilité de l'écoulement transversal qui provoque une transition laminaire-turbulente à la surface de l'aile. On tente de contrôler les modes d'instabilité transverses stationnaires qui se manifestent par des tourbillons en co-rotation constante le long de l'envergure de l'aile ce qui nécessite un forçage EFD tridimensionnel modulé dans le sens de l'envergure. La deuxième configuration d'écoulement choisie pour cette thèse est la couche de mélange plane, dont la dynamique est définie par les structures d'écoulement vortical sousjacentes et leur interaction les unes avec les autres. Les tourbillons à grande échelle sont régis par le mécanisme d'instabilité Kelvin-Helmholtz (KH) qui se manifeste dans cette configuration d'écoulement sous forme d'ondes dépendantes de la fréquence. Les instabilités secondaires se manifestent sous la forme de structures d'écoulement tourbillonnaire dans le sens du courant qui s'étendent le long de la couche avec une longueur d'onde spécifique dans le sens de l'envergure, ce qui nécessite à priori une fois de plus l'utilisation d'un forçage spécifique modulé dans le sens de l'envergure.

Bien que les actionneurs DBD aient été largement utilisés pour leur caractère non intrusif, les techniques de construction conventionnelles généralement suivies en laboratoire aboutissent à des épaisseurs d'électrode de plusieurs dizaines de microns, ce qui peut être prohibitif et avoir ainsi un effet déstabilisant sur le champ d'écoulement. C'est particulièrement le cas ici avec la couche limite sur l'aile en flèche qui est extrêmement sensible aux caractéristiques de surface telles que la rugosité et l'épaisseur des éléments. Dans cette thèse, un prototype de technique de fabrication automatisée a été développé pour construire des électrodes métalliques dont l'épaisseur n'est que de quelques micromètres. La technique proposée repose sur la pulvérisation contrôlée par ordinateur de particules d'argent conductrices sur la surface diélectrique. La technique est capable de produire des actionneurs avec des électrodes extrêmement fines. Elle est également capable de fabriquer des électrodes à géométrie complexe afin de générer un forçage tridimensionnel EFD chaque fois que l'application de contrôle de flux le nécessite. En outre, cette méthode est rentable et facile à réaliser en laboratoire. La caractérisation électromécanique de ces actionneurs suggère un fonctionnement conforme aux attentes, la physique de la décharge plasma n'est pas affectée par la modulation spatiale et l'écoulement produit est bel et bien fortement tridimensionnel. De plus, l'installation sur le modèle aérodynamique des actionneurs DBD ainsi produits n'a montré aucune influence sur les caractéristiques de la couche limite qui s'y développe. Ceci démontre le succès de la technique proposée pour atteindre les objectifs fixés lors de sa conception.

En ce qui concerne le forçage exercé pour manipuler les instabilités dans les configurations de flux correspondantes, les actionneurs du DBD n'ont que partiellement réussi. Les actionneurs employés ont fait preuve d'une grande autorité sur la couche de mélange plane. Les actionneurs correspondants ont réussi à exciter les modes d'instabilité prévus et à manipuler les caractéristiques de stabilité de la couche. Dans la couche limite de l'aile en flèche, les actionneurs déployés ont réussi à retarder la transition laminaire-turbulente. En fait, les résultats présentés ici marquent la première réalisation d'un retard de transition sur une aile balavée dans un cadre expérimental avec un contrôle actif d'écoulement basé sur une décharge de type plasma froid. Cependant, la preuve concluante que le contrôle de l'instabilité basé sur le plasma est la cause du retard de transition réalisé n'a pas été observée lorsque l'actionnement modulé dans le sens de l'envergure a été utilisé pour contrôler les instabilités stationnaires transverses dans la couche limite. Le conditionnement de la couche limite par un ensemble d'éléments de rugosité discrets (DRE) semble rendre les actionneurs utilisés faibles et donc inefficaces pour forcer les modes d'instabilité prévus. Par contre, dans le champ d'écoulement de la couche de mélange, on a observé que les forçages de faible amplitude étaient préférables. Ainsi, la réceptivité du champ d'écoulement semble dicter l'efficacité du contrôle d'instabilité déployé.

Une observation importante provient du déploiement du forçage de l'EFD modulé sur les deux configurations d'écoulement. L'exercice de ce forçage sur la couche de mélange plane a révélé sa capacité à manipuler les instabilités KH primaires qui régissent la dynamique bidimensionnelle de la couche. Ceci, tout en excitant les instabilités secondaires prévues le long de la couche. Cette découverte a des implications notables pour le contrôle des instabilités basées sur le plasma dans d'autres systèmes fluides. Elle démontre que si la dynamique tridimensionnelle du champ d'écoulement est influencée par ces actionneurs modulés comme prévu, la dynamique bidimensionnelle de l'écoulement peut également être influencée. D'une part, cela peut être un problème potentiel car la dynamique bidimensionnelle peut être influencée involontairement. D'autre part, si ce n'est pas un problème ou si l'on préfère, le déploiement de ces actionneurs peut être très avantageux car leur fonctionnement nécessite une dépense d'énergie fortement réduite. L'impact positif de cette moindre consommation d'énergie est bien mis en avant dans les études sur la couche limite de l'aile en flèche. Alors que le délai de transition a été réalisé avec un forçage uniforme et modulé dans le sens de l'envergure, l'efficacité de la commande du forçage modulé dans le sens de l'envergure était nettement supérieure à celle de l'actionnement

uniforme correspondant. De même, dans la configuration d'écoulement de la couche de mélange, si l'objectif final est l'amélioration du mélange, le déploiement de ce forçage modulé dans le sens de l'envergure peut aider à atteindre cet objectif avec une très faible dépense énergétique.

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ACKNOWLEDGEMENT

This booklet narrates part of the story that transpired in the past few years, a story involving questions of science while breathing healthy amounts of ozone. There is however another story that runs parallel to this one and is written on the dark side of the page. This pertains to a unique journey involving two different laboratories, countries, administrations, immigrations, languages and an unhealthy amount of bureaucratic procedures, with of course the doctoral thesis but just as the background score. Ultimately, this work is a testament to the fact that whatever the odds be, much can be achieved when we work together. Much like the flow configurations I have dealt with in this thesis, the past years have been a combination of inflectional profiles giving rise to instabilities and me playing the role of plasma, attempting to control them and manage the turbulence caused. Both these storylines and the experiences they brought wouldn't have been as complete and enjoyable if not for a number of beings, whom I would like to acknowledge here.

Firstly, I would like to thank the jury members of my doctoral thesis defence for the time they spent to read this (rather voluminous) booklet to pose a number of questions and have a valuable discussion. A special thanks to the *Rapporteurs* for writing elaborate reports on my thesis in addition. While being grateful for their many positive feedbacks, I will reflect on their critical comments and improve on them in my future endeavours.

As part of this research I had the opportunity to work under the supervision of two people, Marios Kotsonis and Nicolas Benard (ordered alphabetically only). I met Marios during my MSc. thesis at Delft (which I carried out under his supervision as well) and have worked with him for five years now. He is just a guy playing with a lot of expensive toys (his words, just paraphrased). Keeping in-line with that, he always told me to "play with things" in the lab which made me passionate towards experimental research and in fact lead me to undertake this path. His clear and straightforward communication (with both good and bad news) is a trait I have tried to cultivate in myself and will continue to work on.

In Poitiers, my research efforts were scarred with delays and a number of issues at the beginning. A force that kept me grounded, motivated and zealous towards the topic came from Nico. He gave me independence in a number of aspects which allowed me to make this part of research my own and experience the life of a researcher fully. While the journey of generating a mixing layer that is quite canonical and creating the analytical methodology was turbulent, it would not have been possible if not for Nico's constant guidance and support. His very expressive nature and his plethora of unconventional ideas always kept me at the edge of my seat and made this period of life very adventurous. Finally, both of them have guided and supported me not only in scientific pursuit but also personally on a number of occasions. For these and much more, I am greatly indebted to them.

This prologue will be incomplete without thanking my thesis director Eric Moreau. His energy, enthusiasm and curiosity are admirable and thanks to him, I will never process any PIV data before staring at the raw fields for some time. I want to thank Frank Kerherve for greeting me every time with a big smile and for helping me with modifying the wind-tunnel in Poitiers initially. I wish to also thank Jean-Paul Bonnet for accepting our invitation to be part of my defence as a guest and also looking over my doings from time-to-time (both in the lab and in Poitiers' cafés). I also want to thank Markus Kloker for many fruitful discussions while we were working on an article together and also his supportive words to a young researcher when we received a relatively uninspiring review for the same.

Having done a predominantly experimental research, I must thank a number of technicians/engineers for being patient with me in these years. Thank you Stefan and Leo from Delft for allowing me to call the low-speed lab my home for many months. Thanks Christian for understanding my meagre French and modifying the wind-tunnel at Poitiers for me. A special thanks to Patrick and Romain for making the perfect PIV set-ups for me at Poitiers and for also educating me with many practical PIV tricks. I thank Francis and Pierre-Francois for all the rapid help with my computer needs at Poitiers. I must also acknowledge Colette, Catherine, Marie-Line and Nadia for taking care of all my over-complicated administrative procedures in these four years.

In this period, I also had the chance to work closely with a few research beginners like myself who have also become close friends outside the dark alleys of the labs. The first one is of course Jacopo Serpieri, a supervisor at some point and then a colleague, but always a mentor. His sophisticated and indelicate jokes served as tonic during the late-night (/early-morning) experimentation campaigns. I also had the opportunity to "supervise" Marc Hehner when he joined us at Delft to do his MSc. thesis. In reality though, we always worked together and I must thank him for turning a blind eye when I was stealing silver from the lab by spraying it on my clothes. Finally, although it was a small stint, I was lucky enough to also supervise Tom Fridlender, who became a close confidant in Poitiers eventually. While I would like to thank all the researchers at both the labs for the good times together, I want to name a few personalities specially. These would be Alberto, Houhua, Koen and Theo from Delft and David, Etienne, Gad, Melissa, Michelle, Nishant, Thomas and Ugo from Poitiers. The past four years would not have been the same without you all, so Cheers!!!

There is not even a remote possibility of me completing this laborious journey without my folk outside the laboratories. These sapiens of my private life have supported me in every step of the way and have played a major role in my efforts to remain sane and mentally healthy. For being my co-conspirators of innumerable evenings to suddenly becoming chefs and inviting me to a nice dinner while I am toiling away in the labs, I want to thank Arjun, Chiara, Jacob, Magali, Mannat, Marco, Meghana, Nora, Pranav, Reshal, Rimit and Salil. While I can write essays to thank each of you, I prefer to leave it unspoken. I also want to thank many of my other friends and my dearest family spread around the world for the constant support in many manners and your positive thoughts.

Finally, I want to conclude this by thanking two people who mean the world to me. Staying thousands of kilometres away, they have sailed through this journey with me. Always encouraging, accepting my choices and supporting me through thick and thin. For all this and much more that cannot be described with words, I thank my parents. None of this would be the same without the two of you on my side.

Poitiers & Delft last days of 2020

1

INTRODUCTION

The discipline of flow control has garnered considerable attention from both researchers and engineers alike, as a means to improve the efficiency of many fluid-mechanic systems. Flow control pivots around introducing small changes to an existing configuration, through the use of certain devices, to realize a large engineering benefit. Depending on the desired result, flow control can be employed to avert/cause boundary layer separation in wall-bounded flows, delay/advance laminar-to-turbulent transition or to diminish/augment turbulence. It can effectively lead to drag reduction, lift enhancement, mixing enhancement or suppression of flow-induced noise. The different methods and devices employed for flow control in the past has been synthesized in a number of review works (Gad-el-Hak et al., 1993; Joslin, 1998; Gad-el-Hak, 2007; Cattafesta III & Sheplak, 2010).

Flow control can be broadly classified as passive and active. Passive control includes devices such as roughness elements, vortex generators, turbulators, etc. which are permanent installations on to the surface of an aerodynamic body. As such, these devices employ a steady "brute-force" approach towards flow control and do not require any auxiliary power source or energy expenditure. However, a major drawback of passive devices is that they are always operational, notwithstanding the necessity for exercising flow control. Furthermore, given that passive methods almost always involves geometric modifications, the exercised control cannot be regulated based on the flow conditions and thus, they can potentially have a detrimental effect over the flow field.

Active flow control on the other hand involves the use of actuators which impart energy into the flow and thus require some energy expenditure in order to operate them. At the least, these actuators can be turned on and off whenever needed and some operational parameters can be modified depending on the type of actuator. Active control can once again be classified as closed-loop and open-loop. In closed-loop control the employed actuator is part of a feedback loop which also includes certain sensors placed in the fluid. This renders the entire control loop as a dynamic system. In open-loop control, no such feedback signal exists and the actuator is operated by setting its working parameters manually.

A number of devices have been deployed to realize active flow control in the past. Surface heaters can be used to influence the flow by varying the flow viscosity or density gradients (see Liepmann, Brown, et al., 1982). Fluidic actuators that carry out fluid injection (blowing) or suction can have a significant influence on the flow field by modifying the mean velocity profile of the flow. similar to these there are also zero-net-mass-flux or synthetic jet actuators which ingest and expel the fluid consecutively (Glezer & Amitay, 2002). Acoustic waves can be used

to influence the dynamics of shear layers. Vibrating ribbons, oscillating flaps or piezoelectric actuators can be also employed to introduce some periodic fluctuations into the flow field. While these actuators (particularly the mechanical devices) have demonstrated effective flow control in various flow configurations, they have some inherent shortcomings. These actuators are composed of moving parts which undergo wear-and-tear and are thus prone to breakdown. Their installation on aerodynamic bodies is complicated by the necessity for geometric modifications and housing of the driving parts. Furthermore, they add weight and can be a source of noise and other unwanted vibrations.

A device that that is free from these disadvantages and thus has attracted much recognition from the flow-control research community in the past two decades, is the alternating current (AC) dielectric barrier discharge (DBD) plasma actuator. Most of the advantages of these devices stem from the pertinent electrical nature of their operation and their ability to apply the necessary control without the use of any moving parts. They exert an unsteady, fluctuating forcing based on the frequency content of the input electrical waveform and feature a fast response in terms of amplitude and frequency. Finally, being electrically operated, inclusion of these devices in closed-loop control algorithms is also possible. A number of review articles published in the past two decades have synthesized the different flow control applications of these actuators (Moreau, 2007; Corke, Enloe, et al., 2010; Wang et al., 2013; Kriegseis, Simon, et al., 2016; Corke & Thomas, 2018; K. Choi & Kim, 2018). These actuators have demonstrated proficiency in realizing active flow control in a number of occasions in low-velocity, sub-sonic conditions. However, scaling plasma actuators for high Reynolds number flows, which can be encountered in practical applications, is still a challenge due to their low-momentum energy production. A possible method to put these actuators to use in practical flow conditions such as flight Reynolds numbers is to employ their low-momentum-transfer-based forcing for a sub-category of active flow control that can be referred to as *flow instability control*.

In fluid configurations that have an inflectional velocity profile, certain small-amplitude wavy disturbances can produce large effects as demonstrated first by Lord Rayleigh (1880). These small-amplitude disturbances are what is known as flow instabilities. They can be temporal/frequency dependent waves or spatially occurring waves. These waves, in certain bands (of frequency or spatial wavelength) depending on the conditions of the flow, can amplify rapidly and disrupt the hydrodynamic stability of the fluid, thereby causing a large effect. The amplification of these instability waves can be controlled with very low amplitude forcing which affect the stability characteristics of the flow field, thereby producing the necessary outcome. This approach of flow control is in fact more efficient than a "brute-force" approach as the flow dynamics can be effectively governed with lower energy expenditure.

Among many, two flow configurations whose dynamics are predominantly governed by hydrodynamic instability mechanisms prevalent within, are the three-dimensional boundary layer forming around a swept wing of modern aircraft and the plane mixing layer. Effective control of both these flow configurations can have significant economical and more importantly ecological benefits. In swept-wing boundary layers, exercising laminar flow control by manipulating the instabilities within can aid in achieving a delay in laminar-to-turbulent transition which will result in a direct reduction of friction drag (Joslin, 1998). This has the potential to reduce flight operational costs, fuel consumption and emissions which will lead to a more eco-friendly flying machine. On the other hand, the plane mixing layer configuration is representative of a free-shear layer (Brown & Roshko, 1974). As such, manipulating these instabilities and thereby controlling the mixing layer has a very wide range of applications. Additionally, the shear layer has an influence on a number of fluid systems. Thus, understanding and developing capabil-



Figure 1.1.: Schematics showing instabilities occurring in (a) swept-wing boundary layer and (b) plane mixing layer.

ities to control this can potentially aid in improving flow control applications in other fluid configurations.

This shows the importance and necessity to device methods to control the swept-wing boundary layer and the plane mixing layer flow configurations. Given that the dynamics and evolution of both these fluid systems are governed by hydrodynamic instability mechanisms, instability control can be applied over them. While these two flow configurations are very different from one another, the nature of the prevalent instabilities have certain commonalities. The instabilities in both these flow fields can be classified into two categories. As shown in the schematics in figure 1.1, they are either spatially-occurring instabilities, i.e. instability modes that are differentiated based on their spatial wavelength/wavenumber, which necessitates the applied forcing amplitude to be three-dimensional/modulated in the spatial domain. Or they are instabilities that scale according to temporal wavelengths and thus are frequency-dependent. In which case, the imparted forcing must be unsteady or fluctuating. Modulations of the applied forcing in space and time are a common approach for instability control.

1.1. MOTIVATION AND OBJECTIVES OF THE CURRENT THESIS

The core motivation of this thesis is to investigate the capabilities and limitations of DBD plasma actuators in realizing active flow-instability control. These actuators inherently exert an unsteady/fluctuating forcing. Through geometric features of their construction, spatially-modulated, three-dimensional forcing can also be realized. These features of DBD plasma actuation, combined with the low-momentum-transfer-based forcing exerted by these devices make them an ideal candidate to be deployed for active instability control. Apart from scientific pursuit, realization of control of these two flow configurations can have a significant impact on their engineering applications as well. Different DBD plasma actuators that can exert forcing of two-dimensional and three-dimensional nature will be designed and fabricated. After a thorough characterisation of the these actuators, they will be deployed for flow/instability control purposes over the swept-wing boundary layer and the plane mixing layer flow configurations. The results that will be reported and discussed here are predominantly acquired through experimental means. Certain theoretical computations are also carried out in order to support the experimental investigations. Furthermore, both these flow configurations are investigated at sub-sonic flow conditions. The specific objectives of this thesis are detailed below.

1.1.1. DBD plasma actuators

The core of this thesis being the forcing amplitudes imparted by DBD plasma actuators, certain tasks are necessary for their effective implementation in the respective flow configurations. Two types of DBD actuators are deployed on both these flow fields. The first being a spanwise-uniform actuator that imparts a two-dimensional time-modulated body force and the second being a spanwise-modulated actuator that exerts three-dimensional forcing that is spatially modulated. In order to realize effective instability control with these actuators, two exercises were recognised to be of utmost importance and hence were undertaken. These are listed below.

- The deployed actuators must comply to strict thickness and roughness constraints so as to avoid the mere presence of the actuator delivering a destabilising effect on the flow. This is a general concern related to employing these actuators for flow/instability control applications. To circumvent this issue, a necessity arises to formulate a fabrication technique that produces such actuators. Furthermore, given the desire to impart three-dimensional forcing, the proposed technique must be capable of producing electrodes with intricate geometric features while adhering to strict thickness and roughness constraints. Indeed, this task will be incomplete without a proper validation of the fabricated actuators, which will also be carried out in the course of this thesis.
- The second task is related to the spanwise-modulated actuators. Firstly, a proof-of-concept is necessary to demonstrate the three-dimensionality of the induced flow. Secondly, given the inherent sensitivity of control performance to spanwise-uniform and modulated actuation will be addressed in both flow configurations, it necessitates a direct comparison of these two types of forcing, which further relies on quantification of the performance of spanwise-modulated actuators. Thus, this is one of the key objectives of the current thesis.

1.1.2. Swept-wing boundary layer

This thesis aims at controlling the cross-flow (CF) instability mechanism that is the cause of laminar-to-turbulent transition in three-dimensional swept-wing boundary layers. A control strategy that has garnered much attention from the boundary-layer control community is the use of discrete roughness elements (DRE) at the leading edge to cause a favourable mean-flow deformation along the wing's span (Saric, Carrillo, et al., 1998). This was later generalized as upstream flow deformation (UFD) to encompass the use of any device that can cause this meanflow deformation (Wassermann & Kloker, 2002). While this strategy did show positive results in a few experimental and numerical efforts, it is not very robust due to its dependence on a number of parameters. An alternate strategy, mainly from numerical simulation efforts (Messing & Kloker, 2010; Dörr & Kloker, 2015b), is based on attacking the cause of the instability, the CF velocity component in the boundary layer, and is called base-flow modification (BFM). However, this strategy has never been validated experimentally. In recent years, a number of numerical studies and a few experimental efforts that employ DBD plasma actuators as flow control devices in three-dimensional boundary layers have been carried out. Numerical simulations do demonstrate a positive outcome with respect to transition delay, but the same has not been achieved through experimentation. Applying three-dimensional EFD forcing with DBD actuators to realize the UFD control strategy in a swept-wing boundary layer has been experimentally investigated recently Serpieri, Yadala, et al. (2017) and K. Choi & Kim (2018). While the actuators successfully conditioned the boundary layer with the desired stationary

instability mode, it also resulted in the amplification of unsteady CF instabilities which lead to a promotion of transition.

Thus, with regards to the swept-wing boundary layer, the UFD and BFM control strategies will be realized with plasma actuation. The ultimate goal being the realization of a delay in laminar-to-turbulent transition. The different control strategies will be investigated under identical flow conditions so as to better compare the achieved control authority and efficiency with these two control strategies. If successful, this will mark the first experimental effort to achieve delay of CF instability dominated transition in swept-wing boundary layer. The four objectives of the current thesis in this regard is as follows.

- The first objective is to fabricate novel DBD plasma actuators with very thin electrodes (micrometric scales) that can apply EFD forcing at a high frequency so as to mitigate the amplification of unsteady primary travelling CF instabilities.
- The second goal is to realize the UFD control strategy through DBD actuators that generate a spanwise-modulated forcing. Furthermore, with certain forcing configurations (actuator design and operating conditions), the spanwise-modulated body force being exerted to realize the UFD strategy can be expected to perform a BFM action locally. However, the effect of this action and how the two strategies contribute to controlling transition in such a situation is not well understood and as to date, not been explored experimentally in a systematic manner. This task will also be undertaken in the course of this thesis.
- The third goal is to realize and experimentally study the the BFM control strategy using DBD plasma actuators. Prior to the experimental investigation, a preliminary theoretical study will be performed to investigate the stability of the boundary layer, under the influence of the EFD body force produced by the actuator, using a simplified numerical model based on linear stability theory. Following the predictions of the model, experimental demonstration of the plasma-based base-flow modification will be achieved by installing DBD plasma actuators that produce a two-dimensional spanwise-uniform EFD body force, close to the leading edge of a swept-wing model.
- Finally, this thesis does not only provide the first experimental proof-of-concept of the BFM control strategy, but also is the first comparison of the two control strategies at identical flow conditions. Thus, it provides favourable conditions to investigate the collaboration of these two strategies. Furthermore, these two strategies will also be compared in terms of the achieved control efficiency based on their effect on the total friction drag.

1.1.3. Plane mixing layer

The next part of this thesis concerns the application of DBD plasma actuators on the plane mixing layer. This flow configuration is firstly characterised by large-scale spanwise vortices that arise for the *Kelvin-Helmholtz* instability mechanism (Brown & Roshko, 1974). The dynamics and interaction of these vortices have a direct impact on the growth of the mixing layer. Secondly, the secondary instability mechanism leads to the formation of hair-pin vortices that loop around the spanwise vortices. The interaction of these two vortical structures results in the production of small-scale turbulence and enhanced mixing at molecular scales. Mixing layer is the simplest representation of a free shear layer which is prevalent in every type of fluid flow. Thus, analysing the effect of DBD actuator forcing on the mixing layer holds significant importance to improve the flow-control applications of these actuators. The impact of the spanwise-uniform and modulated plasma forcing on the hydrodynamic stability characteristics of the flow field will be experimentally investigated. A number of analytical tools will be generated and applied on the acquired experimental data in order to discern the influence of manipulating the flow field's stability characteristics over its evolution and dynamics. The objectives of the current research in this regard are as follows.

- The first goal is to assemble an experimental set-up to realize a mixing layer arising due to the interaction of two streams with different free-stream velocities. This also necessitates proper characterisation of the ensuing mixing layer flow field by comparing it with previous experimental efforts.
- The second goal involves generation and validation of analytical tools and formulating a methodology to study this flow configuration. Measurement of flow velocity will be obtained from time-resolved particle image velocimetry. This experimental dataset provides the basis for the first spatio-temporal analysis of the mixing layer flow field. The identified analytical tools include linear stability analysis, power spectral density to investigate the spectral content within the flow field, spectral proper orthogonal decomposition to obtain the spatial organisation of the fluctuations and vortex identification schemes to observe the dynamics and interactions of the vortical flow structures within the flow field.
- The third goal is to design and fabricate the DBD plasma actuators to apply the necessary EFD forcing to excite different instabilities and to characterise their performance. These actuators will be constructed on the splitter plate and will exert the intended forcing into the boundary layer ensuing on the high-velocity side of the splitter plate, just upstream of its trailing edge.
- This is followed by inspecting the impact of exciting different KH instability modes through spanwise-uniform forcing by the corresponding DBD actuator. This is essentially achieved by taking advantage of the unsteady nature of the forcing applied by DBD actuators. Furthermore, the impact of different operating parameters of the DBD actuator will also be studied.
- The final objective of this part of the thesis is to carry out a preliminary study to manipulate secondary instabilities in the mixing layer. This is done by imparting spanwise-modulated forcing with the corresponding DBD plasma actuator.

1.2. THESIS OUTLINE

This thesis is organised as follows. Chapter 2 contains a brief review of the research efforts carried out in the past to better understand the different concepts investigated in the current thesis. It includes a brief description of the working principle of DBD plasma actuators. The evolution of laminar-to-turbulent transition dominated by hydrodynamic instability mechanics in a sweptwing boundary layer and the different control strategies devised to control this instability to realise a delay in transition will then be introduced. The different instability mechanisms that define the dynamics of vortical flow structures within the plane mixing layer flow configuration is also discussed. Finally, specific objectives of this thesis are also presented.

The methodology followed and the tools employed in this thesis to investigate the correspond-

ing flow configurations is introduced in chapter 3. This chapter lists the different experimental test facilities used during the course of this thesis, the measurement systems deployed and finally the different tools employed in this thesis in order to analyse the acquired data.

Chapter 4 is dedicated to DBD plasma actuators. A novel technique to fabricate DBD actuators' electrodes with intricate geometries and thickness in micrometric scale is proposed and the fabricated actuators are validated. The geometric configurations of the different DBD plasma actuators that will be employed in this thesis are then introduced. The spatially-modulated actuator geometry intended to apply three-dimensional forcing is then characterised thoroughly. This is followed by introducing the details of the different DBD plasma actuators that will be employed for flow/instability control in the swept-wing boundary layer and the plane mixing layer flow configurations.

The efforts carried out to investigate the swept-wing boundary layer is reported in section 5. The investigated flow conditions are first introduced followed by the results from stability analysis of the boundary layer to predict certain information regarding the ensuing instabilities. The condition of the baseline (unforced) boundary layer is then introduced. This is followed by the presentation of results acquired while realizing different control strategies with plasma-based flow control. Finally, an estimate of the achieved control efficiency with the different strategies based on its impact on the total friction drag is made in order to compare them with one another.

Chapter 6 consists of all the acquired results pertaining to the investigations carried out on the plane mixing layer flow configuration. As a first step, thorough characterisation of the unforced mixing layer was undertaken. The obtained results are analysed and compared with previous experimental efforts. This is followed by employing the two different DBD plasma actuators to manipulate the underlying instability mechanisms and analyse its impact on the the flow field.

Finally, a few concluding remarks and perspectives for the future are given in chapter 7.

2

LITERATURE REVIEW

This chapter is dedicated to introducing a few fundamental principles and concepts regarding the three main aspects of this thesis, by reviewing past research efforts on these topics. The working mechanism and characterisation of DBD plasma actuators, which is employed in this thesis to exert the necessary flow control is discussed in section 2.1. An introduction to boundary-layer stability and laminar-to-turbulent transition with a focus on the cross-flow instability prevalent in swept-wing boundary layers is synthesised in section 2.2. The instability mechanisms that govern the evolution and dynamics of vortical flow structures in the plane mixing layer flow configuration is considered in section 2.3.

2.1. DBD PLASMA ACTUATOR

Plasma, named after the Greek word $\pi\lambda\dot{\alpha}\sigma\mu\alpha$ which means 'moldable substance', is generally considered to be the fourth state of matter. When sufficient energy is supplied to a fluid, the electrons in the outermost shell of the atoms break free, generating electrons and ions that coexist in a charge equilibrium. This ionised gaseous state was named *plasma* by Nobel laureate Irving Langmuir (Langmuir, 1928). Plasma in the stars and in the intergalactic space between them compose a significant part of the visible universe, making it the most abundant form of ordinary matter. Ozone generation and degeneration, surface treatment, arc welding, neon signs, fluorescent lamps, vacuum sputtering of thin films and pollution control are a few industrial applications of plasma. From understanding the beginning and working of our universe to manufacturing integrated circuits and flat-screen televisions to nuclear energy, plasma physics has found its place in many areas of science and technology today.

Application of plasma as an active flow control method emerged at the very end of the 20^{th} century. This can mainly be attributed to two research groups. The first comprising of researchers at Université de Poitiers, France and University of Buenos Aires, Argentina who focused on using direct current (DC) corona discharges for flow control. This is a discharge that usually occurs between a sharp needle or a thin wire and a metallic plate. When the needle or wire is supplied with a DC power, it produces an electric wind or motion of charged particles between this high-voltage tip and the grounded plate, which is used for flow control purposes. The second group was from University of Tennessee, USA who perfected and developed a new kind of surface plasma discharge. Through the pioneering work of Roth et al. (1998), they illustrated the ability of a non-thermal surface plasma to generate a thin wall jet. This is the device



Figure 2.1.: (a) Schematic of a typical DBD plasma actuator, reproduced from Kotsonis et al. (2011). (b) Photograph of the plasma discharge, reproduced from Forte et al. (2006).

that is currently referred to as a dielectric barrier discharge (DBD) plasma actuator. The most often used design of DBD plasma actuators is composed of two flat-plate metallic electrodes that are flush mounted on either sides of a dielectric substrate as shown in figure 2.1a. Dielectrics typically used include PMMA, Teflon, glass, ceramics or polyimides such as Kapton for instance, having a thickness ranging between tens of micrometres to a few millimetres. The width of the electrodes are a few millimetres and they are asymmetrically mounted on either sides of the dielectric with no gap (or overlap) or a few millimetres of it.

These actuators work on the principle of air ionisation. When supplied with sufficient electrical energy, electrical breakdown ensues, generating ions and free electrons. These ions move in the produced electric field and collide with neutral air particles in their vicinity, thus transferring momentum and generating a Coulombian electro-fluid-dynamics (EFD) body force. This body force gives rise to a weak jet that is parallel to the dielectric wall, whose velocity is typically less than 10 m s⁻¹. In the past two decades a vast amount of work has been carried out by different research groups on the working of these actuators and their capability for flow control. The basic principles, working mechanism and the application of these actuators for different flow control applications have been synthesized well in a number of review articles produced in the past two decades (Moreau, 2007; Corke, Enloe, et al., 2010; Wang et al., 2013; Benard & Moreau, 2014; Kotsonis, 2015; Kriegseis, Simon, et al., 2016; Corke & Thomas, 2018; K. Choi & Kim, 2018).

The presence of the dielectric barrier leads to charge accumulation on its surface when the actuator is supplied with electrical energy. This results in the reduction of the local electric field, thus quenching the plasma discharge. Theoretically, this can be overcome by continuously increasing the voltage of the input signal. However, this method is not viable practically due to constraints of the power supply unit and the breakdown of the dielectric substrate. A second solution to overcome the quenching is to reverse the polarity of the input signal which necessitates the application of alternating current (AC) signal as input. This results in lesser power consumption while sustaining the plasma discharge. The frequency (f_{ac}) of the AC waveform ranges from 100 Hz to a few tens of kilohertz and the voltage amplitude (V_{ac}) can be a few kV to about 50 kV depending on the thickness and properties of the dielectric substrate. Typically, one electrode is supplied with a high-voltage AC signal (referred to as active electrode) while the other is grounded. This actuator configuration results in the generation of a plasma discharge on either sides of the dielectric, at the edge of the electrodes that is in the vicinity of an edge



Figure 2.2.: Typical discharge current and fast-imaging of the developing plasma along one sinusoidal AC input waveform. Reproduced from Benard & Moreau (2012).

of the other electrode. However, as the actuator is mounted on to a surface in all applications, one electrode is encapsulated from the surrounding air by the dielectric substrate. Hence, once the plasma ignition voltage (V_0) is passed (i.e. electrical breakdown has occurred), the plasma discharge is observed only on the air-exposed side of the dielectric substrate, a photograph of which is depicted in figure 2.1b.

2.1.1. Discharge current and plasma extension

The discharge current of a DBD plasma actuator is generally measured by connecting a noninductive shunt resistor in series with the grounded electrode. The discharge current has two distinct regimes, one during positive-going half of the AC cycle while the other during the negative-going half cycle (see for example Enloe, McLaughlin, VanDyken, Kachner, Jumper, & Corke, 2004; Pons et al., 2005; Forte et al., 2007; Benard & Moreau, 2012; Debien, Benard, & Moreau, 2012). While the discharge seems to be a uniform glow to the human eye as in figure 2.1b, it is composed of numerous micro-discharges that exist for a few nanoseconds (Fridman et al., 2005; Laurentie et al., 2009). Concurrent to the current, the spatial structure of the plasma discharge is different during the two halves of the AC cycle as well. Investigating the structure of the plasma discharge provides more insight into streamer formation and the physical mechanisms that govern the operation of plasma actuators. This is carried out by optical imaging of the ionization process using an intensified charged-coupled device (iCCD) camera (see Hoskinson et al., 2008; Benard & Moreau, 2012; Debien, Benard, & Moreau, 2012). These cameras are capable of acquiring images of fast processes with temporal gate-width in the order of a few nanoseconds which are the typical time scales of plasma discharges.

As observed figure 2.2, during the positive-going half of the AC cycle, the discharge current is composed of many high-current peaks which is a consequence of streamer formation in the discharge. The amplitude of these peaks is a function of a number of parameters such as the geometry and construction of the actuator, performance characteristics of the power supply and environmental conditions among others. During this half of the AC cycle, the plasma layer is composed of streamers stochastically spread along the edge of the air-exposed electrode and hence is sometimes referred to as *streamer regime* (A and B in figure 2.2). Each of these microdischarges have a bright spot at the electrode's edge (corona zone) and then extend towards the dielectric substrate with very erratic propagation paths. These streamers are the conductive channels through which electrons move in the plasma layer.

During the negative-going half of the AC cycle, the current signal shows negative peaks whose amplitudes are one order of magnitude lower than that observed in the positive-going half cycle. Correspondingly, the discharge appears as a diffuse jet-like structure. The plasma layer is composed only of the corona regions that extend from the edge of the air-exposed electrode to dielectric surface. The micro-discharges bear a plume shape as seen in C and D of figure 2.2 and hence this half-cycle is generally referred to as the *glow regime*. Additionally, the length of each micro-discharge remains approximately the same. The extent of the plasma discharge along the dielectric surface increases linearly with the applied voltage amplitude (Enloe, McLaughlin, VanDyken, Kachner, Jumper, & Corke, 2004; Forte et al., 2007; Kriegseis, Grundmann, et al., 2011).

2.1.2. Power consumption

Electrical power is an important parameter to be considered when characterising plasma actuators as it directly affects the efficiency of the actuator and the flow-control system as a whole. While the voltage (V(t)) applied across the DBD actuator is usually available from the power supply unit itself or can be measured using a high-voltage probe, measurement of instantaneous current through the actuator is complicated by the fast micro-discharges in the plasma jet region and the resulting current peaks discussed in previously. A simpler method is based on the work of Manley (1943), which was adapted for the surface DBD actuator by Pons et al. (2005) and perfected by Kriegseis, Grundmann, et al. (2011). It is often used due to its reliability, relatively higher accuracy and ease of implementation (Ashpis et al., 2012). This method involves connecting a monitor capacitor (C_M) in series between the actuator and the ground. The voltage drop $(V_M(t))$ across this capacitor is measured and the charge in the capacitor is calculated by

$$Q_M(t) = C_M V_M(t) \tag{2.1}$$

The energy consumed by the DBD plasma actuator is then estimated by calculating the area under the charge-voltage $(Q_M(t) - V(t))$ Lissajous curve for a full AC cycle (see figure 2.3a). The average power consumed (P_{avg}) is then obtained by dividing this energy by the time period of the AC cycle (T_{ac}) as shown in equation 2.2.

$$P_{avg} = \frac{1}{T_{ac}} \oint V(t) dQ_M \tag{2.2}$$

The power consumed by the DBD plasma actuator is mainly due to that dissipated by the discharge in the form of heat and dielectric losses. With voltage amplitude, the input power variation is different for thick and thin dielectrics. For a thick dielectric (see figure 2.3c), the average power consumed varies according to the relation $P_{avg} = f_{ac} \cdot A \cdot (V_{ac} - V_0)^2$ (Pons et al., 2005; Forte et al., 2007; Jolibois & Moreau, 2009). The proportionality constant A (in W kV⁻²) depends on the dielectric properties, geometry of the electrodes and environmental conditions. In the case of thin dielectrics (see figure 2.3d), the consumed power variation with voltage amplitude is well approximated by $P_{avg} \propto = V_{ac}^n$ (Enloe, McLaughlin, VanDyken, Kachner, Jumper, & Corke, 2004), with n ranging from 2 to 3.5. The difference between thick and thin



Figure 2.3.: (a) A sample of a charge-voltage Lissajous curve, reproduced from Pons et al. (2005). (b) and (c) show the variation of consumed power as a function of input AC waveform frequency and voltage amplitude respectively for an actuator constructed on a thick dielectric, reproduced from Forte et al. (2007). (d) Power consumed as a function of voltage amplitude for a thin dielectric, reproduced from Enloe, McLaughlin, VanDyken, Kachner, Jumper, & Corke (2004).

dielectrics is attributed to the characteristics of the discharge. The power dissipated due to dielectric losses is proportional to the AC frequency and the square of the voltage amplitude. Thus, increase of power with respect to voltage amplitude according to a power law greater than two is a result of the discharge becoming filamentary, which for a thin dielectric occurs at lower voltages compared to thick dielectrics. The average power consumed increases linearly with the frequency of the input AC waveform provided the discharge is not filamentary as demonstrated by Forte et al. (2007) (see figure 2.3b). This linear increase is mainly attributed to the dielectric losses which increase linearly with AC frequency.

2.1.3. Induced jet

DBD plasma actuators generate an EFD body force through air ionisation, as mentioned previously. This essentially transfers momentum to the flow field, thus producing a jet, sometimes referred to as ionic (or electric) wind. Initially, measurements of the induced jet velocity were acquired using a glass pitot tube (total pressure probe) connected to a manometer (see for example Pons et al., 2005; Forte et al., 2007; Jolibois & Moreau, 2009). These probes can be fixed to a displacement system to obtain the spatial distribution of velocity within the jet. While this method is simple, fast and highly accurate in measuring the velocity within the induced jet, it is intrusive and only the horizontal (wall-parallel) component of velocity can be measured. Furthermore, the low sampling rates of this measurement prohibits the acquisition of



Figure 2.4.: (a) Schlieren image of a starting vortex generated by a DBD plasma actuator, reproduced from Moreau et al. (2008). (b) Time-averaged ionic wind generated by a DBD plasma actuator with a 3 mm thick dielectric substrate ($V_{ac} = 22$ kV, $f_{ac} = 1.5$ kHz), reproduced from Debien, Benard, & Moreau (2012).

time-resolved data. A measurement technique that provides velocity data within the induced jet with excellent temporal resolution is hot-wire anemometry. However, this method to quantify the induced velocity bears some inherent limitations as the measurements can be affected by the electromagnetic interferences due to the actuator. Furthermore, generation of an electric arc between the air-exposed electrode edge and the hot-wire is a possibility if the probe is very close to the actuator and would destroy the probe and the anemometer. Due to these limitations, the use of optical measurement techniques such as particle image velocimetry (PIV, see for example Kriegseis, Grundmann, et al., 2011; Kotsonis et al., 2011; Debien, Benard, & Moreau, 2012) have been employed in the past decade to study the topology of the induced jet. This method is non-intrusive, provides at least two components of velocity and measurement of the induced jet with both high spatial and/or temporal resolutions can be realized.

These measurements have revealed much about the topology of the flow induced by DBD plasma actuators. The flow is accelerated as soon as the DBD actuator is operational. This acceleration lasts for tens of milliseconds and generates an initial vortex (see figure 2.4a). Post this acceleration stage, a thin wall jet is produced from the air-exposed electrode edge, along the grounded electrode, which is as depicted in figure 2.4b (also see figure 2.6a). This induced flow is parallel to the dielectric surface, with thickness of a few millimetres. Additionally, a weak suction effect is observed close to the air-exposed electrode edge (just upstream) which induces a vertical velocity component that is oriented towards the dielectric surface. The magnitude of wall-normal velocity component. Finally, the time-resolved PIV acquisitions of the induced flow have revealed that the actuator produces a jet that fluctuates at the same frequency as the input AC waveform. Both AC half-cycles (i.e. positive- and negative-going) generate a horizontal velocity. However, the negative-going half cycle induces more velocity compared to the positive-going half (see changes in velocity during one AC cycle in 2.5b) which leads to a fluctuating jet.

In many situations, amplitude controlled perturbations at a specific frequency is required for flow control. Based on the frequency response of the flow, it might be more effective to apply amplitude modulation over the AC carrier signal. An analysis of the effect of different kinds of AC signal modulation can be found in Benard & Moreau (2010). Here, the effect of burst modulation of the AC carrier frequency will be discussed briefly. The burst-modulation essentially



Figure 2.5.: (a) Variation of DBD plasma actuator power consumption with burst-modulation duty cycle $(f_b = 10 \text{ Hz}, f_{ac} = 1 \text{ kHz})$, reproduced from Benard & Moreau (2009). (b) Time-resolved velocity at a point in the ionic wind when the DBD plasma actuator is operated at a burst modulation $(f_b = 150 \text{ Hz}, 50\% \text{ duty cycle}, f_{ac} = 1.5 \text{ kHz})$, reproduced from Benard & Moreau (2010).

turns the plasma on and off at a particular frequency, the duration of plasma discharge within one cycle being dictated by the duty cycle. The power consumed by the DBD actuator is a linear function of the burst-modulation duty cycle as shown in figure 2.5a, reproduced from Benard & Moreau (2009). The power spectral density of the input burst modulated AC voltage signal computed by Benard & Moreau (2010) showed frequency content not only at the AC carrier frequency and the burst-modulation frequency, but also at harmonics of the burst-modulation frequency.

An example of the evolution of the ionic wind velocity is as shown in figure 2.5b. During the on phase, the induced velocity increases. If the number of AC cycles in this phase is adequate, it reaches maximum time-averaged velocity and fluctuates about it (not the case in the shown figure however). During the off phase, the flow essentially relaxes. Therefore, if the off phase is too small/short for the flow to completely relax, a lesser number of AC cycles is necessary in the next burst-modulation cycle to reach the maximum velocity. This off period dictates the amplitude of the imposed fluctuations. Therefore, as the off period reduces, either by increasing the burst-modulation duty cycle or frequency, the level of fluctuations within the induced walljet decreases. Finally, Benard & Moreau (2010) observed that the frequency content of the induced flow field was similar to that of the electrical signal. Although, energy of fluctuations at the harmonics of the burst frequency is low initially and increases downstream.

2.1.4. Body force and integrated thrust

Production of an electro-fluid-dynamic (EFD) body force is the driving mechanism that culminates in the DBD plasma actuator inducing a wall-jet. The simplest method to measure the time-averaged spatially integrated value of the EFD body force, which is commonly referred to as generated thrust, is to employ a high precision laboratory balance or load cells (see for example Debien, Benard, & Moreau, 2012). The actuator is mounted vertically on a balance such that the induced jet is oriented upwards (away from the balance). This measurement is based on Newton's 3^{rd} law of motion. When the actuator is operational, it applies the EFD force on the flow inducing a jet. As a consequence, the flow applies an equal reaction force on



Figure 2.6.: (a) Control volume (dashed rectangle *abcd*) over time-average induced velocity field for to compute thrust generated through the momentum balance, reproduced from Kotsonis et al. (2011). (b) and (c) Variation of wall-parallel thrust generated as a function of voltage amplitude and average power consumed respectively, reproduced from Debien, Benard, & Moreau (2012).

the actuator. This force being oriented downward will add to the load the actuator applies on the balance. Thus, by comparing the load measured when the actuator is off and on, the thrust generated is estimated. It is worth noting that the measured values correspond only to the force in the wall-parallel direction. Furthermore, the acquired measurements are not the total force generated by the DBD actuator, but the integrated EFD body force along the wall-parallel direction (T_x) subtracted by frictional forces (τ_w) incurred due to the presence of the dielectric substrate $(T_x - \tau_w)$.

A second method to estimate the thrust generated is the momentum balance approach (see Kotsonis et al., 2011). This is computed from the captured time-averaged velocity fields of the induced jet through PIV. A control volume is defined on the time-averaged velocity field as shown in figure 2.6a. The momentum flux along the boundaries of this control volume is used to estimate the thrust generated by the DBD actuator, both in the wall-parallel and wall-normal directions, according to equation 2.3.

$$T_x = T_x^F + \tau_w = \rho \int_{da} u^2 \cdot dy + \rho \int_{cd} uv \cdot dx - \rho \int_{bc} u^2 \cdot dy$$
(2.3a)

$$T_y = \rho \int_{da} uv \cdot dy + \rho \int_{cd} v^2 \cdot dx - \rho \int_{bc} uv \cdot dy$$
(2.3b)

Here, ρ is the density of air, *abcd* represents the vertices of the control volume, (u, v) are the components of the induced velocity in the wall-parallel and wall-normal directions, (T_x, T_y) are



Figure 2.7.: (a) Temporal evolution of the different NS terms governing the development of the induced flow in the wall-parallel direction, reproduced from Kotsonis et al. (2011). (b) Wall-parallel component of the mean EFD body forced distribution (F_x in N m⁻³), reproduced from Benard, Debien, et al. (2013).

the computed thrust along the wall-parallel and wall-normal directions respectively. It is worth noting that the computed wall-parallel thrust contains not only the contribution of the DBD actuator's EFD body force (T_x^F) but also the shear force (or skin-friction) between the induced flow and the dielectric surface. Furthermore, pressure gradient due to the acceleration of the flow by the DBD actuator is not readily available from the velocity data obtained from PIV. Thus, the control volume boundaries must be set far away from the bulk of the EFD body force, so that the pressure can be considered to be uniform and constant.

Similar to the power consumed, the thrust generated too increases with applied AC voltage amplitude according to a power law $T_x \propto (V_{ac} - V_0)^2$, as demonstrated by Debien, Benard, & Moreau (2012) (see figure 2.6b). As a function of power, the thrust generated increases linearly as shown in figure 2.6c. This has been observed by a number of works on different actuators (Enloe, McLaughlin, VanDyken, Kachner, Jumper, Corke, et al., 2004; Thomas, Corke, et al., 2009; Kriegseis, Grundmann, et al., 2011; Debien, Benard, & Moreau, 2012). A deviation of the thrust generated from this linear increase signifies the onset of a filamentary discharge. In this situation the power consumed by the actuator increases with very little change to the thrust generated. The slope of the thrust-power plot in the linear rise region gives the effectiveness of the plasma actuator in performing electro-mechanical conversion.

Determining the spatio-temporal distribution of the EFD force experimentally, holds great value. It not only helps understand the working mechanism of DBD actuators better, but also aids in modelling and designing plasma actuators for flow control applications. Generally, the body force distribution is estimated by solving the Navier-Stokes (NS) equation using the velocity vector fields acquired through PIV (Kotsonis et al., 2011; Benard, Debien, et al., 2013).

$$\frac{F}{\rho} - \frac{\Delta p}{\rho} = \frac{\partial U}{\partial t} + U \cdot \nabla U - \nu \nabla^2 U$$
(2.4)

The full Navier-Stokes equation is as shown in equation 2.4. Here, U represents the acquired velocity data, p is the static pressure, ν is the kinematic viscosity of the fluid, ρ is its density and F is the generated EFD body force. The unsteady acceleration, convection and diffusion terms on the right side of the equation can be computed conveniently using the velocity data. The challenge faced here is, the pressure and the EFD force terms on the left are both unknown

and cannot both be determined with only one equation. Generally, it is assumed that the pressure gradient is negligible compared to the body-force term (Maden et al., 2013). Benard, Debien, et al. (2013) in fact demonstrated that the pressure term on average does not contribute significantly to the total body force. This simplifies the equation and leads to a rough estimation of the body force. However, the magnitude of pressure gradient is comparable to the other terms of the Navier-Stokes equation (see figure 2.7a). Thus, to get an accurate estimate of the EFD body force produced, the pressure term must be resolved. Recently, Kuryachii et al. (2018) came to a similar conclusion through their efforts to model the wall-jet induced by a DBD plasma actuator.

The shape of the EFD body force field is as shown in figure 2.7b. It is found to be an elongated ellipse with maximum force at the air-exposed electrode edge. Furthermore, increasing the voltage is observed to increase the extent of the body force distribution along the wall-parallel direction consequent to the increasing extent of the plasma discharge itself. Finally, a comment has to be made on the temporal dynamics of the body force. Benard, Debien, et al. (2013) synchronized the operation of the actuator and the low-speed (high spatial resolution) PIV acquisition to obtain a phase-locked measurement of the induced flow field. They carried this procedure at different phases along the AC cycle to resolve the unsteady component of the induced jet (similar to Debien, Benard, & Moreau, 2012). The EFD body force was then computed at each of these phase-averaged fields. While the orientation of the force was always along the wall-parallel direction (positive), the magnitude of the force during the negative-going half of the AC cycle was significantly higher compared to the positive-going half. This confirmed the unsteadiness of the forcing exerted by a DBD plasma actuators.

2.2. SWEPT-WING BOUNDARY LAYER

The year 1905 marked, among other things in fundamental physics, the seminal work of Ludwig Prandtl, which introduced the concept of boundary layer. It is the thin region of fluid flow close to the surface of any body where the velocity reduces from the value in the external inviscid flow to relative rest (no-slip condition) at the wall. Thus, the effect of viscous forces is non-negligible when compared to the inertial forces of the flow within the boundary layer. As such, the dynamics of the flow in this region has a significant influence on the aerodynamic performance of any given body. Hence, a tremendous amount of research efforts have been performed in the last century to better understand this relatively small region of flow.

Like any flow, boundary layers can be either laminar or turbulent. Laminar boundary layers are characterised by parallel layers of fluid flow that exchange momentum between each other through viscous shears. No particle motion perpendicular to the direction of this parallel flow occurs. On the other hand, turbulent boundary layers are composed of chaotic motion of particles. These boundary layers are characterised by very strong velocity gradients close to the wall compared to its laminar counterpart. Thus, turbulent boundary layers lead to higher levels of skin-friction. Boundary-layer transition is the process through which an initially laminar boundary layer transform to a turbulent state. As such, understanding the mechanisms that cause this transition and possibly controlling them holds utmost importance as it can lead to improving the efficiency of aerodynamic bodies through the reduction of friction drag.

The second half of the 20^{th} century marked a new beginning to the aviation industry. The induction of jet engines resulted in aeroplanes being capable of flying at higher speeds. As



Figure 2.8.: The Airbus A321neo aircraft with a backward swept wing.

aircraft approach the speed of sound, pressure waves are generated over the wings which increase wave drag, thus negating the beneficial effects of the jet engines to a certain extent. A solution to overcome this problem is to employ swept-wings, as is the case with a majority of modern aircraft. The geometry of this wing is such that it angles backward (occasionally forward) from its root at the aircraft's fuselage (see figure 2.8), rather than extending in a straight sideways direction. While the aircraft itself moves at transonic speeds, the aerodynamic phenomenon over the wing surface is lower as it is a function of the flow component perpendicular to the leading edge. Thus, the critical Mach number of the aircraft where shock waves appear is increased, allowing the aircraft to travel at higher speeds.

Four basic instability mechanisms can contribute to transition in three-dimensional boundary layers ensuing on swept wings. These are Görtler instability, attachment-line contamination and instability, streamwise instabilities related to Tollmien-Schlichting (TS) mechanism and crossflow instability. The Görtler instability is caused due to concave surfaces and can be avoided by using appropriate profile design (Saric, 1994). Attachment-line contamination is only noticed when the leading-edge radius is above a certain critical value and adhering to this limit helps avoid this instability. The streamwise TS instabilities can be controlled using favourable pressure gradients and minimizing the extent of the pressure recovery region (Reed, Saric, & Arnal, 1996). However, favourable pressure gradients that are used to stabilize the TS instabilities destabilize the cross-flow instability in swept-wing boundary layers (Saric, Reed, et al., 2003).

2.2.1. Cross-flow instability

At the swept-wing's leading edge, where the pressure gradients are high, both the surface and flow streamlines are highly curved. These, with the wing sweep, lead to inviscid streamlines deflecting towards the wing-tip for a backward-swept wing and towards the wing's root for a forward-swept wing. Close to the aerodynamic body, a force equilibrium is established along the direction perpendicular to the outer inviscid streamline, between the pressure and centrifugal forces. However, this force equilibrium is broken within the boundary layer owing to the lower momentum. This results in a weak secondary flow, perpendicular to the local inviscid streamline direction called cross-flow (CF), as shown in figure 2.9.



Figure 2.9.: Sketch of a generic swept wing. The inviscid streamline (ISL) and the corresponding coordinate system $([xyz]_{ISL})$ are shown. The tangential and cross-flow velocity profiles are depicted in a local inviscid-streamline-oriented coordinate system.

2.2.1.1. Primary cross-flow instability

The cross-flow velocity component vanishes both at the wall and the outer edge of boundary layer, but has a finite maxima at some distance from the surface within the boundary layer (see figure 2.9). Thus, the cross-flow component has a boundary layer velocity profile that features an inflection point which gives rise to an inviscid instability named cross-flow instability (Saric, Reed, et al., 2003). As inferred from a number of efforts involving surface flow visualization (see for example Dagenhart et al., 1989; Deyhle & Bippes, 1996; Downs & White, 2013; Serpieri & Kotsonis, 2016; Serpieri, Yadala, et al., 2017), the swept-wing boundary layer has a distinctive pattern, characterised by streaks aligned within a few degrees with the flow direction (see figure 2.10a). These streaks are the footprint of the co-rotating vortices amplified by the primary CF instability modes shown in the tomographic PIV measurement of the swept-wing boundary layer in figure 2.10b. These primary instability modes can be of stationary (steady) or travelling (unsteady) nature, depending on the boundary layer's receptivity to external disturbances.

Receptivity models explain how external disturbances interact and influence with the disturbances within the boundary layer (Reed & Saric, 1989). This process is generally neglected in transition prediction models but a number of experiments have shown that transition in three-dimensional boundary layers demonstrate strong dependence on the external disturbance environment. An overview of the receptivity of three-dimensional boundary layer can be found in Bippes (1999), Saric, Reed, et al. (2003), Tempelmann et al. (2012), Downs & White (2013), Kurz & Kloker (2014), and Rizzo et al. (2019).

In his article, Bippes (1999) synthesizes the results from a number of experiments conducted at *DLR*, *Germany*. Through this it was shown that, primary travelling CF instability modes are generated due to vortical disturbances in the free-stream that get entrained into the boundary layer. Hence, they are observed in medium to high free-stream turbulence conditions. They are excited at higher amplitudes, hindering the development of the stationary modes and thus dominate transition. At low levels of free-stream turbulence ($Tu/U_{\infty} < 0.15\%$), stationary CF modes prevail in the boundary layer and dominate transition. These observations were confirmed by Downs & White (2013) through their extensive experimental efforts. They confirmed that



Figure 2.10.: (a) Fluorescent oil-flow visualization on the pressure side of a 45° swept-wing, reproduced from Serpieri & Kotsonis (2015). (b) Time-average velocity magnitude along the stationary CF axis measured with tomographic PIV, reproduced from Serpieri & Kotsonis (2016).

free-stream turbulence enhanced travelling instability modes while surface roughness, especially close to the leading-edge, played a pivotal role in the amplification of stationary CF instabilities.

The efforts of Prof. Saric and co-workers concentrated mainly on the receptivity of the swept-wing boundary layer to surface roughness. Reibert et al. (1996) carried out experimental investigations in the boundary layer developing on a 45° swept-wing model featuring a laminar airfoil at a chord Reynolds number of $2.4 \cdot 10^6$. They observed the dominance of stationary CF instability in the transition process even when the root-mean-square value of surface roughness was only about 0.25 μ m. This non-uniform distribution of roughness led to the amplification of different instability modes which complicated the analysis of the boundary layer. In order to simplify the flow arrangement and to be capable of comparing the obtained results with the available numerical solution, they employed small cylindrical roughness elements (3.7 mm diameter and 6 μ m height) equally spaced along the span close to the leading-edge, to force specific stationary CF instability modes. The discrete roughness elements (DREs) fixed the wavelength of the amplified CF instability mode, resulting in a more uniform transition pattern. This way, they forced the naturally most-amplified mode predicted through linear stability theory (LST) and a mode with three times the wavelength of the critical mode. In both cases, they observed that stationary CF vortices grew monotonically and reached a given saturation amplitude followed by the onset of transition. Amplification of instabilities with shorter wavelengths than the forced mode was observed but no modes with longer wavelengths were amplified. However, they observed that the naturally most-amplified (or critical) stationary mode dominated the transition process in both scenarios.

The effect of the geometric parameters (width and height) of the roughness elements on the transition process was investigated by Radeztsky et al. (1999). They demonstrated a promotion of transition with roughness element height if its diameter is lower that 0.08 times the wavelength of the forced mode. For larger diameters, roughness element height did not influence the onset of transition. Tempelmann et al. (2012) studied the boundary layer receptivity to surface roughness through parabolised stability equations (PSEs) and adjoint-PSE, together with direct numerical simulation (DNS). The amplitude of CF modes predicted by their receptivity model was lower compared to those observed in the experiments of Reibert et al. (1996). They suggest that this discrepancy is probably caused by experimental uncertainties arising from imperfections in roughness element application and geometry, which results in the additional receptivity observed in the experiments. Given the micrometric size of the installed roughness elements, such imperfections are indeed likely in laboratory conditions. These research efforts, among many others, show that at low free-stream turbulence levels, the ensuing stationary CF instabilities are highly receptive to the surface roughness close to the leading-edge of the swept wing.

2.2.1.2. Secondary cross-flow instability

Although the CF vortices resulting from the primary instabilities are relatively weak, they cause a mean flow distortion and modify the stability of the boundary layer. The upwelling side displaces low-momentum fluid in the lower regions of the boundary layer to the high-momentum regions at the top, and the vice-versa is carried out by the downwelling side of the vortices. These actions cause a momentum modulation in the spanwise direction which results in the development of strong velocity shears within the boundary layer along the spanwise and wall-normal directions. These shear layers are inflectional and unstable to secondary instability mechanisms (White & Saric, 2005). These secondary instabilities grow rapidly along the streamwise direction, triggering transition to turbulence. Laminar breakdown dominated by stationary CF instabilities has a characteristic wedge-shaped or jagged transition front (Reibert et al., 1996; White & Saric, 2005; Serpieri & Kotsonis, 2016). If, on the other hand, primary travelling instabilities are amplified initially, they tend to degrade the stationary structures within the boundary layer and thus lead to a more spanwise-uniform transition front (Downs & White, 2013).

Malik et al. (1999) employed linear and non-linear parabolised stability equations (PSE) to analyse the amplification of primary CF instabilities and a temporal two-dimensional eigenvalue problem to asses secondary instabilities within a swept-wing boundary layer. Through this, they distinguished two types of secondary instability modes. The first, referred to as type-I modes, were produced by the streamwise flow velocity gradient along the span $(\partial U/\partial z)$. These modes were composed of high-energy, low-frequency fluctuations and were observed to exist at the outer side of the upwelling region of the primary CF vortices, as depicted in figure 2.11a (reproduced from Bonfigli & Kloker, 2007). The second type referred to as type-II modes are a result of large gradients of the streamwise velocity along the wall-normal direction $(\partial U/\partial y)$ and thus are located on the top of the primary vortices (see figure 2.11b). Compared to type-I secondary modes, these modes contain high-frequency fluctuations but with lower energies. A third type of secondary instabilities, or type-III modes, were also identified later (Högberg & Henningson, 1998; Janke & Balakumar, 2000). These modes are a product of the spanwise modulation of the primary travelling instabilities caused by stationary CF vortices. Thus, their role in the transition process necessitates the amplification of both, primary stationary and primary travelling instabilities, and is dictated by free-stream turbulence levels (Downs & White, 2013).



Figure 2.11.: Location of secondary instability modes along the stationary CF vortices (dotted black curves), reproduced from Bonfigli & Kloker (2007). (a) *type-I* modes, (b) *type-II* modes and (c) *type-III* modes.

In their experiment, White & Saric (2005) aimed to provide detailed measurements of the secondary instability. The swept-wing model employed and the flow conditions were the same as Reibert et al. (1996). The boundary layer was forced with the critical stationary instability mode using DREs fixed along the span. The *type-I* and *type-II* modes were clearly identified and their sudden growth caused laminar breakdown. Additionally, the location of these secondary modes on the primary CF vortices was in congruence with the finding of Malik et al. (1999). An important outcome of this study was, it revealed that the occurrence of *type-I* or *type-II* secondary instability mechanisms depended on the wavelength of primary CF instabilities present in the boundary layer with respect to that of the critical CF instability. The *type-I* mode is more common and is observed when the primary CF vortices have a wavelength similar to the critical primary instability or lower (sub-critical), owing to the high spanwise shear generated. When super-critical stationary modes are forced (i.e. modes with wavelength greater than the critical mode), it results in the reduction of the spanwise shear and hence the amplification of the *type-I* mode. Thus, the *type-II* modes trigger transition to turbulence.

The DNS study of Wassermann & Kloker (2002) demonstrated the nature and topology of type-I secondary modes developing around primary CF vortices. One important finding of this simulation is that secondary instability modes are triggered only with the presence of an unsteady background disturbance, which in a physical sense is the turbulence in the free-stream. Thus, in practical applications, secondary instabilities are always triggered. However, in their simulation, the authors demonstrate that when the unsteady disturbance is switched-off, the structures associated with secondary modes just advect downstream. This confirms that secondary instabilities are inherently convective (as also confirmed by Janke & Balakumar, 2000; Serpieri & Kotsonis, 2016). Bonfigli & Kloker (2007) investigated secondary instability mechanism through DNS and secondary linear stability theory (SLST). Both techniques captured the structures related to type-II modes, SLST did predict the growth of structures corresponding to these modes. Additionally, they demonstrated that in the absence of stationary CF instability modes, the type-III modes did not destabilize, which is expected as discussed earlier. A final important

outcome of this work is that the characteristic vortical flow structures associated with secondary instabilities (especially *type-I* and *type-II* modes) bear close resemblance to those generated by the Kelvin-Helmholtz (KH) instability mechanism.

Serpieri & Kotsonis (2016) experimentally investigated the development of CF instability modes and flow transition on the pressure side of a 45° swept-wing model. The experiments were carried out at a wing incidence angle of 3° and chord Reynolds number of $2.17 \cdot 10^{6}$ (it is worth noting that the current study is carried out on the same swept-wing model and at similar flow conditions, see section 5.1). The authors estimated the stability of the boundary layer through linear stability theory (LST) and investigated the flow topology qualitatively using fluorescent oil. Through this, the spanwise velocity modulation due to forcing the critical stationary mode by DREs installed at the leading edge and the consequent jagged transition front was visualized. They observed that the stationary mode grows along the chord and reaches a certain saturation amplitude beyond which, it starts to decay. This decay is accompanied by the amplification of unsteady fluctuations related to secondary instability mechanisms. Tomographic PIV was then employed to study the evolution of the primary CF instabilities close to chordwise station where they saturate. This also led to the three-dimensional description of the spatial organisation (through POD) and evolution of unsteady modes within the boundary layer. The time-averaged vectors fields measured here showed good agreement with the hot-wire measurements.

Serpieri & Kotsonis (2016) found that at the investigated flow conditions, fluctuations corresponding to the type-I secondary instability was centred in the frequency band 5 kHz $\leq f \leq$ 6 kHz. These fluctuations were organised as narrowly-spaced high and low-speed fingers, just outside the upwelling region of the stationary CF vortices. As these structures almost span the entire boundary layer thickness, they are advected at different velocities along the wallnormal direction and hence seem to have a helicoidal shape as reported by in the DNS study by Wassermann & Kloker (2002). The type-II secondary modes were observed in the band 7 kHz $\leq f \leq 8$ kHz but had lower amplitudes compared to the type-I modes. These occurred in the top region of the stationary CF vortices as observed through hot-wire measurements. Finally, travelling CF instability modes and hence, the low-frequency type-III secondary modes was contained in the frequency band 350 Hz $\leq f \leq 550$ Hz which corroborated with the predictions of LST. These fluctuations were organised as elongated high and low-speed pairs on the inner side of the upwelling region of the stationary CF vortices. Serpieri & Kotsonis (2018) installed a DBD actuator applying two-dimensional forcing close to the location of stationary mode saturation. The unsteady forcing of the actuator was exploited to condition the boundary layer with type-I and type-III secondary modes. Boundary layer velocity measurements were obtained through hot-wire anemometry and the acquired results confirmed the observations made with the tomographic-PIV fields.

2.2.2. Control strategies for cross-flow instability

Skin friction in a turbulent boundary layer is significantly higher than in its laminar counterpart at the same Reynolds number (Joslin, 1998). Ensuring that the boundary layer over aircraft wings remain laminar has a number of advantages that predominantly stem from the reduction of friction drag. Thus, delaying transition to turbulence can aide in significantly improving the performance efficiency of an aircraft. This can ultimately results in cutting-back of flight costs due to lower fuel consumption and will lead to lower emissions and curtail the detrimental effect of aircraft on the environment. Laminar flow control (LFC) is a broad term that encompasses the deployment of both, passive and active flow control techniques, to retain laminar flow over aerodynamic bodies by delaying laminar-to-turbulent transition. The main driver of laminar-to-turbulent transition in a threedimensional swept-wing boundary layer is the cross-flow (CF) instability mechanism. This instability is the only flow component that is holding back total laminar flow control in sweptwing boundary layers. A number of techniques have been proposed and investigated to control CF instability and the corresponding laminar breakdown. Some of these strategies, that are relevant to the current investigation and summarized here. It is worth noting here that research efforts have mainly focussed on controlling stationary CF instability modes in the boundary layer. This is because, conditions during cruise-flight present a low free-stream turbulence scenario and as such, stationary modes dominate the boundary layer.

Wall cooling and suction have been predominantly used for boundary-layer flow control. While both these techniques are very effective in controlling TS instabilities (Liepmann, Brown, et al., 1982), thermal flow control has minimal effect on CF instabilities (Mack, 1980). As this mechanism arises from an inflection point within the boundary layer, thermal perturbations have no influence on it. Surface suction is based on reducing the velocity magnitude within the boundary layer. For many years, it was considered to be the only technique to control CF instabilities and the consequent transition to turbulence in three-dimensional boundary layers. However, the necessity for moving parts to realize suction and the amplification of CF instabilities due to imperfections in the suction orifices are major drawbacks of this flow control method. The main contributions towards control of CF instability and the different strategies proposed to realize this control are summarised in the following sections.

2.2.2.1. Upstream flow deformation (DRE/UFD)

As discussed earlier in section 2.2.1.1, Prof. Saric and co-workers, especially through the experimental investigation of Reibert et al. (1996), observed the high sensitivity of stationary CF instabilities to leading-edge surface roughness. Following this observation, Saric, Carrillo, et al. (1998) proposed inducing sub-critical stationary modes, i.e. modes with spanwise wavelength lower than the critical mode wavelength, to delay laminar-to-turbulent transition in three-dimensional boundary layers. Named sub-critical discrete roughness elements (DRE) strategy, it is based on using an array of equi-spaced roughness elements (as done by Reibert et al., 1996) to amplify sub-critical stationary modes. This leads to a mean-flow deformation within the boundary layer, which delays the onset and destabilisation of the critical mode. Consequently, this postpones the amplification of secondary instabilities and thus, delays the eventual flow transition to turbulence.

Investigations in the theoretical and numerical domains that followed (Malik et al., 1999; Wassermann & Kloker, 2002; Hosseini et al., 2013) confirmed the positive effect of this strategy in delaying transition. An instantaneous velocity field from the DNS by Hosseini et al. (2013) investigating the mechanism of this strategy is reproduced in figure 2.12. In the natural case shown (figure 2.12a), the transition process begins close to x/c = 0.5. When a sub-critical stationary mode is forced (figure 2.12b) on the other hand, the mean-flow deformation caused delays the onset of the critical mode and the transition is postponed to x/c = 0.7. This strategy was later generalised as upstream flow deformation (UFD) by Wassermann & Kloker (2002), to encompass the use of any device or method that can induce the desired sub-critical stationary modes.



Figure 2.12.: Isosurfaces of the instantaneous velocity field of the natural case (a) and with a sub-critical CFI mode being forced (b). Reproduced from Hosseini et al. (2013).

Several investigations following this strategy involving wind-tunnel experiments, numerical simulations and free-flight experiments have been conducted ever since this method was first proposed. White & Saric (2000) employed pneumatically actuated roughness elements in their wind-tunnel experiment. The motive behind this study was to realize a flow control system based on DRE/UFD that can potentially be used for a wide range of flow conditions. This system allowed the variation of both spanwise wavelength of forcing and its amplitude (DRE height). While this system did not provide positive results due to technological limitations, it enhanced the interest towards active LFC based on the DRE/UFD strategy.

Ever since, a number of methods and devices have been employed to force monochromatic stationary CF instability modes to realize the UFD control strategy. While the motive was not transition control, Deyhle & Bippes (1996) used localized surface heating to condition the boundary layer. Lohse et al. (2016) experimentally investigated the idea of conditioning the primary CF instability with steady, localised suction and blowing. To this effect, orifices were made in a swept flat-plate model and the amplitude of the forced instability modes was controlled with the mass-flow rate through each orifice. Both suction and blowing showed effective conditioning of the boundary layer with the forced CF instability mode, with suction being a more robust technique. Schuele et al. (2013) demonstrated the application of the UFD strategy on a yawed circular cone in supersonic flow conditions. They employed leading-edge surface indentations placed at the critical and sub-critical (azimuthal) wavelengths to realize this strategy. When compared to the smooth cone tip, critical mode forcing promoted transition while sub-critical mode forcing delayed transition. DBD plasma actuators that apply three-dimensional EFD forcing were also employed towards the same goal. This, and other investigations employing DBD plasma actuators are discussed separately in section 2.2.3.

The global working mechanism of the DRE/UFD control strategy has been sufficiently assessed, through extensive experimental and numerical investigations. However, eventual transition delay, especially with respect the unforced flow, has not always been demonstrated experimentally (see for example Downs & White, 2013). Later flight test too have been inconclusive (Saric, West, et al., 2015). This suggests the high sensitivity of the DRE/UFD control strategy to a number of factors. The effectiveness of this strategy in delaying transition seems to depend, among others, on boundary layer receptivity to background surface roughness (Lovig et al., 2014) and free-stream turbulence (Downs & White, 2013), and boundary layer stability characteristics due to Reynolds number and pressure distribution. The stability parameters are in fact very crucial to the DRE/UFD strategy as it is only effective if firstly, a sub-critical mode is forced and secondly the amplitude of the forced mode is considerably higher close to



Figure 2.13.: λ_2 visualisation of the primary CF vortices without (a) and with active flow control performing base-flow modification (b). Reproduced from Dörr & Kloker (2015b).

the leading edge compared to the critical mode. Thus, it is necessary to further investigate each of these factors in order to successfully and robustly implement the DRE/UFD strategy for the realization of effective and efficient LFC.

2.2.2.2. Base-flow modification (BFM)

An alternate control strategy to achieve LFC in three-dimensional boundary layers named baseflow modification (BFM), was proposed by Prof. Kloker and co-workers through their extensive efforts using DNS. This strategy is based on the suppression of the cause of the instability, which in this case is the CF velocity component within the boundary layer. As shown in figure 2.13, compared to the reference (no forcing) case, performing base-flow modification through the application of a force results in significant attenuation of the CF vortex strength and hence the mean-flow deformation within the boundary layer.

Messing & Kloker (2010) employed (spanwise) homogenous wall suction to realize this strategy. This lead to fuller streamwise velocity profiles close to the wall and weaker CF velocities. This results in enhanced boundary-layer stability, ultimately leading to transition delay. However, suction orifices can cause amplification of unsteady modes due to imperfections in orifice size, position or strength of suction (i.e. the mass-flow rate). The authors also investigated a combination of distributed suction and UFD (referred to as 'formative suction'). It is reported that the combined effect is a reduction of the required suction-flow rate which results in a more efficient control system.

Bonfigli & Kloker (2007) demonstrated that small induced vortices that were orthogonal to the local shear were very effective in improving boundary layer stability. This observation resulted in other control methods such as 'pinpoint suction' proposed by Friederich & Kloker (2012). This is based on directly weakening the stationary CF vortices and showed beneficial effects in delaying transition. A number of efforts involving EFD body force generated by a DBD plasma actuator was also employed by this group. These are discussed in section 2.2.3.

An essential difference of the BFM strategy from the DRE/UFD approach is that BFM depends on the suppression of the mean boundary layer. Therefore, its performance is relatively independent of the nature of CF instabilities, including the modes' wavelength and if they are stationary or travelling type. This renders the BFM strategy to be more robust. Although, a disadvantage of such mean-flow manipulating strategies is the higher power input necessary to realize the necessary control. To this point though, the concept of base-flow modification for suppression of CF instability has not been demonstrated experimentally.

2.2.3. Cross-flow instability control with DBD plasma actuators

Use of the EFD body force generated by a DBD plasma actuator for controlling the CF instability in three-dimensional boundary layer has garnered considerable attention in the last decade. One of the first reported efforts of this pertains to the work of Schuele et al. (2013) on the boundary layer of a yawed circular cone at supersonic flow conditions. Apart from using surface indentation near the cone's tip to realize the UFD control strategy as discussed in section 2.2.2.1, they also employed DBD plasma actuators to do the same. The actuators were designed and installed such that, the applied three-dimensional EFD forcing would generate boundary layer perturbations similar to the surface indentations. The flow arrangement with critical mode forcing was presented, which showed great similarity to the effect of surface indentations. This shows that DBD actuator forcing can cause the mean-flow deformation according to the forced stationary mode. Unfortunately, the flow arrangement with sub-critical mode plasma forcing and also the eventual transition front displacement due to the applied plasma forcing were not reported.

Several DNS investigations employing the EFD body force in these boundary layers have been reported by the group of Prof. Kloker in recent years (Dörr & Kloker, 2015b, 2016, 2017; Dörr, Kloker, & Hanifi, 2017; Guo & Kloker, 2019, 2020). In most of these studies the the effect of the plasma actuators are modelled as a local steady wall-parallel body force. The magnitude of this body force distribution is retrieved from the induced velocity fields according to the empirical model proposed by Maden et al. (2013). The volumetric body force model used in these DNS investigations and its limitations are presented by Dörr & Kloker (2015a).

Dörr & Kloker (2015b) and Chernyshev et al. (2016) numerically investigated the application of EFD body force to perform base-flow modification. They observed that application of plasma body force against the CF velocity component leads to its reduction and hence, the stabilisation of the boundary layer and a delay in transition. This is similar to the case of homogenous suction demonstrated by Messing & Kloker (2010). Dörr & Kloker (2015b) also studied the effect of counter-blowing. In this case, the boundary layer became thicker with related near-wall momentum loss in the streamwise direction and increased the CF velocity component. Thus, it destabilized the boundary layer and caused promotion of transition.

Dörr & Kloker (2016) then undertook the task of employing spanwise-modulated plasma forcing in this boundary layer, akin to 'pinpoint suction' proposed by Friederich & Kloker (2012). The effect of placement of spanwise-modulated plasma forcing with respect to the oncoming CF vortices was extensively investigated. When plasma forcing was applied against the incoming flow, a transition delay was always achieved, whatsoever be the forced stationary CF instability mode (critical or sub-critical). This can be observed in figure 2.14 reproduced from Dörr & Kloker (2016), where application of one plasma body force distribution per critical CF instability wavelength is seen to delay the onset of secondary instabilities and thus transition. The authors



Figure 2.14.: Cross-flow vortex visualisation reproduced from Dörr & Kloker (2016). The unforced case is shown in (a). Effect of applying one plasma body force distribution (filled black bars) per critical CF instability wavelength is depicted in (b).

conclude this to be the case due to two mechanisms acting on the boundary layer collaboratively. The first is the direct attenuation of the oncoming CF vortices, when the maximum body force counteracts the upwelling region of the CF vortices (most ideal position). This weakening of the CF vortices reduces the velocity shears and damp the secondary instability modes, thus delaying transition. In all other cases, the plasma forcing reduces the CF component in the boundary layer, essentially performing a localised BFM action. Therefore, with certain designs of DBD actuators configured for UFD strategy (mainly the finger-tip actuators used by Schuele et al., 2013; Serpieri, Yadala, et al., 2017), the applied three-dimensional plasma forcing, while conditioning the boundary layer with specific CF instability modes, can be expected to perform BFM action locally. However, the effect of this action and how the two strategies collaborate in such a situation has not been investigated experimentally.

Dörr & Kloker (2017) employed the same steady body force model to investigate the applicability of DBD plasma actuators for UFD control. Three volumetric force distributions per fundamental wavelength was applied on the numerical boundary layer, resulting in forcing a highly sub-critical stationary mode. Introduction of the EFD body force against the CF component was found to be most effective, with successful onset of narrowly-spaced CF vortices corresponding to the forced mode. While these vortices were observed to vanish downstream, no amplification of secondary instabilities were observed in the simulated domain (which was observed in the no forcing case), suggesting a delay in transition. Furthermore, it is shown that at higher forcing amplitudes, transition delay is achieved, but due to the stabilization of the boundary layer through BFM action rather than UFD control.

Experimental investigation of employing DBD plasma actuators to realize the UFD control strategy in a swept-wing boundary layer was performed by Serpieri, Yadala, et al. (2017). The swept-wing model and flow conditions were similar to that in the experimental efforts of Serpieri & Kotsonis (2016) (and also the current thesis). A novel DBD plasma actuator that generated a three-dimensional spanwise-modulated forcing was designed and installed close to the swept-wing leading edge to force the necessary stationary CF instability modes. Sub-critical, critical and super-critical stationary modes were forced in this study with the respective actuators. Successful re-organisation of the boundary layer with respect to the forced stationary mode was observed in both, the flow topology captured through infrared thermography and PIV. IR thermography fields depicting this is reproduced in figure 2.15. The unforced boundary layer (figure 2.15a) shows strong homogeneity in both the spanwise and chordwise directions and no stationary modes appear in the captured FOV. The bright spots close to the swept-wing leading edge in figure 2.15b are the thermal signatures of the plasma discharge of the novel actuators.


Figure 2.15.: IR thermography fields showing the thermal signature of the unforced boundary layer (a) and the stationary CF vortices successfully forced by the employed actuator (b). Bright spots close to the swept-wing leading edge in (b) are the thermal signature of the plasma discharge.

A streaky pattern is observed in this field further downstream which are the forced stationary CF vortices.

Although forcing was applied at a frequency of 2 kHz which was about five times the frequency band of the travelling CF instability modes in the investigated flow conditions (see section 2.2.1.2), excitation of unsteady travelling instabilities was observed through the fluctuating fields captured using PIV. While it is not reported, the authors conclude that transition to turbulence might be promoted due to the actuation. Hence they suggest that forcing with plasma actuators need to be applied at much higher frequencies to avoid such an effect. Furthermore, the authors demonstrated the effect of the time-invariant flow induced by DBD plasma actuators on the swept-wing boundary layer. The time-averaged EFD body force applied by the actuators was observed to affect the receptivity or stability characteristics of the boundary layer, thus influencing the amplification of instability modes. To study these effects, a simplified numerical model based on LST was proposed which is presented in more detail in section 3.3.1.2.

In all the DNS studies discussed until now, the plasma body force is assumed to be a steady one. However, this is not the case in reality as DBD actuators produce an unsteady walljet at the frequency of the supplied AC signal (see Benard & Moreau, 2014). This inherent unsteadiness of the applied forcing causes the amplification of travelling instabilities as observed by Serpieri, Yadala, et al. (2017). Dörr & Kloker (2015b, 2016) studied the effects of this force unsteadiness and confirmed the amplification of travelling CF instabilities corresponding to the specific forcing frequency. They work under the assumption that the forcing frequency was well beyond that of secondary instabilities and hence has no detrimental effect. However, they acknowledge that, at higher Reynolds number compared to their simulations the frequency band of secondary instabilities will be greater as well, requiring application of forcing at much higher frequencies. Finally, in the investigation of Serpieri & Kotsonis (2018) mentioned in section 2.2.1, the authors employed DBD actuator forcing successfully to condition the boundary layer with the desired secondary modes. This once again shows that the EFD forcing being applied has an influence on unsteady modes in the current boundary layer. In fact, in the case with the most downstream installation of the actuator, Serpieri & Kotsonis (2018) observed the highfrequency forcing triggering laminar-to-turbulent breakdown. Thus, it can be concluded that to achieve transition delay with DBD plasma actuators, the forcing frequency must be much higher compared to secondary instability fluctuations in the boundary layer. However, forcing at such high frequencies (especially for higher Reynolds number flows) could be technologically unfeasible and thus, introduction of undesirable and detrimental flow fluctuations is expectable. In their experimental investigation, Kim et al. (2016) investigated the applicability of ringtype DBD plasma actuators for CF instability control in a swept-wing boundary layer. Unlike classical DBD actuators discussed until now, ring-type actuators induce a wall-normal jet and act as virtual roughness elements. Array of these actuators with different inter-element spacing were installed on the model different chordwise stations. Here as well, the applied forcing successfully re-organised the boundary layer according to the forced stationary instability. Location of transition was in fact reported to be promoted due to the applied forcing. Here too, the authors conclude that the amplification of travelling CF instabilities and their corresponding interaction with the stationary vortices leads to transition promotion. Shahriari et al. (2018) investigated these experimental conditions using DNS and observed that application of higher magnitude force attenuates unsteady fluctuations and successfully delay transition. However, it is worth noting that knowledge regarding the working mechanism of ring-type actuators is rather limited. Thus, the employed method of modelling the body force distribution might not have encompassed all the flow structures that these actuators induce. Furthermore, this DNS study too, is based on a steady EFD body force assumption.

2.3. PLANE MIXING LAYER

The interface of two fluid streams parallel to one another and having different velocities or densities is commonly observed in many natural phenomenon and practical applications. Intense mixing of the two streams occurs in the region with a gradient of the specific parameter and hence, such fluid systems are generally referred to as mixing layers. The interaction of the two streams results in strong shearing action and hence, mixing layer is fundamentally a representation of a wider class of inviscidly unstable free-shear flows. This flow configuration dominates a number of phenomenon including flow in jets and wakes, pollutant dispersal in atmosphere, meteorology in connection with the formation of cyclones, among others.

Mixing layers arising due to the interaction of two parallel streams with different free-stream velocity (U_1 and U_2 , U_1 being the lower of the two velocities, velocity ratio being $r = U_1/U_2$) at the trailing-edge of a splitter plate was comprehensively investigated first by Liepmann & Laufer (1947). They generated a mixing layer by discharging a single stream into a quiescent surrounding fluid ($U_1 = 0 \text{ m s}^{-1}$, r = 0) and observed a mixing layer that spreads preferentially into the quiescent fluid. Once the flow was fully developed, the velocity profiles at all measured streamwise locations were well approximated by Görtler's error function profile (Görtler, 1942). More importantly, the results showed that this flow configuration is self-preserving or attains self-similarity (also confirmed by Wygnanski & Fiedler, 1970).

Self-preservation implies that the flow field is in a state of dynamic equilibrium and is not affected by the initial conditions any more (Townsend, 1976). Therefore, the mean and fluctuating velocity parameters of the flow field are independent of the Reynolds number and become geometrically similar along the downstream direction when non-dimensionalised with a single velocity and plotted along a non-dimensional (similarity) coordinate (Townsend, 1976; Bell & Mehta, 1990; Mehta, 1991). However, the location downstream of the splitter-plate trailing edge where self-similarity is realized varies depending on the initial state of the mixing layer, i.e. if the incoming boundary layers are laminar or turbulent. Mixing layers originating from laminar initial conditions are characterized by ordered motion of large-scale flow structures and attain self-similarity much later compared to those ensuing from turbulent initial conditions. Truly turbulent mixing layers are characterized by small-scale motions which enable efficient



Figure 2.16.: (a) Vorticity distribution in a free-shear layer theoretically computed by Michalke (1965b). (b) Smoke visualisation showing vortex formation in a jet flow (flow from left), reproduced from Freymuth (1966).

mixing at molecular levels. In fact the mechanisms that lead to laminar-to-turbulent transition in mixing layers have been widely researched. While the state of the incoming boundary layer is a prime factor in the subsequent development of a self-similar mixing layer, a number of other parameters (such as trailing-edge thickness, free-stream turbulence level, etc.) can also have an effect (Dziomba & Fiedler, 1985).

2.3.1. Instability mechanisms in the mixing layer flow configuration

2.3.1.1. Primary instability

Free-shear flows, including mixing layers and jets, present an inflectional velocity profile. It is well known through the theorem first formulated by Lord Rayleigh (1880) that, an inviscid fluid with an inflection point in its velocity profile is unstable to certain small-amplitude wavy disturbances. Free-shear flow fields are unstable to perturbations arising from the inviscid *Kelvin-Helmholtz* (KH) instability mechanism that is susceptible to three-dimensional secondary instabilities which can facilitate rapid onset of turbulence. Indeed, the KH instability mechanism is a fundamental phenomenon that is abundantly observed in nature.

Until the 1960s however, free-shear flows were considered to have two regions: the first being the turbulent flow featuring random three-dimensional motions and vorticity fluctuations, and the second being the irrotational fluid surrounding it, separated by a sharp interface. As shown by Corrsin & Kistler (1955), viscous forces within this interface ("laminar superlayer" as they call it) play a major role in propagating vorticity from the turbulent flow to the outer potential flow and resulting in the expansion of the turbulent region. The motion within the turbulent region itself was considered to be homogeneous and isotropic, as postulated by Taylor (1935). However, the theoretical efforts of Michalke (1965a,b) indicated towards the existence of discrete vortices within free-shear layers. In fact, Michalke (1965b) used non-linear Helmholtz vorticity equation on a hyperbolic-tangent velocity profile to understand vortex formation in free-shear flows. With time, the free-shear layer was observed to become thinner in certain regions and subsequently thicker in the others, with a vorticity extrema located in the thicker region (figure 2.16a). This corroborated very well with smoke visualisation experiments of the time such as that of Freymuth (1966) in jet flow shown in figure 2.16b.

Michalke (1964, 1965a) numerically determined the stability characteristics of a shear layer by solving the linearised Rayleigh equation on the hyperbolic-tangent velocity profile. In the first work, it was assumed the instability waves were growing with time while in the second work,



Figure 2.17.: Shawdowgraph of mixing layer forming between a helium (upper) and nitrogen (lower) streams at velocities $U_2 = 5 \text{ m s}^{-1}$ and $U_1 = 1.9 \text{ m s}^{-1}$ respectively. The organised motion of vortical flow structures resulting from the Kelvin-Helmholtz roll-up is clearly visible. Reproduced from Brown & Roshko (1974).

stability of spatially growing instabilities were investigated. By comparing the obtained results with experimental data acquired by Freymuth (1966), it was demonstrated that spatially growing disturbances describe the instabilities in a shear layer more precisely than temporally growing ones. This was also confirmed by Huerre & Monkewitz (1985), who classified the instabilities as absolutely unstable (disturbances grow everywhere with respect to time) and convectively unstable (disturbances are convected away from the origin as they amplify, eventually leaving the basic flow undisturbed). Through linear stability analysis, they indicated that shear layers with a velocity ratio $R = (U_2 - U_1)/(U_2 + U_1) = (1 - r)/(1 + r) < 1.315$ are convectively unstable, and the instabilities within develop spatially and not temporally.

The shadowgraphy images of the plane mixing layer at high Reynolds number $(O(10^5))$ captured by Brown & Roshko (1971, 1974) showed (see figure 2.17) that this flow configuration is dominated by large-scale, predominantly two-dimensional, spanwise vortex-like coherent structures (or eddies) that are well organised. This was an important discovery as it revealed the occurrence of ordered motion within the mixing layer, which contradicts classical turbulence theory (Taylor, 1935). These structures were found to be similar to those observed in the late stages of instability growth in a laminar shear layer, hence suggesting that the driving phenomenon which governs the growth of the mixing layer must be hydrodynamics instabilities arising from the *Kelvin-Helmholtz* (KH) instability mechanism. Thus, the plane mixing layer is generally considered to be the simplest representation of a free-shear layer that attains self-similarity (Wygnanski & Fiedler, 1970; Brown & Roshko, 1974). Prior to these findings, experimental investigations on free-shear layers were based on measurements at the edge of jet flows where shear layers are known to occur (for example, Freymuth, 1966). The mixing layer flow field being a representative of free-shear flows provided the research community a simpler flow configuration to investigate shear-layer instabilities. Thus, the findings from investigating this flow configuration can be applied to the shear layers at the edges of jets, and vice versa.

The interaction of these large-scale vortices and its effect on mixing layer growth and transition was studied by Winant & Browand (1974), using a mixing layer formed by two streams of water at moderately low Reynolds numbers (45 to 850). Dynamics of the large-scale vortical structures was studied by introducing a dye into the flow field. They observed that initially, a laminar region was evident where the mixing layer grew as the square root of downstream distance. In



Figure 2.18.: Sequence of photographs showing vortex pairing in a mixing layer where upper and lower streams move at velocities $U_1 = 1.44$ cm s⁻¹ and $U_2 = 4.06$ cm s⁻¹ respectively. The camera moves at the average velocity of the two streams $\overline{U} = (U_1 + U_2)/2 = 2.75$ cm s⁻¹. Downstream distance to the centre of each frame is indicated to the right. Reproduced from Winant & Browand (1974).

this region, small amplitude waves form due to the linear instability (left side of first four frames in figure 2.18). These waves distort the flow field, concentrating the vorticity, thus resulting in the thinning of the layer in certain regions and thickening it in others. This overall picture corroborates very well with the stability calculations of Michalke (1965b) discussed earlier (figure 2.16a).

Winant & Browand (1974) found that these vortical lumps grow further downstream and form a periodic train of two-dimensional vortices. Due to small variations in either the strength of these vortices or the spacing between them, neighbouring vortices roll around each other and through viscous diffusion, form a single larger vortex as depicted in their photographs reproduced in figure 2.18. It was found that this *vortex pairing* or *amalgamation* was responsible for the growth of the mixing layer. Furthermore, the results suggested that entrainment of irrotational fluid does not occur only due to diffusion of vorticity from the mixing layer to the outer potential flow. But a higher amount of entrainment occurs during vortex-pairing, where irrotational fluid between pairing vortices get engulfed into the resulting vortex. In an unforced mixing layer, these vortex interaction occur stochastically in space and time resulting in its linear growth.

Ho & Huang (1982) investigated the influence of sub-harmonic instabilities on the evolution of the mixing layer, in particular on the process of vortex pairing. These experiments were carried out in a water tunnel. The authors observed that coherent structures in the forced mixing layer developed periodically due to streamwise and transverse velocity perturbations. The vortices form at the response frequency (f_r) which is the closest harmonic of the forcing frequency (f_f) to the frequency of the fundamental KH instability (f_0) in the mixing layer. The interaction between these three frequencies lead to vortex pairing. As f_f is a sub-harmonic of f_r , $M = f_r/f_f$ adjacent vortices are grouped together due to lateral displacement (caused due to phase difference with fundamental) and the convection velocities of these vortices are slightly different. This eventually leads to merging of these vortices. The sub-harmonic acts as the catalyst for vortex pairing and is observed to be essential for this process. The subharmonic grows linearly and saturates close to the location of vortex merging post which, it decays. Furthermore, Ho & Huang (1982) propose that evolution of the coherent structures are governed by local stability and global feedback mechanisms. The local stability provides a band of possible unstable frequencies while the global feedback fine tunes the mixing layer to a specific frequency through the selection of the response frequency.

In their DNS, Mansour et al. (1988) studied the effect of the phase difference between the fundamental frequency in the mixing layer and its sub-harmonic. They demonstrate that changing the phase angle between the fundamental and its sub-harmonic has an effect on the evolution of the mixing layer and the vortices in it. Both the fundamental and sub-harmonic modes are observed to grow, saturate and then decay. At 0° phase difference, as the fundamental saturates, the growth of the sub-harmonic is enhanced due to resonance, leading to earlier amalgamation. At a 90° phase difference, the sub-harmonic growth is diminished as the fundamental mode saturates. The sub-harmonic growth is recovered post the saturation of the fundamental and the amalgamation process begins, but later compared to the previous case. Furthermore, the effect of feedback mechanism is also acknowledged.

2.3.1.2. Sub-harmonic resonance and sidebands

Having presented the importance of sub-harmonic instabilities in the vortex-pairing process, it is now imperative to address the *sub-harmonic resonance* phenomenon. First proposed by Kelly (1967) using a weakly non-linear temporal theory of a parallel flow, sub-harmonic resonance is a result of the non-linear interaction between the fundamental instability of the mixing layer and its first sub-harmonic. Depending on the relative phase between these two modes, close to the streamwise station where the fundamental instability becomes neutrally stable, the sub-harmonic instability can be reinforced resulting in its higher growth rate. Monkewitz (1988) later extended this to spatially growing waves in a mixing layer with a hyperbolic-tangent velocity profile with the assumptions that the flow is weakly non-linear and the fundamental instability is neutrally stable. However, it is important to note that physical flows are not weakly non-linear and the fundamental is not neutrally stable. Thus, these results are only valid in a small streamwise region around the location where the fundamental instability attains its maximum/saturation amplitude. Besides these assumptions, the existence of a critical fundamental mode amplitude $(\approx 0.0153 \cdot R^2)$, scaled with \overline{U}) required for it phase-locking with the sub-harmonic and modifying its growth rate was demonstrated by Monkewitz (1988). This was later confirmed by Thomas & Chu (1989) who measured the critical fundamental mode amplitude to be 0.015 in their jet-flow experiment (R = 1).

Monkewitz (1988) also studied the effect of sub-harmonic sidebands, which are a pair of modes with frequencies slightly different from the sub-harmonic. It was suggested that when the subharmonic modes enter the region where the fundamental mode amplitude is greater than the critical value, the resonant interaction amplifies a narrow frequency band centred around the sub-harmonic which combine to form an amplitude modulated sub-harmonic instability wave. These sidebands are confined within the frequency band $St_0/2\pm 0.08 \cdot St_0$ (St_0 being the Strouhal number of the fundamental instability). This is also corroborated by the experimental studies of Laufer & Zhang (1983) and Thomas & Chu (1989).

Husain & Hussain (1995) investigated the sub-harmonic resonance mechanism through the



Figure 2.19.: Power spectrum of streamwise velocity component under detuned excitation to study subharmonic sideband frequencies, reproduced from Husain & Hussain (1995). The f in the figure (not axis) is the frequency of the fundamental instability and Δf is the detuning applied to the sub-harmonic (2.5% of f).

dynamics of vortices within the shear layer of a jet flow configuration. The shear layer was excited acoustically, with the driving signal input to the speaker consisting of two harmonically related frequencies. This way, they were able to investigate the effect of the phase-difference between the fundamental and sub-harmonic instability waves. The results show that vortex pairing is stabilized at a broad spectrum of phase differences, but is attenuated at a critical value when the initial amplitude ratio between the forced sub-harmonic and fundamental instability modes is low. At higher initial amplitude ratios, forcing at the critical phase difference results in vortex shredding.

By introducing a slight detuning to the sub-harmonic, Husain & Hussain (1995) investigated the effect of sideband frequency forcing. Perturbations at the sub-harmonic were still observed due to the feedback from pairing as can be seen in figure 2.19, although it now had an amplitude and phase modulation (as proposed by Monkewitz, 1988). As such, pairing was observed to now be dependant on the phase-difference between the forced fundamental instability and the modulated sub-harmonic. Furthermore, the authors measured the critical fundamental instability amplitude to be ≈ 0.03 which is twice compared to prediction of Monkewitz (1988). This difference is mainly attributed to the parallel-flow assumption of the theoretical computations. However, through their results, Husain & Hussain (1995) suggest that the amplitude of the fundamental instability needs to be greater than the theoretical critical value in order to counteract the influence of feedback from pairing vortices.

2.3.1.3. Secondary instability

Apart from the large-scale spanwise vortices, Brown & Roshko (1974) also observed distinctive streaks in the plane view of the mixing layer, similar to that shown in figure 2.20, which did not form any distinguishable pattern on the edge views (also see in figure 2.17). These streamwise streaks are caused by secondary instabilities prevalent in the mixing layer flow field. More shadowgraph images of these streamwise streaks were presented by Konrad (1976, PhD thesis). The spanwise vortices create a positive strain in the braid region (between two spanwise vortices) which stretches the residual vorticity in this region. This stretching mechanism concentrates the



Figure 2.20.: Edge-view of mixing layer showing the large-scale vortical flow structures and plane view showing the streamwise vortices arising due to the secondary instability mechanism, reproduced from Konrad (1976, PhD thesis)

vorticity, resulting in the formation of counter-rotating pairs of vortex loops oriented along the streamwise direction (Lasheras & H. Choi, 1988; Nygaard & Glezer, 1991; Moser & Rogers, 1991; Rogers & Moser, 1992).

The interaction of these two vortical flow structures within the mixing layer plays an important role to cause mixing transition, which was first reported by Konrad (1976, PhD thesis). Near the splitter-plate trailing-edge, the mixing layer contains highly organised spanwise structures. However, farther downstream, small-scale turbulence is superimposed upon the large-scale flow structures by the streamwise-oriented secondary vortices. This change in the state of the mixing layer is referred to as transition or *small-scale mixing transition*, as it has a strong effect on mixing at the molecular level.

In their experiments, Wygnanski, Oster, et al. (1979) acquired correlated measurement of velocity and temperature within the mixing layer. The authors hypothesized the occurrence of cylindrical eddies aligned with the streamwise direction. These observations supported the notion that, at all initial flow conditions, a periodic structure (which was later realized to be the secondary instabilities) does exist along the spanwise direction. In their flow visualization too, they observed large eddies and smaller lumps superimposed on them, giving the former a rough, corrugated appearance. Furthermore, they observed that two-dimensional forcing improves the spanwise regularity of the large eddies but does not remove the corrugations. This suggests that the periodic structures in the spanwise direction are not influenced by two-dimensional forcing. This is expectable as such a forcing does not exercise control over smaller-scales which seem to induce these structures along the span of the mixing layer.

Breidenthal (1981) studied the structures in a mixing layer using chemical reaction. They agreed that large-scale two-dimensional structures dominated both mixing and the distribution of the mixed fluid in the shear layer. However, in their flow visualisation, they observed a spanwise-sinuous disturbance which they called *wiggle*. The average wavelength of the wiggle was measured to be about 1.1 times the wavelength of the initial KH instability (primary fundamental instability). While the shear layer is two-dimensional upstream of this wiggle, streamwise lines were visible downstream and the flow was three-dimensional. These streamwise streaks appeared at the onset of the wiggle, with their amplitude-maxima a short distance downstream. Furthermore, they observed that the outer sheet of the two-dimensional vortices were corrugated while the inner sheet was relatively unperturbed. While the presence of streamwise vortices was confirmed by this work, the cause of these structures were still unclear.

The two- and three-dimensional stability characteristics of the shear layer was computed by Pierrehumbert & Widnall (1982). They employed a non-parallel flow field unlike Michalke (1964, 1965a), and defined it as a periodic array Stuart (1967) vortices to represent the spanwise vortices. The stability equations used was an extension of the Rayleigh problem to account for the strong non-parallel flow field and was solved numerically using spectral methods. They found two classes of instabilities. The first being the sub-harmonic instability which demonstrate highest growth rate when there is no spanwise variation. Hence this class of instability is twodimensional and corresponds to vortex pairing of the spanwise vortices. The second class of instability repeats itself in the streamwise direction, at the same periodicity as the undisturbed spanwise vortex array. These disturbances lead to an infinitesimal translation of the primary vortex cores and hence was recognised by the authors as a translative instability which results in the streamwise streaks observed in experimental investigations. These disturbances translate the primary vortex cores variably along the spanwise direction, resulting in the sinusoidal wiggle observed by Breidenthal (1981). Finally, they found that instabilities with a spanwise wavelength that is $2/3^{rd}$ of the separation between vortices arising from the primary instability mechanism is most unstable. Although, growth rate as a function of spanwise wavelength shows the broadband nature of these instabilities.

Bernal & Roshko (1986) investigated the dynamics of these streamwise streaks observed along the span caused by secondary instability mechanisms, their impact on the development of threedimensionality in the mixing layer and their contribution to the growth of the mixing layer itself. The onset of these streaks and their spanwise wavelength were affected by initial conditions as previously observed (Wygnanski, Oster, et al., 1979; Breidenthal, 1981). Furthermore, these instabilities were observed to grow in amplitude in the streamwise direction. These secondary instability modes generate a counter-rotating streamwise vortex structure superimposed on the primary spanwise vortices, that is seen in the plane view as streamwise streaks in figure.

The experiments of Bernal & Roshko (1986) showed that the large-scale spanwise vortices still resembled those generated by the KH instability (as in figure 2.17 and edge-view in figure 2.20), while the streamwise vortices had a hairpin structure similar to that shown in figure 2.21a. The authors proposed that these secondary instabilities form a warped vortex line, threading its way up- and downstream of adjacent spanwise vortices, thus changing its streamwise direction on each pass. Additionally, fluid from the high and low velocity streams are entrained into the primary vortex core whenever the induced velocity of these secondary vortices is oriented towards the former. As a result, the primary vortex also grows between subsequent instances of vortex pairing and amalgamation. It was also observed that streamwise vortices play a role in the mixing transition process. These secondary vortices become part of the primary vortices in subsequent amalgamations and expedite the aforementioned mixing. Thus, a combination of



Figure 2.21.: Plane view of the mixing layer after the secondary vortices wrap around the large-scale spanwise vortices. (b) Topology of the vorticity field between two consecutive spanwise vortices. Reproduced from Lasheras & H. Choi (1988).

the three-dimensional secondary instabilities and the process of spanwise-vortex amalgamation lead to transition in mixing layers.

Lasheras & H. Choi (1988) experimentally investigated the three-dimensional development of of a plane mixing layer subjected to small-amplitude sinusoidal spanwise perturbations. This was achieved by modifying the splitter-plate trailing-edge to give it a sinusoidal shape with a particular wavelength and amplitude. As the primary KH instability was not forced, the shear layer developed naturally, forming two-dimensional spanwise vortex tubes. Through flow visualisation, the authors observe that the development of both these vortical flow structures are independent initially. However, further downstream, the streamwise vortices induce a wavy undulation on the spanwise vortex cores. This wavy undulation of the spanwise vortices is nothing but the *wiggle* observed previously by Breidenthal (1981). The wavelength of this undulation is the same the applied spanwise perturbation, although with a 180° phase shift. Eventually, the streamwise vortices get entangled around the spanwise vortex as shown in figure 2.21. Spanwise perturbations with a wide range of wavelengths were applied but no particular mode was found to be most unstable. The authors thus conclude that secondary instabilities have a large bandwidth of unstable modes that all grow at a comparable rate.

Spanwise phase distributions through an array of heaters was employed in the experiments of Nygaard & Glezer (1991), to excite secondary instabilities into the their mixing layer flow field. Upstream of this spanwise array, they employed an array of spanwise-uniform heaters to excite the fundamental (primary) KH instability simultaneously. The obtained results demonstrated that when the spanwise wavelength of applied excitation was greater than the streamwise wave-

length of the fundamental KH instability, undulation of the spanwise vortices as a result of the translative instability was observed. Furthermore, the bending of the spanwise rollers resulted in concentration of vorticity on either the top or bottom of the roller at the bends which is caused by the stretching or compression of the vorticity at these locations. When the spanwise wavelength was lower than the streamwise wavelength, this undulation was not observed.

Huang & Ho (1990) investigated the generation of small-scale turbulence in an initially laminar mixing layer experimentally. They studied the evolution of both spanwise and streamwise vortical flow structures and how each of them influence mixing transition. They observed that the interaction of these two vortical flow structures during a vortex pairing resulted in the production of small-scale turbulence. Additionally, the results also demonstrated that pairing resulted in the doubling of both, the spanwise wavelength of the streamwise vortices and the streamwise wavelength of the spanwise roller. Thus, the ratio between them remains constant. Finally, they also indicate that the local Reynolds number must be high enough to withstand small-scale turbulence.

Moser & Rogers (1991) and Rogers & Moser (1992) performed DNS of a time-developing mixing layer with different initial conditions to study the effect of these on mixing transition and the breakdown of flow structures prevalent in a three-dimensional mixing layers. As postulated by Bernal & Roshko (1986), the authors observe here that when the spanwise vortices pair, the ribs between them are engulfed into the resulting vortex. This results in the breakup of the resulting vortex followed by secondary roll-up of thin vortex sheets, ultimately resulting in transition to small-scale turbulence. Thus, similar to Huang & Ho (1990), they concluded that three components are necessary for transition to occur, which include a Reynolds number must be sufficiently high, a set of spanwise vortices and rib vortices between them, and finally a pairing event. The authors also comment that if the three-dimensional disturbances are not strong enough initially, mixing transition might occur after the second pairing event or more. Furthermore, they also observe that if the circulation of rib vortices is high enough, it can induce non-linear effects. This non-linearity results in the collapse of rib vortices into streamwise streaks in experimental investigations. It also causes concentration of vorticity into thin cup-shaped structures at the bends of the undulated spanwise rollers that was observed in the experiments of Nygaard & Glezer (1991). The authors also show the appearance of spanwise vorticity with opposite sign compared to that of the mean vorticity.

2.3.2. Mixing layer growth

The cross-stream width of the mixing region, often referred to as spread or growth of the mixing layer has been investigated thoroughly and holds utmost importance. The different shear layer instabilities, by governing the evolution and interaction of large-scale vortices and by modulating the rate of entrainment, control the manner in which the mixing layer spreads and also its rate. Thus, this parameter is an indication of the size and evolution of the large-scale vortices and as such can give some information about their dynamics (especially when some control is applied, discussed later in section 2.3.3). Furthermore, naturally occurring mixing layers spread linearly along the streamwise direction. This linear spread is one of the prerequisites for the mixing layer to have achieved self-similarity as has been reported by many research efforts (see for example Liepmann & Laufer, 1947; Wygnanski & Fiedler, 1970; Spencer & Jones, 1971; Brown & Roshko, 1974; Oster & Wygnanski, 1982; Ho & Huang, 1982; Huang & Ho, 1990).

This spatial evolution of the mixing layer has been measured through a number of parameters.

The most rudimentary one is the visual thickness δ_{vis} (see Brown & Roshko, 1974, figure 2.17) which is obtained by measuring the width of the layer from flow visualisation pictures. This is very similar to another width parameter $b = y_{0.95} - y_{0.05}$. Here, $y_{0.95}$ and $y_{0.05}$ are the distance along the cross-stream direction measured from the splitter-plate trailing edge where the streamwise velocities of the flow are $U_1 + 0.95\Delta U$ and $U_1 + 0.05\Delta U$ respectively, with $\Delta U = U_2 - U_1$.

A parameter to measure the growth of the mixing layer that is independent of the similarity of velocity profiles at different streamwise locations like the previous two parameters is the mixing layer integral thickness which is often referred to as *momentum thickness*, computed as shown below, with \overline{u} representing the velocity profile at a particular streamwise location.

$$\theta = \int_{-\infty}^{\infty} \frac{\overline{u} - U_1}{\Delta U} \cdot \left[1 - \frac{\overline{u} - U_1}{\Delta U} \right] dy$$
(2.5)

The growth of the mixing layer can also be quantified with the velocity-profile maximum slope thickness (see Spencer & Jones, 1971; Brown & Roshko, 1974). This parameter can also be interpreted as the *vorticity thickness* (δ_{ω}) and is computed according to equation 2.6.

$$\delta_{\omega} = \frac{\Delta U}{(d\overline{u}/dy)_{max}} \tag{2.6}$$

This parameter is used extensively in the current investigation to quantify the growth of the mixing layer flow field. The motivation for this is that the development of the mixing layer is a product of the unstable motion caused by the vorticity distribution, which manifests itself as large-scale coherent vortices in the flow field (Brown & Roshko, 1974). Given that these vortices, by entrainment of irrotational fluid or by amalgamation (Winant & Browand, 1974), influence the growth of the mixing layer, it is only appropriate to use the vorticity thickness for this purpose.

Finally, another parameter that is used to quantify mixing layer growth is σ (see Birch & Eggers, 1972; Bell & Mehta, 1990; Mehta, 1991). This is obtained by fitting a standard curve (such as the *tanh* or error-function) on to the velocity profile. The parameter σ is then chosen for an overall best fit. This parameter σ is in fact an inverse measure of the spreading rate. Compared to the vorticity thickness, this method does take into account features of the entire velocity profile. However, given that the thickness of the mixing layer is estimated through a relation to a shape function, it is dependent on the quality of the fit. Thus, this method is tedious and can be more prone to errors.

A method to quantify and compare the spreading/growth rate of the mixing layer is to equate it to the velocity ratio $R = (U_2 - U_1)/(U_2 + U_1) = \Delta U/2\overline{U}$ based on the Abramovich (1963)-Sabin (1965) rule given by

$$\frac{d\delta_{\omega}}{dx} = c \cdot \frac{U_2 - U_1}{U_2 + U_1} = c \cdot R \tag{2.7}$$

where the constant c is a fitting parameter. This operation can be carried out on the layer width b (or visual thickness), the vorticity thickness δ_{ω} or the thickness estimated by fitting a standard

curve on the measured velocity profile, but not on the integral momentum thickness θ .

A large variation of this spreading parameter c has been reported by different experimental efforts, ranging from 0.16 (Liepmann & Laufer, 1947; Spencer & Jones, 1971) to 0.27 (Abramovich, 1963; Huang & Ho, 1990). This large spread in the mixing layer growth rate can be attributed to the receptivity of the mixing layer flow field to a number of experimental parameters. These include size of the test section which, if small, could lead to the interaction of the mixing layer with the test-section walls. Free-stream turbulence level too has a significant influence. A higher level of turbulence results in greater entrainment and thus a higher growth rate (Chandrsuda et al., 1978; Patel, 1978; Pui & Gartshore, 1979). The state of the incoming splitter-plate boundary layers, being laminar or turbulent, also affects the growth rate of the ensuing mixing layer. Mixing layers ensuing from turbulent boundary layers grow at a lower rate but are known to achieve self-similarity sooner (further upstream) compared to those originating from boundary layers at a laminar state (Browand & Latigo, 1979; Wygnanski, Oster, et al., 1979; Bell & Mehta, 1990). Vibrations of the splitter-plate itself (Pui & Gartshore, 1979) and the trailing-edge geometry (Dziomba & Fiedler, 1985; Braud et al., 2004; Laizet et al., 2010) also have an influence on the mixing layer development.

2.3.3. Mixing layer flow control

The research efforts reported until now in this review have shown that the interaction of the fundamental KH instability wave with its sub-harmonic leads to sequential merging of the spanwise large-scale vortices in the mixing layer flow configuration. The occurrence of this phenomenon, in the presence of streamwise vortices arising from the secondary instability mechanism, results in the occurrence of small-scale transition. This eventually can enhance molecular mixing. The logical extension of this coherent flow structure research is to control them in order to reap some specific benefits.

Application of flow control over the mixing layer configuration is based on introducing smallamplitude perturbations to manipulate the stability characteristics of the flow field. This way, the evolution and dynamics of coherent flow structures can be controlled, thereby realizing the desired end results. A number of devices have been employed in the past to control these instabilities. These include mechanical devices such as vibrating ribbons or oscillating flaps, acoustically by means of loudspeakers, thermal perturbations by means of heating elements and suction or blowing devices, to name a few. A synthesis of different methods to control the mixing layer flow field can be found in the reviews of Ho & Huerre (1984) and Fiedler (1998).

The effect of two-dimensional excitation through a thin oscillating flap at the splitter-plate trailing edge, on the development of the mixing layer, was experimentally studied by Oster & Wygnanski (1982). They demonstrated a successful control of the mixing layer flow field with this excitation method. Such mechanical means of forcing hinges around the input energy of the mechanical device (either an oscillating flap or vibrating ribbons) being converted into instability waves, which in turn manipulates the stability of the mixing layer flow field. While the working mechanism of such mechanical devices are simple, the main constraint is that the forcing frequencies that can be realized is limited and hence, is not widely applicable. Furthermore, only single-frequency excitation can be performed. Finally, only control of the primary KH instability mechanism can be achieved with these devices as the applied forcing is two-dimensional. Thus, while the large-scale spanwise flow structures can be controlled, the streamwise vortices arising from the secondary instability mechanism is not affected.



Figure 2.22.: Growth of mixing layer represented by the momentum thickness, reproduced from Ho & Huang (1982). Growth of the unforced mixing layer (dashed curve), and two mixing layers forced at $f_f/f_m = 0.38$ (\bigcirc) and $f_f/f_m = 0.43$ (\triangle). f_m is the most-probable frequency of vortex passage initially and is equal to the frequency of the fundamental KH instability f_0 . Regions I, II and III indicated are with respect to $f_f/f_0 = 0.38$ forcing case.

Disregarding these drawbacks, Oster & Wygnanski (1982) observed that, when sub-harmonic instabilities are excited, the evolution of the forced mixing layer can be divided into three regions as shown in figure 2.22. In region I and III, the mixing layer grows linearly due to amalgamation. While in region II, the mixing layer does not grow due to inhibition of vortex amalgamation. By comparing the results from other research efforts, especially Ho & Huang (1982, discussed previously in section 2.3.1.1), a few additional conclusions can be drawn. In region I ($St \cdot R < 1$, St being the Strouhal number), a phase difference exists between the sub-harmonic and the fundamental frequencies in the mixing layer, which promotes vortex pairing and leads to linear growth of the mixing layer as discussed briefly by the authors. Region II ($1 < St \cdot R < 2$) is characterized by no vortex pairing or amalgamations. The authors observe that in region II, the spectral energy in the sub-harmonic frequencies are very low compared to the fundamental instability which impedes amalgamation. The sub-harmonic starts to grow again close to $St \cdot R > 1.5$ and becomes prominent at the start of region III ($2 < St \cdot R$), re-starting vortex amalgamation. However, the spectra now shows energy in a broad frequency band suggesting the occurrence of random amalgamations.

A second method that has been widely used for exciting small disturbances within the shear layer is through acoustic waves generated by a loudspeaker. The loudspeaker is fed with a sinusoidal signal with the necessary excitation frequency which results in the generation of acoustic waves/perturbations. These waves are guided through tubes and introduced directly into the flow, close to the splitter-plate trailing edge (or at the nozzle lip in jet flows). Such a method has proved useful to excite the desired instabilities in both, the mixing layer flow field (see Huang & Ho, 1990) and jet flows (see Laufer & Zhang, 1983; Husain & Hussain, 1995). A less common but still effective method is to introduce the acoustic perturbations into the free-stream, at an upstream location (see Thomas & Chu, 1989).

Instability excitation by acoustic means has a few advantages over the use of mechanical devices. Firstly, instabilities in a broad frequency spectrum ranging from just a few hertz to tens of kilohertz can be excited through acoustic perturbations. A wide range of forcing amplitudes too can be realized. Furthermore, double-frequency excitation is also possible as demonstrated by Husain & Hussain (1995, discussed previously in section 2.3.1.2). However, similar to mechanical devices, acoustic perturbations can only control the primary KH instability mechanism and the corresponding large-scale coherent flow structures but cannot manipulate the secondary instabilities.

The common shortcoming of both the above control methods is their inability to manipulate secondary instability mechanisms. One method to excite these instabilities is by modifying the shape of the splitter-plate trailing edge. As discussed in section 2.3.1.3, Lasheras & H. Choi (1988) employed a splitter plate whose trailing edge bore a sinusoidal shape (with a particular wavelength and amplitude) along the span and successfully manipulated the secondary instability mechanisms in their mixing layer flow field (see figure 2.21a). However, this is a passive flow control method and the disadvantage here is the opposite of what was observed with mechanical and acoustic devices. While secondary instability mechanisms can be influenced, no control is realized over the primary KH instability mechanism which is equally (if not more) important.

The first approach that has been instrumental in controlling both these instability mechanisms is the use of surface-mounted heating elements. Surface heaters have a high frequency response and employ no moving parts. They can be easily mounted on any surface while maintaining minimum intrusion to the flow and are able to impart both, temporal and spatial perturbations. This method is based on introducing vorticity perturbations by exploiting the dependence of fluid viscosity on temperature (demonstrated by Liepmann, Brown, et al., 1982). As discussed previously in section 2.3.1.3, Nygaard & Glezer (1991) employed surface-heating elements to control mixing layer instabilities. They applied two arrays of surface heaters over the splitter plate on the high-velocity stream's side, just upstream of the trailing edge. The first array comprised of spanwise-uniform heating elements whose purpose was to excite the primary KH instabilities. While the second array consisted of spanwise row of smaller (spanwise-width) heating elements placed along the span with a particular spanwise spatial wavelength. These arrays were operated by supplying them an electric signal with an arbitrary temporal waveform to manipulate the primary instability mechanism.

While Nygaard & Glezer (1991) demonstrated the capability of such a forcing to manipulate the vorticity distribution in the high-velocity boundary layer and thus the vortical structures arising from both the primary and secondary instability mechanisms, use of thermal excitations to perturb the shear layer instabilities does have certain limitations. Perturbations induced into the splitter-plate boundary layer will amplify or decay based on the layer's stability characteristics. Hence, care must be taken while selecting the location of heating elements, forcing frequencies and amplitudes so as to ensure that the introduced perturbations do not decay before reaching the splitter-plate trailing edge. Therefore, these devices might require a significant amount of power to perform effective control.

Mathis et al. (2007) attempted controlling the mixing layer by inducing a boundary-layer separation just upstream of the splitter-plate trailing edge. A pneumatic injection system that produces a steady (blowing) jet perpendicular to the splitter-plate surface introduced the necessary perturbations to cause boundary-layer separation on the high-velocity side of the splitter plate. The controlled separation was observed to result in high turbulent quantities and longitudinal vorticity. The steady blowing leads to strong unsteady three-dimensional perturbations



Figure 2.23.: Smoke/laser visualisation of the mixing layer reproduced from Parezanovic et al. (2014). The unforced flow is shown in (a) while the re-organisation of large-scale vortices when low-frequency (10 Hz) fluctuations are imparted through controlled blowing is shown in (b).

which enhances mixing. While this is an interesting approach, controlling the instability growth might not be feasible as this method is based on a steady input.

Controlled blowing was employed by Parezanovic et al. (2014) to excite the primary instabilities in the mixing layer flow field. In their experiment, the splitter-plate trailing edge contained 96 circular nozzles along the span. Compressed air for these nozzles were supplied from a pressurised plenum chamber. The pressure difference between the plenum chamber and the flow field controlled the mass-flow rate through each nozzle and thus defined the amplitude of the forcing. Between each nozzle and the plenum chamber, a micro-valve was installed. These micro-valves were driven simultaneously with a square-wave signal, which controlled the frequency and duty cycle of the actuation. The authors applied forcing at different frequencies and demonstrated effective control of the evolution and interaction of the large-scale coherent vortices arising from the primary KH instability mechanism (see figure 2.23). While not attempted by the authors, manipulation of secondary instabilities can also be achieved by controlling the activation of different micro-valves along the span. After the success with controlling the primary instabilities with these actuators, the authors also employed these actuators in a closed-loop control algorithm of feedback type (also see Parezanović et al., 2016). This is not elaborated upon here as closed-loop flow control is outside the scope of the current thesis.

Finally, recent efforts of Little & co-workers (Ely & Little, 2013; Singh & Little, 2020) on the application of AC-DBD plasma actuator forcing to control the primary instability and thus, the dynamics of spanwise vortices, in an incompressible turbulent mixing layer needs to be mentioned. While the frequency of the fundamental KH instability of their mixing layer flow configuration is not clearly reported, the authors excite a low-frequency primary instability that is most unstable at a specific streamwise location from the splitter-plate trailing edge. They demonstrated the mixing layer flow configuration is sensitive to momentum-based forcing applied by the DBD plasma actuator. The actuation alters the growth of the layer which was indicated by the momentum thickness (see figure 2.24a). It was also concluded to impact the vortical flow structures through the phase-averaged traverse velocity fluctuation fields (figure 2.24b). It is worth noting here that most of the results presented by Singh & Little (2020)pertains to imparting momentum into the low-velocity splitter-plate boundary layer. Although, the AC-DBD actuator was shown to have more authority when the momentum is input into the high-velocity boundary layer by Ely & Little (2013) (figure 2.24a). This is however expectable given that this boundary layer dominates the stability characteristics of the flow field (Ho & Huang, 1982). Furthermore, given the importance of vortex pairing or amalgamation in the dynamics of mixing layers, the influence of sub-harmonic forcing received surprisingly little attention as well.



Figure 2.24.: (a) Momentum thickness variation along the stream with different actuation location and voltage amplitudes, reproduced from Ely & Little (2013). (b) Phase-averaged transverse velocity fluctuation fields when plasma body forcing is applied into the low-velocity splitter-plate boundary layer, reproduced from Singh & Little (2020).

3

METHODOLOGY

The research presented in this thesis is mainly derived from experimental investigations which is backed by some theoretical computations. This thesis comprises of investigating the effect of plasma-based flow control, more precisely instability control, over two different flow configurations. The investigations pertaining to the swept-wing boundary layer and the corresponding cross-flow instability control was carried out at AWEP Department, Faculty of Aerospace Engineering, Delft University of Technology, The Netherlands. Investigations related the plane mixing layer flow configuration on the other hand was performed at Institut PPrime, Université de Poitiers, France.

In this chapter, the different tools employed to carry out the current research is introduced. The experimental test facilities, which include the different wind-tunnels utilized and the models employed are first described (section 3.1). The different systems deployed to measure the necessary physical quantities are introduced in section 3.2. These include infrared (IR) thermography to analyse the evolution of laminar-to-turbulent transition over the swept-wing, particle image velocimetry (PIV) to obtain quantitative data of respective flow field's velocity and the different methods followed to characterise the performance of DBD plasma actuators. Finally, various analytical tools used to investigate the acquired experimental data is discussed in section 3.3.

3.1. EXPERIMENTAL FACILITIES

3.1.1. Wind-tunnel and model for swept-wing boundary layer investigation

Swept-wing boundary layer experiments were carried out in the closed-loop, Low-Turbulence Tunnel (LTT) at Delft University of Technology, The Netherlands. The tunnel's testing chamber dimensions are 1.25 m × 1.80 m × 2.6 m in height, width and length, respectively. The tunnel is equipped with seven anti-turbulence screens followed by a convergence with a contraction ratio of 17 : 1, which results in a flow with very low levels of free-stream turbulence. A recent characterisation of the flow in this wind-tunnel revealed the free-stream turbulence to be $Tu/U_{\infty} = 0.024\%$ at a free-stream velocity $U_{\infty} = 24$ m s⁻¹ (measured with a single hot-wire probe, bandpass filtered between 2 and 5000 Hz, see Serpieri, 2018, PhD thesis).

The model (referred to as M3J) used in this thesis is a 45° swept-wing. The airfoil used is a modified version of the NACA 66018 shape named 66018M3J and is as shown in figure 3.2a.



Figure 3.1.: The low-turbulence wind tunnel at at Delft University of Technology, The Netherlands.



Figure 3.2.: (a) The 66018M3J airfoil shape (orthogonal to the wing's leading edge) of the swept-wing model employed in the current thesis. Spatial coordinates are non-dimensionalised with the chord length perpendicular to the leading edge ($c_x = 0.898$ m). (b) Schematic of the swept-wing model (not to scale), chord length in the streamwise direction (c) and perpendicular to the leading edge (c_x), the inviscid streamline (*ISL*) and the coordinate systems are shown.

The wing model is such that the bottom wing-tip is upstream compared to the wing-tip on the top. The model has a of span 1.25 m and streamwise chord of c = 1.27 m. The chord length perpendicular to the leading edge is $c_x = 0.898$ m. The model is made of fibreglass with a polished surface, featuring a surface roughness standard deviation of 0.2 μ m. The wing model is equipped with two arrays of 46 surface pressure taps along the free-stream direction, each of which was positioned 300 mm from the top and bottom walls of the wind tunnel. These taps were connected to a transducer to measure the pressure distribution along the free-stream direction.

This model is designed such that, it can facilitate an experimental study of the CF instability. The wing model (schematic in figure 3.2b) was designed with a small leading-edge radius of about 1% of chord to avoid attachment-line instability (Bippes, 1999). At zero-incidence, the airfoil shows accelerating flow and a favourable pressure gradient $(\frac{\partial p}{\partial x} > 0)$ until about 63% of chord. Thus the amplification of Tollmien-Schlichting (TS) waves is avoided. Furthermore, the



Figure 3.3.: Pressure distribution on the swept-wing model's pressure side measured by two arrays of 46 pressure taps installed 300 mm from the bottom and top walls of the wind-tunnel at $\alpha = 3^{\circ}$ and $U_{\infty} = 25$ m s⁻¹ ($Re_c = 2.1 \cdot 10^6$). X represents distance from the leading edge along the streamwise direction.

model has no concave surfaces to avoid Görtler vortices. The isolation of these three instability mechanisms, combined with low surface roughness of the employed model and the low turbulence level in the wind tunnel makes this experimental set-up very appropriate for the investigation of CF instability and the laminar-to-turbulent transition caused by it.

Finally, a fundamental study on CF instability requires the model to represent an infinite swept-wing, with the flow being invariant along the span (Deyhle & Bippes, 1996; Reibert et al., 1996; Radeztsky et al., 1999). At a wing incidence angle of $\alpha = 3^{\circ}$ and free-stream velocity of $U_{\infty} = 25 \text{ m s}^{-1}$ ($Re_c = 2.1 \cdot 10^6$, which correspond to the flow conditions predominantly investigated in this thesis), no difference is noticeable in the pressure distribution on the wing's pressure side, measured by the two arrays of surface pressure taps, as shown in figure 3.3. It is worth noting that the horizontal axis in this plot represents the non-dimensional distance from the swept-wing leading edge along the streamwise direction. This figure indicates a high degree of uniformity of the flow along the span, signifying the realization of the infinite swept wing hypothesis at these flow conditions.

3.1.2. Wind tunnel for plane mixing layer investigation

The plane mixing layer experiments were carried out in a closed-loop wind-tunnel, specially modified to accommodate mixing layer studies at *Institut PPrime*, *Université de Poitiers*, *France*. The dimensions of the test-section of this wind-tunnel was 30 cm \times 30 cm \times 1.2 m in height, width and length (streamwise direction) respectively. A 15 mm thick aluminium (splitter) plate extended from the entrance of the settling chamber, just after the honeycomb, through the convergent and until the entrance of the test section, as illustrated in figure 3.4a (thick red line). The aluminium plate split this section of the wind-tunnel into two equal halves along the vertical direction, thereby generating two parallel flow streams. The set-up was designed such that the free-stream velocity of the top stream (U_2) was different compared to that of the bottom stream (U_1), so as to realize a plane mixing layer configuration arising between two parallel flow streams.

One metallic grid (with wire diameter of 0.18 mm and solidity of 35%) was installed just upstream of the convergent in the top flow stream to breakdown the incoming vortices and



Figure 3.4.: (a) Wind-tunnel set-up. Location of the foams and perforated plate (dark rectangle), metallic grids (four vertical lines) and the splitter plate (red) are depicted. (b) Schematic (not to scale) of PMMA attachment (light grey) fixed to the aluminium plate (black) to complete the splitter-plate. Spanwise incision (dark grey) to house an electrode of the DBD plasma actuator is shown. Coordinate system used in this study is presented in both figures.

reduce the free-stream turbulence in this stream. Turbulence intensity was computed from velocity measurements acquired with a single hot-wire sensor (bandpass filtered between 10 and 1000 Hz). The turbulence intensity of the high-velocity (top) stream at $U_2 = 10 \text{ m s}^{-1}$ was found to be $Tu_2/U_2 = 0.22\%$. In order to realize a difference in the velocity between the two streams, flow manipulators were fitted in the bottom stream, just upstream of the convergent (black rectangle thick vertical lines in figure 3.4a). A combination of foams and a perforated plate was installed here to reduce velocity of this stream. A set of four metallic grids (same as in top stream), separated from each other by 8 mm, was installed downstream of the perforated plate to reduce the level of free-stream turbulence in this stream. These manipulators resulted in the velocity of the bottom stream being half that of the top stream. Thus, the velocity ratio between the two streams was measured with the same hot-wire. When the velocity of the bottom stream was measured with the same hot-wire. When the velocity of the bottom stream was $U_1 = 10 \text{ m s}^{-1}$ ($U_2 = 20 \text{ m s}^{-1}$), the turbulence intensity (bandpass filtered between 10 and 1000 Hz) was $Tu_1/U_1 = 1.5\%$. At $U_1 = 5 \text{ m s}^{-1}$ ($U_2 = 10 \text{ m s}^{-1}$), the measured turbulence intensity was $Tu_1/U_1 = 1.3\%$.

The end of the aluminium plate was attached with a PMMA attachment (schematic shown figure 3.4b) which completed the splitter plate. This attachment had a 120 mm long taper of 3.1° on both the top and bottom sides, which resulted in a 3 mm thick trailing-edge. The use of PMMA as the material for this attachment is motivated by its additional application as the dielectric for the DBD plasma actuator. To this effect, a 40 mm wide (along the streamwise direction) spanwise incision was produced on its bottom side (dark grey area in figure 3.4b) such that 2 mm thick PMMA remained on the top side, which served as the dielectric substrate. Details of the design and construction of the different DBD plasma actuators employed are introduced in section 4.5. The incision was filled with an epoxy resin before performing surface treatment to obtain a smooth and uninterrupted surface.

3.1.3. Coordinate systems

A number of coordinate/reference systems are utilized in this thesis based upon the flow fields being investigated. These different coordinate systems are introduced here.

3.1.3.1. DBD plasma actuator

The coordinate system corresponding to the induced flow of the DBD plasma actuators and other derived quantities is represented in this thesis by x_p , y_p and z_p . It is such that the x_p axis is along the induced jet/flow, y_p is perpendicular to the dielectric substrate and z_p is along the span of the actuator. The geometrical origin is at the plasma-generating edge of the air-exposed electrode, at mid-span. The corresponding velocities are represented by u_p , v_p and w_p .

3.1.3.2. Swept-wing boundary layer

The coordinate systems employed in the investigation of the swept-wing boundary layer control are as follows (see figure 3.2b). The (wing) coordinate system used in this study is represented by x, y and z. It is such that its x axis is perpendicular to the leading edge of the wing and along the chordwise direction, y axis is perpendicular to the chord plane and z axis is parallel to the leading edge. The geometrical origin with respect to the x and y axis is set at the leading edge, while the coordinate z = 0 can be moved freely along the leading edge due to the infinite swept wing condition. The corresponding velocity components are represented by u, v and w. It is noteworthy that, as the swept-wing model has an incidence angle of $\alpha = 3^{\circ}$ for the entire experiment, the wind-tunnel coordinate system does not coincide with the wing coordinate system. Thus, the free-stream velocity in the wind-tunnel test section is referred to as U_{∞} .

A second coordinate system which is locally aligned with the flow inviscid streamline is employed to express the velocity components within the boundary layer. This system is represented with the subscript $_{ISL}$, where x_{ISL} is along the local inviscid streamline, z_{ISL} is orthogonal to it and positive towards the leading edge (opposing the CF component) and y_{ISL} is orthogonal to the local surface curvature. The velocity components are represented by $(uvw)_{ISL}$, with w_{ISL} being the magnitude of the CF within the boundary layer. The inviscid streamline angle φ_{ISL} is defined with respect to x (from the wing coordinate system) and is such that, it becomes 90° when the flow is along the z direction.

3.1.3.3. Plane mixing layer

The coordinate system used in this thesis to investigate the plane mixing layer flow configuration is represented by x, y and z, where x is along the free-stream direction, y is perpendicular to the plane of the mixing layer and z is along the span. The geometrical origin of the mixing layer is fixed at the trailing-edge of the splitter plate as shown in figure 3.4b, with z = 0 being at the centre of the splitter plate along the span (15 cm from both walls). The corresponding velocity components are represented by u, v and w respectively.

3.2. EXPERIMENTAL DATA ACQUISITION

A number of measurement systems were deployed during the course of this thesis to quantify the relevant physical quantities. These include various velocity field measurements, laminar-toturbulent transition location over the swept wing and different performance parameters of DBD



Figure 3.5.: (a) IR thermography set-up installed outside the wind-tunnel test section. Six halogen lamps irradiating the model (dashed red rectangles) and the IR camera (dashed white rectangle) are also shown. (b) Schematic of the swept-wing model with the FOVs of IR thermography. Solid red rectangle: *global IR-FOV*; dashed red rectangle: *zoomed-in IR-FOV*. The coordinate system is also shown.

plasma actuators. The instrumentation and set-up of the measurements systems employed for these purposes are introduced here.

3.2.1. Swept-wing boundary transition topology through infrared thermography

A key necessity in this thesis is to visualize and investigate the topology of the boundary layer ensuing over the swept wing and its eventual laminar-to-turbulent transition. An often used technique to visualize flow transition topology is oil-flow visualisation (see for example figure 2.10a, reproduced from Serpieri & Kotsonis, 2015). This method involves the application of some luminescent viscous fluid over the aerodynamic body's surface. The shear forces in the flow displaces the viscous fluid accordingly, which is then photographed and studied. While this method is simple and relatively fast, longer tunnel time is necessary to obtain evident displacements of the viscous fluids. Care must be taken while selecting the viscous fluid as the effect of other forces such as gravity, pressure forces or surface tension must be significantly lower compared to the flow shears. Finally, this is an intrusive technique in a certain manner as it necessitates constant entry into the test section to re-apply/re-distribute the viscous fluid. These reasons make this technique relatively cumbersome to visualise the flow organisation.

Infrared (IR) thermography on the other hand pivots around employing specific sensors to capture the IR radiation emitted by a body. The aerodynamic body being tested is heated continuously by an external source (such as printed circuits heating the model by Joule effect, halogen lamp radiations, etc.) and cooled down by the flow through convection. In low-speed conditions, flow visualisation through IR thermography can be considered to be an active technique. Furthermore, it is also non-intrusive which is an important attribute for utilisation in swept-wing boundary layer investigations, as even minute particles introduced into the test facilities can lead to significantly different boundary layer evolution and transition due to the receptivity mechanism. Thus, in the current thesis, IR flow visualisation is extensively used to inspect the swept-wing boundary-layer's transition evolution. It is worth noting however that heat conduction through the model surface must be negligible as this helps ascertain the effect of the flow better, by increasing the thermal sensitivity of the measurement system with respect to the flow. The swept-wing model employed in this thesis is made of fibreglass which is not a good thermal conductor. Thus, it favours the use of infrared thermography for flow visualisation.

It is known that the wall shear stress resulting from a turbulent boundary layer is higher compared to its laminar counterpart. Following Reynolds analogy, surface heat transfer coefficient is higher in the regions where the flow is turbulent. In the case where the airflow is at a lower temperature than the wing surface this will result in enhanced cooling (i.e. lower surface temperature) in turbulent regions compared to regions on the model where the flow is laminar. Thus, bright areas in the obtained IR fields correspond to the regions on the swept wing where the flow is in laminar state while the darker areas correspond to a turbulent boundary layer (see figures 3.6b and c). Additionally, stationary CF vortices cause regions of low shear at their upwelling region and high shear close to the downwelling region. Thus, if the thermal sensitivity and spatial resolution of the measurement system is high enough, these vortices appear as alternating streaks of high and low temperature as seen in the IR field in figure 2.15b. Thus, the spanwise-modulation of the boundary layer by the stationary CF instability and its corresponding wavelength can be extracted from these fields.

The set-up of the IR thermography system is depicted in figure 3.5a. The swept-wing model was irradiated with six 1 kW halogen lamps (dashed red rectangles in figure 3.5a) through the glass windows of the test section to maintain a constant heat flux and enhance contrast. The IR camera (dashed white rectangles in figure 3.5a) used during the experiment was an Optris PI640. It features a $640 \times 480 \text{ px}^2$, un-cooled focal plane array that responds to the electromagnetic spectral range between $7.5 - 13 \ \mu m$. Its thermal sensitivity (NETD) is 75 mK. The IR camera too was placed outside the wind tunnel, about 1 m from the swept-wing model and imaged it through an opportune window. Two different fields-of-view (FOVs) were captured to inspect different boundary layer features which are depicted in figure 3.5b. For the first FOV, the IR camera was equipped with a $15^{\circ} \times 11^{\circ}$ wide angle lens with a focal length of 41.5 mm in order to capture the flow over the entire swept-wing model and detect laminar-to-turbulent transition (solid red rectangle in figure 3.5b, henceforth referred to as *global IR-FOV*). However, the achieved spatial resolution of this measurement night not be adequate to inspect the wavelength of the stationary CFI modes being considered here. To this effect, a second zoomed-in FOV was captured by equipping the IR camera with a $60^{\circ} \times 45^{\circ}$ lens with a focal length of 10.5 mm (dashed red rectangle in figure 3.5b, henceforth referred to as *zoomed-in IR-FOV*).

The camera acquired 100 IR fields at a rate of 5 Hz for every tested flow case which is averaged to increase the signal-to-noise ratio. The IR image processing and detection of transition location were performed using an in-house MATLAB package (AWEP Department, Delft University of Technology). The procedure is depicted in figure 3.6 (also see Rius Vidales et al., 2018). Firstly, a calibration sheet was designed and installed on the wing. The wing model was then set at the investigated angle of attack and the calibration sheet was imaged with the IR camera. One such calibration image corresponding to the global IR-FOV is depicted in figure 3.6a. The calibration sheet was then removed and the wing-surface cleaned, before acquiring the necessary IR fields, an example of which is shown in figure 3.6b. To plot these IR images in the xz plane, dewarping/deskewing of the images is applied using the predefined calibration images, thus shedding the wing surface curvature and parallax (viewing angle) errors. The result of applying this calibration procedure over the captured IR field shown in figure 3.6b is presented in figure 3.6c.

The difference in surface temperature of the swept wing between the regions with laminar and turbulent flow can be seen in figure 3.6c (bright and dark regions respectively). This difference in temperature is exploited in this thesis and the transition location is identified as the average location of maximum temperature gradient along the streamwise direction (see blue points in 3.6d). To increase the robustness of this identification scheme, outliers are removed



Figure 3.6.: (a) Captured calibration image corresponding to the *global IR-FOV*. (b) Raw image of the thermal signature of the boundary layer over the swept-wing model's surface. (c) Full IR field, after the application of dewarping/deskewing. (d) Detected transition front in the region enclosed by the dashed-white rectangle in (c).

based on the spanwise standard deviation of the detected positions. The typical jagged pattern of the transition front necessitates the use of a linear fit through the transition locations along the span which is represented by the solid red line in 3.6d. The standard error of this linear fit is used as an estimate of the transition front width (dashed red line in 3.6d). Wherever reported, the average streamwise location of this linear fit is used as an approximation to the spanwise-averaged transition front location. It is noteworthy however that, since quantification of pertinent features such as transition location and CFI wavelength is only dependent on thermal gradients, the results presented in this thesis are the raw output of the IR camera and no quantitative temperature data was used.

3.2.2. Flow velocity measurement through particle image velocimetry

Particle image velocimetry (PIV) is one of the most successful techniques to obtain quantitative data of the flow velocity, that has emerged in the past decade. It is currently the most commonly used technique for velocity measurements. The primary reason for this is that it is built upon optical observations of the flow and thus is considered to be non-intrusive. Additionally, it can provide spatially and temporally correlated velocity data along two or three spatial dimensions. An excellent overview of this measurement technique can be found in the book by Raffel et al. (2007).

The principle of PIV is based on measuring the displacement of small tracer/seeder particles that are illuminated with a pulsating light source. The illumination is provided by laser light sources. The use of lasers for this purpose is motivated by the fact that they emit high energy collimated light that can be shaped as necessitated by the measured flow fields through a combination of spherical and cylindrical lenses. The tracers used should be small enough so that they passively displace with the flowing fluid without interfering or altering the flow characteristics. In low Reynolds number flows droplets of water-glycol mixture or dielectric oil can be used as tracers. While in high Reynolds number flows where compressibility effects are prominent, solid tracer particles such as titanium dioxide powder is commonly used. The light scattered by these particles are imaged using digital cameras and used to compute velocity vector fields.

In planar PIV measurement systems, which provides two-dimensional velocity data, the light sheet is shaped into a thin sheet by means of the optical lenses. Digital cameras are installed such that they are normal to and focused on the illuminated plane. The tracer particles in the flow manifest themselves as high-intensity pixels on the images recorded by the cameras. For each representative time instant, two images are captured, separated by a pre-defined time separation that is significantly lower than the typical time-scales in the flow configuration being investigated. The captured images are divided into small pieces called interrogation windows. The average displacement of particles contained within these interrogation windows in consecutive frames is estimated by applying a normalised cross-correlation operator over the light intensity in the chosen interrogation windows according to equation,

$$C(\Delta x, \Delta y) = \sum_{1}^{l} \sum_{1}^{m} \frac{I_1(x, y) - \bar{I}_1}{\sigma_{I_1(x, y)}} \cdot \frac{I_2(x + \Delta x, y + \Delta y) - \bar{I}_2}{\sigma_{I_1(x + \Delta x, y + \Delta y)}}$$
(3.1)

Here, I_1 and I_2 are the light intensity distribution in a particular interrogation window of size $l \times m$ px² in subsequent images. The intensity I_2 is used as a sliding template for estimating the normalised cross-correlation with I_1 . \overline{I} and σ_I are the average and standard deviation of the light intensity in the windows and are used for normalisation. The value of Δx and Δy , for which the cross-correlation operator C is maximum defines the average particle displacement within this interrogation window between the two images. Dividing these values with the time separation between the images provides the corresponding components of velocity. This process is repeated in the entire captured image to obtain the two-component velocity vector field.

Image acquisition can either be made in a low-speed or time-resolved manner. Although lowspeed PIV lacks temporal resolution, the spatial resolution of the vector field is significantly higher compared to the time-resolved measurement, for the same FOV size. Additionally the PIV acquisitions can also be classified based on the manner in which consecutive images are captured. In *single-frame* mode, images are captured continuously at a particular sampling rate (f_s) . As such, consecutive images are correlated with one another in order to compute the vector field. This type of acquisition is usually chosen for low-speed flows that necessitates high sampling frequencies and always provides a time-resolved velocity data. Instead in *double-frame* mode acquisition, two images separated by a known lag are captured for every representative time instant and used to compute the vector field. The sampling frequency is then the reciprocal of the time-separation between adjacent image-pairs.

This work-flow can be scaled up to acquire three-dimensional velocity vector field data through the tomographic PIV technique proposed by Elsinga et al. (2006). Some more recent developments in this regarded are reported by Scarano (2012). This technique capitalises on the volumetric reconstruction of particle illumination captured through a number of conventional PIV cameras (at least three) simultaneously. Thus, the first necessity is that the illumination must



Figure 3.7.: Schematic of the experimental set-up to characterize the wall-jet induced by the DBD plasma actuator with planar PIV. DBD plasma actuator (black: dielectric, orange: electrodes, grey: grounded-electrode encapsulation), laser sheet (green), FOV (red-dashed rectangle) and coordinate system $(x_p y_p)$ are shown.

be volumetric instead of a plane. This can be achieved by expanding the laser sheet through spherical lenses. The cameras are arranged to capture the volumetric FOV from different angles. Through the use of Scheimflug adapters, the lenses of all the employed cameras are focused on the same plane. However, spatial calibration of this measurement technique is not as simple as the planar PIV set-up and as such, is done in two steps. The first iteration of calibration is done by introducing a physical multi-plane calibration plate, the relative location of whose surface markings are known in all three spatial dimensions. The second calibration iteration is referred to as *self-calibration* (Wieneke, 2008) and involves introducing tracer particles. Their positions are determined through optical triangulation and used to correct the outcome of the first calibration iteration.

Once the calibration procedure is completed, the volumetric reconstruction of the flow field being measured is carried out from the two-dimensional images captured by the different cameras. This essentially results in locating the tracer particles in space through the scattered light intensity in different *voxels* (the three-dimensional counterpart of a pixel). This three-dimensional definition of particle locations is achieved by means of some Multiplicative Algebraic Reconstruction Techniques (MART). Following this, the computation of the flow velocity vectors is carried out, the procedure of which is similar to that in planar PIV. The volumetric reconstruction is divided into three-dimensional interrogation boxes with certain number of voxels in each dimension. Cross-correlation of these interrogation boxes in subsequent volumes (separated in time) is computed to arrive at the velocity vector components. The cross-correlation computation is similar to that shown in equation 3.1, Although, this equation is modified to account for the displacements in the third dimension as well. The different PIV measurement systems employed in this thesis are now introduced which is followed by a brief description of the method followed to estimate the measurement error.

3.2.2.1. High-resolution time-averaged PIV to measure DBD actuator's induced flow field

In order to measure the average velocity field of the ionic wind/induced jet generated by the DBD plasma actuator with a high spatial resolution, a low-speed planar PIV measurement system was utilized. These results were acquired in quiescent conditions to separate the effect of the actuator from other hydrodynamic and pressure phenomena. Thus, the DBD actuator itself was covered with a transparent optical box during this measurement (and in other DBD actuator's induced velocity measurements through PIV). The box was then adequately filled with particles



Figure 3.8.: Tomographic PIV set-up to acquire three-component velocity data of the DBD plasma actuator's induced flow field. The x_p direction is represented by the red arrow.

of dielectric oil (Shell Ondina) which served as the seeder particles.

An Evergreen Nd-YAG multi-pulse laser illuminated the FOV. A combination of spherical and cylindrical lenses was used to convert the laser beam into a 1 mm thick light sheet, which was oriented perpendicular to the dielectric surface to capture the ionic wind in the $x_p - y_p$ plane (see figure 3.7). The PIV fields were captured using a 16 megapixel Imager LX camera with a 4920 × 3280 px² sensor (with 5.5 μ m pixels). The camera was equipped with a Nikon-Nikkor 200 mm micro lens ($f_{\#} = 8$) which resulted in a final FOV of 30 × 20 mm². Image acquisition, pre-processing and correlation were performed using LaVision Davis 8.4 suite. The acquisition was carried out at a rate of 3.7 Hz in double-frame mode. The time interval between the two laser pulses was fixed at 12 μ s so as to realize a particle displacement of 6 – 10 pixels. Thus, cross-correlation was performed with a final interrogation window of 16 × 16 px² with relative overlap of 75%. As such, the final vector spacing was 25 μ m in both x_p and y_p directions.

3.2.2.2. Three-dimensional measurement of DBD actuator induced flow field with tomographic-PIV

Time-resolved tomographic PIV was employed in the current thesis to acquire three-component velocity data of the DBD plasma actuator's induced flow field. The set-up of this experiment is depicted in figure 3.8. This acquisition was carried out in quiescent conditions as well. The actuator was placed inside an optical box in an upright position such that the induced flow is oriented towards the right edge of the figure (red arrow). Air in the optical box was seeded with droplets of paraffin oil, 1 μ m in average diameter, generated by a *TSI* atomizer nozzle. Illumination of the particles was provided by a *Continuum Mesa PIV* (18 mJ per pulse) Nd:YAG laser. The laser light was introduced from outside the optical box. The light beam was shaped into a 4 mm thick sheet by using light sheet optics (combination of cylindrical and spherical lenses) and was oriented parallel to the dielectric surface.

The imaging system comprised of four *Photron FastCAM SA1* digital cameras featuring a $1024 \times 1024 \text{ px}^2$ sensor each (20 μ m pixels and 12 bits of digital resolution). The active sensor of the cameras were reduced to $960 \times 440 \text{ px}^2$. The cameras were equipped with a *Nikon-Nikkor 200* mm micro lens which were operated at a numerical aperture of $f_{\#} = 8$. The imaged volume was

 $(x_p \times y_p \times z_p =)$ 33 × 4 × 16 mm³. A lens shift mechanism was fixed between the cameras and the objectives in order to comply with the Scheimflug condition and adjust the focal plane alignment on the camera sensors. The laser and the four cameras were synchronized using a *Lavision High Speed Controller* and *Lavision Davis 8.2* suite. Image acquisition, volume reconstruction and cross-correlation were carried out through the same software suite. The acquisition rate was $f_s = 10$ kHz in single-frame mode, resulting in a pulse separation of 100 μ s between consecutive frames.

A dual layer target was first used for the calibration of this measurement system and the obtained mapping function was corrected using the volume self-calibration technique (Wieneke, 2008). As the current measurement is time-resolved, the simultaneous motion-tracking-enhanced (SMTE-) MART (Lynch & Scarano, 2015) operator was used for volumetric reconstruction of the acquired fields. Cross-correlation to obtain the velocity vector fields was performed with a final interrogation volume size of $48 \times 32 \times 24$ voxel³, with 75% overlap. This resulted in a final vector resolution of 0.22 mm in all three directions.

3.2.2.3. Time-resolved PIV to measure DBD actuator's induced flow field

Time-resolved (high-speed) planar PIV was also employed in some instances for the quantification of the wall-jet induced by the employed DBD plasma actuators. The actuator was placed in a transparent optical box which was seeded with droplets of olive oil generated by a *TSI* atomizer. A *Continuum Mesa PIV* (18 mJ per pulse) Nd:YAG laser was used to illuminate the necessary FOV. The laser beam was shaped into a 1 mm thick light sheet and was oriented perpendicular to the dielectric substrate to capture the $x_p - y_p$ plane as in section 3.2.2.1. Image acquisition was carried out using a *Photron FastCAM SA1* camera featuring a 1024 × 1024 px² sensor (20 µm pixels and 12 bits of digital resolution). For this experiment, the active sensor of the camera was reduced to 640 × 368 px². The camera was equipped with a *Nikon Nikkor* 200 mm micro lens operated at $f_{\#} = 8$ which resulted in a final FOV of 16 × 9 mm².

The laser and the two cameras were synchronized using a Lavision High Speed Controller and Lavision Davis 8.3.1 suite. Image acquisition was carried out at a rate of 19.9 kHz in single-frame mode. Thus, the pulse separation between subsequent frames was 50.3 μ s and the corresponding particle displacement between adjacent frames was 6 – 8 pixels. Pre-processing and correlation were performed using the same software package. Cross-correlation was carried out with a final interrogation window of 24 × 24 px² with a relative overlap of 75%, resulting in a final vector spacing of 0.16 mm. It is noteworthy that this time-resolved measurement has a significantly better temporal resolution but at the cost of spatial resolution when compared to the low-speed PIV set-up described in section 3.2.2.1.

3.2.2.4. PIV measurement system for plane mixing layer experiment

In this thesis, time-resolved (high-speed) planar PIV was employed for the quantification of the mixing layer flow field velocity. The experimental set-up is as shown in figure 3.9. The inspected plane is along the x - y directions. The flow was seeded with particles of dielectric oil (*Shell Ondina*) injected upstream of the wind-tunnel convergence. A Nd:YAG high speed *Continuum Mesa PIV* laser (18 mJ per pulse) was employed for particle illumination from downstream as shown in figure 3.9a. The laser head was placed outside the wind tunnel and the light beam was shaped into a 1 mm thick light sheet. The imaging of the necessary FOV was carried out using



Figure 3.9.: (a) Photograph of the PIV system to acquire velocity vector fields of the mixing layer. (b) Schematic of PIV experimental set-up (not to scale). Wind-tunnel est section, splitter plate, PIV cameras, laser-head and laser sheet, free-stream velocities and coordinate system are depicted. The two captured fields: near-wake (dashed red rectangle) and far-wake (dash-dotted blue rectangle) are also presented.

two Photron FastCAM SA-Z cameras featuring a $1024 \times 1024 \text{ px}^2$ CMOS sensor (with 20 μ m pixels and 12 bits of digital resolution). Both cameras were installed outside the wind-tunnel test section at about 0.5 m from the measurement plane which they imaged through the optical window of the wind-tunnel test section. The cameras were equipped with a Nikon Nikkor 50 mm lens operated at $f_{\#} = 2$.

The two cameras were placed adjacent to one another such that their respective FOVs overlap by about 26 mm. A calibration plate was employed to stitch the fields captured by the two cameras using *Lavision Davis 8.4.0* suite which resulted in a final FOV of $318 \times 169 \text{ mm}^2$. The plate was 600 mm × 300 mm in length and height respectively and contained 30×15 black dots (5 mm diameter) separated from each other by 20 mm in both directions. This arrangement (2 cameras and laser) was traversed along the streamwise direction to capture two fields. The first field (near-wake, dashed (red) rectangle in figure 3.9b) spanned from 14 mm upstream of the splitter-plate trailing-edge to about x = 300 mm while the second field (far-wake, dash-dotted (blue) rectangle in figure 3.9b) spanned from about x = 280 mm to x = 598 mm. These two fields overlap each other by about 20 mm and were stitched together using the same calibration plate with an in-house *MATLAB* script. An example of this stitching is shown in figure 3.10. The size of the resulting FOV was $612 \times 169 \text{ mm}^2$.

The laser and the two cameras were synchronized using a Lavision High Speed Controller and



Figure 3.10.: Stitching of near- and far-wake PIV u-velocity (m s⁻¹) fields. The dashed black lines represent the overlapping region along the streamwise direction.

Lavision Davis 8.4.0 suite. The acquisition rate during all the tested cases was 9.5 kHz in singleframe mode. Thus, the pulse separation between subsequent frames was 105 μ s and the particle displacement between adjacent frames was about 6 – 8 pixels. Image acquisition, pre-processing and cross-correlation were performed using the same software package. Cross-correlation was carried out with a final interrogation window of 12×12 px² with a relative overlap of 50%. The final vector spacing is 1 mm in both x and y directions.

3.2.2.5. Measurement error estimation

Correlation error in planar PIV vector computations is estimated to be 0.1 pixels (Raffel et al., 2007). The relative error in the measurement of the displacement can be estimated by comparing the above correlation error with the particle displacement in pixels of the particular acquisition. The relative error on the time separation is very low and can be neglected. Therefore, the relative error of the measured velocity (instantaneous and/or fluctuating fields) is approximated by $[\epsilon_{V(t)} = 0.1/\text{pixel displacement}]$, where V(t) is a time-series of instantaneous velocity data.

The uncertainty on the time-averaged velocity data is estimated using the linear error propagation technique discussed by Sciacchitano & Wieneke (2016). The uncertainty $\epsilon_{\overline{V}}$ of any average velocity quantity \overline{V} is determined from the standard deviation of the corresponding measurement ($\sigma_{V(t)}$) and the number of samples (*ES*) that are uncorrelated in time, as shown in equation 3.2.

$$\epsilon_{\overline{V}} = \frac{\sigma_{V(t)}}{\sqrt{ES}} \tag{3.2}$$

3.2.3. Electrical systems to operate and measurement systems to characterise DBD plasma actuator

The exercise of quantifying the performance of DBD plasma actuators is necessary for a number of reasons. Indeed, this exercise assists in better understanding the working mechanism of the actuators. But more specifically for flow control applications, it helps to ascertain the effectiveness of the applied control and also facilitates the comparison of different control strategies. Furthermore, in order to generate theoretical models to estimate the effect of the applied flow control before-hand, the amplitude of the applied forcing and its spatial distribution is imperative. The different instrumentation employed to operate and to characterize the DBD plasma actuators is introduced here. As already indicated, velocity measurements of the wall-jet induced by the DBD plasma actuators was measured using PIV systems. The set-ups of these measurement systems were already introduced in section 3.2.2 and thus, will not be repeated here once again.

3.2.3.1. Power supplies utilized

The operation of DBD plasma actuators necessitates the use of high-voltage amplifiers to provide the AC input signal. Two different amplifiers are used in this thesis extensively. These devices and the method of their operation is introduced in the here briefly.

3.2.3.1.1. Low-frequency high-voltage amplifier and driving method

The model 30/20A high-voltage amplifier developed by *Trek Inc., USA* (see figure 3.11a) is one of the power supply units widely used in this thesis. This device applies a voltage gain of 3 kV/V over any signal that is fed into it through a function/signal generator. Thus, it can potentially generate high-voltage AC output with any waveform type and with additional amplitude modulation. It has a wide output voltage range of 0 to ± 30 kV amplitude (0.1% distortion) and output current range between 0 to ± 20 mA. It is best suited for generating high-voltage signals in the frequency band of 0 to 2.5 kHz (with only 2% distortion) but higher frequency signals can also be produced. It features a high slew rate of 550 V/ μ s. These features, combined with the low-noise (1.5 V rms) output signal it produces makes it a suitable candidate to the be employed as the power supply to operate DBD plasma actuators.

The input waveform to this power supply was generated with either of the two devices in this thesis. The first is the *Teledyne LeCroy WaveStation 3082* waveform generator. It is a 14 bit waveform generator with an output sampling rate of 500 MS/s. This device has two output channels and can produce signals with a frequency of up to 80 MHz. It can generate signals with all standard waveforms and can additionally apply amplitude or frequency modulation over the output signal. As such, the availability of two channels whose internal clocks are synchronised is especially useful for the current thesis as one of the channels can be used to operate the power supply (and thus the DBD actuator) while the other can be used to provide the trigger signal for iCCD imaging or PIV acquisitions.

The second system employed to supply the amplifier with the necessary input signal was the *National Instruments PXI* controlled with a *LabVIEW* interface. This system was especially used when measuring the thrust generated by the actuators using the laboratory balance (see section 3.2.3.5) or the consumed power (see section 3.2.3.3), as these measurements were also



Figure 3.11.: The power supplies deployed in this thesis to operate the DBD plasma actuators.

acquired through other LabVIEW interfaces.

3.2.3.1.2. High-frequency high-voltage amplifier and driving method

The second power supply unit that is used in this thesis is the *Minipuls* 4 high-voltage amplifier developed by *GBS Electronik, Germany.* This device enables the generation of only sinusoidal voltage signals in the frequency range of 5 - 20 kHz and amplitudes of up to 20 kV. As input, it requires a pulsed signal which defines the frequency of the output signal and a DC voltage which establishes the amplitude of the output signal. The device is as shown in figure 3.11b. It consists of a full-bridge converter (red rectangle in figure) which generates a low-voltage square wave signal from the two inputs. This square-wave signal is then fed into the transformer cascade (blue rectangle in figure) which converts it into a sine-like waveform and amplifies the voltage. As can be seen in the operational range, this power supply is very useful when high frequency, high-voltage AC sine-like output signal is necessitated.

Whenever this power supply was deployed to operate the corresponding DBD plasma actuators, the pulsed signal to fix the AC frequency was supplied to the *Minipuls 4* was supplied from a *National Instruments PXI* module controlled by a *LabView* interface. The DC voltage was supplied through a standard laboratory DC power supply.

3.2.3.2. Electrode thickness measurement with interferometry

Thickness and roughness of the DBD plasma actuator's electrodes can have a significant influence on the flow dynamics if not adhered to certain strict constraints. This is especially the case in the swept-wing boundary layer that will be investigated here, which is highly receptive to surface roughness features (see section 2.2.1). Thus, fabrication of actuators with minimal thickness and roughness is a priority in the current thesis. For a thorough validation of the fabricated actuators, measurement of the electrode thickness is absolutely necessary. These surface measurements were obtained using a white light interferometry profiler.

When there is a difference in the optical path lengths of different rays of white light, the phase difference between them causes interference and results in a pattern where light intensity is high in certain regions (fringes) compared to others. This principle is exploited by interferometers to investigate and quantify small-scale surface features or irregularities. In interferometers, light from a source is split into two beams by a partially reflecting mirror (beam splitter). One of these beams (reference) is oriented towards a reference mirror while the other (measurement beam) is directed towards the sample. The reflected beams from both these surfaces results in an interference pattern that is captured by a CCD image sensor. As the reference mirror and the optical path length of the reflected reference beam is fixed, the interference pattern observed can be used to estimate the optical path of the reflected measurement beam and thus the size of the surface features.

In this thesis, a Talysurf CCI 6000 white light interferometer is employed to investigate the surface features of the DBD plasma actuator's electrodes. This machine is equipped with a $1024 \times 1024 \text{ px}^2$ low-noise CCD image sensor which captures the interference pattern and the corresponding software suite to control the machine reconstructs the surface features. The machine was first calibrated with a dual-layer target (a plate with grooves of pre-defined depth). The sample was then placed on the work table and the working distance range was defined. The interferometer automatically traverses the defined measurement area and provides the results pertaining to the surface features of the sample.

3.2.3.3. Electrical characterisation through discharge current and power measurements

Electrical characterisation of DBD plasma actuators involves the measurement of two parameters, the discharge current and the power consumed by (or delivered to) the actuator. In this thesis the discharge current is measured by connecting a non-inductive shunt resistor, with resistance $R_s = 100\Omega$, in series between the encapsulated electrode of the plasma actuator and the ground. The voltage $(V_s(t))$ across this shunt resistor is measured using an oscilloscope. The discharge current is then computed by dividing the measured voltage by the resistance of the shunt resistor $(I = V_s(t)/R_s)$. The oscilloscope employed in this thesis is a *Teledyne LeCroy HDO6054*. This scope has a bandwidth of 500 MHz and can acquire data at a rate of 2.5 gigasamples per second. The resulting temporal resolution of the acquired electrical data ($\Delta t = 400$ picoseconds) is more than adequate to capture the characteristics of the discharge current which have nanosecond time-scales.

The power consumed by the different actuators is estimated through the integral charge measurement technique adapted by Pons et al. (2005) for DBD plasma actuators, described in section 2.1.2. To this effect, a monitor capacitor with capacitance C_M was connected in series between the encapsulated electrode and the ground. Unless otherwise mentioned in the thesis, the monitor capacitor employed in this thesis had a capacitance of $C_M = 47$ nF. The voltage across the monitor capacitor is measured using the same *Teledyne LeCroy HDO6054* oscilloscope. By multiplying the capacitance of the monitor capacitor with the measured voltage, the charge accumulated (Q_M) in the capacitor is computed as shown in equation 2.1. The charge-voltage Lissajous curve (see figure 2.3a for an example) is then plotted and the area within this curve corresponds to the energy dissipated by the plasma discharge. This value is multiplied with the AC frequency to obtain the power consumed per AC cycle (see equation 2.2). At every tested



Figure 3.12.: Schematic of the experimental set-up to characterize electrical parameters of DBD plasma actuators. DBD plasma actuator (black: dielectric, orange: electrodes, grey: grounded-electrode encapsulation), monitor capacitor (C_M) , shunt resistor (R_S) and the oscilloscope to measure the voltage across the capacitor/resistor are represented.

voltage amplitude, a number of AC cycles were used to converge to an average value of power (P_{avg}) . The measurement was in fact carried out through a *LabVIEW* interface.

3.2.3.4. Discharge characterisation through fast imaging of the plasma layer

Discharge characteristics provide more insight into streamer formation, the underlying structure of the ensuing plasma layer and the physical mechanisms that govern the operation of plasma actuators. This was carried out by optical imaging of the ionization process using an iCCD camera, as discussed in section 2.1.1. The plasma discharge of the DBD actuators was captured using a *Princeton Instruments PI-MAX4* iCCD camera featuring a 1024 × 1024 px² sensor. The camera was equipped with a *Nikon Nikkor 200 mm* lens. The DBD actuator was placed in an upright position and the camera was placed in front of it so as to capture the discharge in the $x_p - z_p$ plane. The supplied voltage and discharge current signals were also measured simultaneously and used to synchronize the camera's opening gate-widths accurately with the voltage cycle and capture the discharge at specific time periods along the input high-voltage signal.

3.2.3.5. Generated thrust

The time-averaged, wall-parallel thrust (T_x) produced by the employed actuators was measured using a laboratory balance. The balance, a *Denver Instrument SI-603* (0.02 mN precision), was placed in a Faraday box to prevent any electromagnetic interference. A rack was fabricated and placed on the balance which assisted in mounting the DBD actuator in an upright position, such that the induced jet was oriented upwards. Thus, the reactionary force measured by the balance can be equated to the thrust generated by the actuator along wall-parallel (x_p) direction. Given the sensitivity of the employed balance, the wire connections made to supply electrical power to the actuator can distort the measurement. Thus, poles made of non-conductive material were placed on either sides of the Faraday box. The wiring (i.e. active and ground) were attached to these poles. The final connection, from here to the electrodes of the actuator was made with tungsten filaments with a diameter of 25 μ m. For each test case, time-averaged thrust was computed from an ensemble of 1000 samples acquired at a rate of 50 Hz by the balance through a *Lab VIEW* interface. The generated thrust was also computed by applying the momentum-balance technique (see section 2.1.4) to the time-averaged vector field of the DBD actuator's induced flow, measured through PIV (sections 3.2.2.1 and 3.2.2.3). These velocity vector fields were also employed to compute the spatial distribution of the EFD body force wherever necessary.

3.3. ANALYTICAL TOOLS

Once the data acquisition has been completed, further processing of these data sets is necessary to assist in the analysis of the corresponding flow configurations. A number of analytical tools are deployed in this thesis to either use the acquired experimental data to make some forecasts based on theoretical grounds or to help gain more insight into the underlying physical phenomenon. These tools are introduced in this section.

3.3.1. Linear stability analysis

Linear stability theory (LST) has been traditionally used in studies on hydrodynamic instabilities both in free shear layers and boundary layers. It has been a versatile tool to predict the onset of instabilities, the related wavelengths and/or frequencies and spatial growth rates. It pivots around the assumption that a parallel base flow (i.e. flow does not vary in the direction of the disturbances' amplification) is unstable to infinitesimal wavy disturbances. The formulation of the LST solver is briefly described below.

The instantaneous flow field can be expanded as $U = \overline{U} + \widetilde{U}$ where $\overline{U}(=[\overline{u}, \overline{v}, \overline{w}])$ is the base flow that satisfies the steady incompressible Navier-Stokes equation and $\widetilde{U}(=[\widetilde{u}, \widetilde{v}, \widetilde{w}])$ represents the perturbations or velocity fluctuations. If the base flow is assumed to be parallel and the perturbations are considered to be infinitesimal and wave-like, these disturbances can be represented by the stream function,

$$\widetilde{U}(x, y, z, t) = \psi(y)e^{i(\alpha x + \beta z - \omega t)}$$
(3.3)

where $\psi(y)$ is the perturbation eigenfunction, α and β are the complex streamwise and spanwise wavenumbers respectively, and ω is the angular frequency. As the instabilities grow along the streamwise direction, evolution of the streamwise wavenumber α (= $\alpha_r - i\alpha_i$) is of particular interest, as $-\alpha_i$ represents the growth rate of the perturbations corresponding to the different instability modes. The spanwise wavenumber (β) is complex as well. However, the spanwiseinvariance of the flow field results in $\beta_i = 0$. The angular frequency (ω) is real as the spatial formulation of LST is being applied here as already mentioned.

The development of the flow and hence the perturbations contained within are governed by the incompressible Navier-Stokes equations (see Schlichting & Gersten, 2016) for derivation). The expansion of the instantaneous flow is applied on the Navier-Stokes equation and all the terms with only the base flow is neglected. Furthermore, going along with the assumption of infinitesimal perturbations, all quadratic terms of these fluctuating components are neglected as well. This results in the linearised Navier-Stokes equations which govern the development of
the perturbations and is given in equation 3.4.

$$\nabla \cdot \widetilde{U} = 0 \tag{3.4a}$$

$$\frac{\partial \widetilde{U}}{\partial t} + (\overline{U} \cdot \nabla) \widetilde{U} + (\widetilde{U} \cdot \nabla) \overline{U} = -\frac{\Delta \widetilde{p}}{\rho} + \nu \nabla^2 \widetilde{U}$$
(3.4b)

Here, ρ is the fluid density and ν is the kinematic viscosity. Equation 3.4a and 3.4b represent the conservation of mass and momentum respectively. Substituting the stream function of the perturbation given in equation 3.3 into equation 3.4 and neglecting \tilde{u} , \tilde{w} and \tilde{p} results in the Orr-Sommerfeld equation.

$$\left[\left(\frac{\partial^2}{\partial y^2} - \alpha^2 - \beta^2 \right)^2 - iRe(\alpha \overline{u} + \beta \overline{w} - \omega) \left(\frac{\partial^2}{\partial y^2} - \alpha^2 - \beta^2 \right) + iRe\left(\alpha \frac{\partial^2 \overline{u}}{\partial y^2} + \beta \frac{\partial^2 \overline{w}}{\partial y^2} \right) \right] \psi(y) = 0$$
(3.5)

The eigenfunction $\psi(y)$ with the defined boundary conditions is then mapped on to the Chebyshev domain and correspondingly, the second and fourth order derivatives are expressed as Chebyshev derivative matrices. These computations are carried out with the algorithm (and the *MatLab* suite) provided by Weideman & Reddy (2000). As previously mentioned, the perturbations pertaining to the different instabilities grow along the streamwise (x) direction only and hence both β and ω are provided as input to the equation. As a consequence, equation 3.5 can be re-arranged as an eigenvalue problem with only the streamwise wavenumber α . This problem is then solved according to the companion matrix technique proposed by Bridges & Morris (1984). The amplification N-factor (Van Ingen, 1956, 2008) of each instability mode with a given frequency and spanwise wavelength λ_z (= $2\pi/\beta_r$) is evaluated through the integral in equation 3.6,

$$N(x,\lambda_z,\omega) = \int_{x_0(\lambda_z,\omega)}^x -\alpha_i(x,\lambda_z,\omega)dx$$
(3.6)

where, $x_0(\lambda_z, \omega)$ is the first streamwise station where the instability mode being considered is unstable.

In both the investigated flow configurations, this LST formulation is employed to estimate the stability characteristics of the respective flow field. The procedure followed for the two flow configurations are introduced below. Finally, the simplified model proposed by Serpieri, Yadala, et al. (2017) to predict the effect of the EFD body force on the stability characteristics of the swept-wing boundary layer is also introduced.

3.3.1.1. Stability computations of swept-wing boundary layer

The stability of the swept-wing boundary layer is investigated with the numerical solutions of the boundary layer equations and linear stability theory (LST). The two assumptions that are made here are it is a parallel base flow and the amplitude of the disturbances is small. In two-dimensional boundary layers, LST is often used to estimate the early stages of instability growth (see Schlichting & Gersten, 2016). Application of this computational exercise to three-dimensional boundary layers can be carried out if the flow is invariant along the span (Mack, 1984), resulting in what is commonly referred to as 2.5-dimensional flow. Thus, if the flow conditions present an infinite swept-wing flow (as is the case in the current investigation), LST can be employed to predict the stability characteristics of the boundary layer. LST has been extensively used in CF instability studies to predict the onset of the primary modes, the corresponding wavelengths and frequencies, directions and spatial growth rates (see for example Reibert et al., 1996; Malik et al., 1999; Serpieri & Kotsonis, 2016).

The stability characteristics of the boundary layer being investigated was obtained through an in-house MatLab solver (Aerodynamics, TU Delft). Here, the laminar boundary layer is solved along the x direction using the steady incompressible boundary layer equations. The continuity equation is given in equation 3.7a, and the streamwise and spanwise momentum equations are as in equations 3.7b and 3.7c respectively (Schlichting & Gersten, 2016),

$$\frac{\partial u}{\partial x} + \frac{\partial v}{\partial y} = 0 \tag{3.7a}$$

$$u\frac{\partial u}{\partial x} + v\frac{\partial u}{\partial y} = -\frac{1}{\rho}\frac{\partial p}{\partial x} + \nu\frac{\partial^2 u}{\partial y^2}$$
(3.7b)

$$u\frac{\partial u}{\partial x} + v\frac{\partial w}{\partial y} = \nu\frac{\partial^2 w}{\partial y^2}$$
(3.7c)

where ν is the kinematic viscosity. It is worth noting that based on the flow invariance along the z direction, the z-momentum equation is decoupled from the others. This system of equations is solved on a Cartesian grid that is stretched such that, the x coordinate coincides with the wing's curvilinear abscissa. Apart from the no-slip and no-penetration boundary conditions at the wing surface, the boundary layer edge velocity $(u_e(x))$ computed from the experimentally measured C_p (figure 3.3a), is fed into the solver as the third boundary condition.

The velocity data obtained from this laminar boundary layer solution is then utilized to solve the Orr-Sommerfeld (OS) equation (3.5) using the Chebyshev spectral colocation method as described previously. The spatial formulation is chosen here and thus, the small-amplitude disturbances are expressed in the form of a wave as in equation 3.3. Once again, the flowinvariance along the z-direction renders $\beta_i = 0$. Furthermore, as the current investigation pivots around controlling the stationary CF instability, the angular frequency too is null ($\omega = 0$). Thus, solving the OS equation gives the disturbance amplification along the streamwise direction represented by α_i .

3.3.1.2. Simplified numerical model of plasma actuation in swept-wing boundary layer

A simplified numerical model based on LST was proposed by Serpieri, Yadala, et al. (2017) to analyse the influence of the applied time-average body force on the receptivity/stability characteristics of the boundary layer, which is employed in the current investigation.

This numerical model is based on adding the EFD body force generated by the DBD actuator as a source term in the steady, incompressible laminar boundary layer equations (3.7). As this system of equations represent a spanwise-invariant flow, this model pivots around a spanwiseuniform forcing configuration. However, certain assumptions have to be made to conveniently model and include the volumetric EFD body force into the boundary layer solver, which are as follows. Firstly, it is assumed that the external flow has no influence on the EFD body force being applied. Secondly, although DBD actuators apply unsteady forcing (see Benard, Debien, et al., 2013; Benard, Noté, et al., 2017), the forcing is assumed to be time-invariant (or steady). For this assumption to hold, it is imperative to operate the actuator at very high frequencies which will render the applied forcing comparable to a steady effect (Dörr & Kloker, 2016; Serpieri, Yadala, et al., 2017). Finally, the wall-normal component of force is assumed to be negligible $(F_{\mu} \approx 0)$. This is essential, as the parallel-flow consideration, on which the system of equations 3.7 is based upon, will not hold otherwise. This assumption can certainly be made in the context of DBD forcing, as it is well established that the wall-normal component of the generated EFD body force is two orders of magnitude lower than the wall-parallel component F_x as demonstrated by Debien, Benard, David, et al. (2012).

These assumptions facilitate the modification of only the x-momentum equation shown in 3.7b, which becomes

$$u\frac{\partial u}{\partial x} + v\frac{\partial u}{\partial y} = -\frac{1}{\rho}\frac{\partial p}{\partial x} + \nu\frac{\partial^2 u}{\partial y^2} + \frac{F_x}{\rho}$$
(3.8)

The boundary layer equations 3.7a, 3.8 and 3.7c with the boundary conditions described earlier are then integrated with the same in-house solver to obtain the laminar boundary layer solution. The stability characteristics of this modified boundary layer is then obtained through the linear stability solver and analysed.

3.3.1.3. Stability computations of plane mixing layer

Mixing layers can be conceived, using hydrodynamic stability theory, as superposition of instability waves that augment and/or weaken along the free-stream direction and interact with each other. The mixing layer flow field is unstable to the perturbations arising from the inviscid Kelvin-Helmholtz (KH) instability mechanism which leads to an exponential growth of two-dimensional waves downstream of the splitter plate, ultimately resulting in their roll-up into discrete spanwise vortices. The high value of shear in the time-averaged flow gives ground for the use of linear stability analysis to numerically investigate the perturbations arising from the inviscid KH instability mechanism in the mixing layer. A thorough review of such analysis applied on the mixing layer is synthesized by Ho & Huerre (1984).

Michalke (1965a) numerically determined the stability characteristics of a shear layer (hyperbolic tangent) velocity profile by solving the linear Rayleigh equation. By comparing the



Figure 3.13.: LST solver validation with the solution of Michalke (1965a). (a) and (b) show the evolution of the real (α_r) and imaginary (α_i) parts respectively, of the streamwise wavenumber at different angular frequecies (ω).

obtained results with experimental data, it was demonstrated that spatially growing disturbances describe the instabilities in a shear layer better than temporally growing ones. This is also confirmed by Huerre & Monkewitz (1985), who indicated that mixing layers with a velocity ratio R < 1.315 are convectively unstable and the instabilities within develop spatially and not temporally. As the velocity ratio of the mixing layer in this thesis is R = 0.33 (see section 3.1.2), instability modes are expected to grow spatially downstream of the splitter-plate trailing edge and as such, stability characteristics are computed using the spatial linear stability theory (LST) formulation introduced above.

$$U^{\tanh} = 0.5 \cdot [1 + \tanh(y)] \tag{3.9}$$

The basic flow of a mixing layer is well estimated by a hyperbolic-tangent profile shown in equation 3.9. Many stability calculations have been conducted on this profile (see for example Michalke, 1964, 1965a; Monkewitz & Huerre, 1982; Huerre & Monkewitz, 1985; Monkewitz, 1988). In this flow configuration, velocity fluctuations are assumed to be absent at both extremes of the flow field along the y-direction, i.e. $\tilde{U}(-\infty) = \tilde{U}(\infty) = 0$. Therefore, equation 3.5 is solved here with the boundary conditions,

$$\psi(-\infty) = \psi(\infty) = \frac{\partial \psi}{\partial y}(-\infty) = \frac{\partial \psi}{\partial y}(\infty) = 0$$
(3.10)

The LST solver was validated by comparing the computed streamwise wavenumbers at different angular frequencies with those obtained by Michalke (1965a). For this purpose, the hyperbolic-tangent profile was utilized. A total of 400 Chebyshev polynomials were used to find the solution. The real and imaginary parts of the streamwise wavenumber (α_r and α_i respectively) computed using the solver is shown in figure 3.13. As it can be seen, the computed results are very similar to those presented in the work of Michalke (1965a). The average difference in the computed values of α_r is $\varepsilon_{\alpha_r} = 0.0014$ and that of α_i is $\varepsilon_{\alpha_i} = 0.0026$ which is two orders of magnitude lower than the computed data. It is good to note that the error in the eigenvalue computation of the solver is expected to be approximately 0.0016. Thus, the differences observed between the computed results and those of Michalke (1965a) is mainly attributed to the machine precision.

3.3.2. Spectral distribution of the flow through power spectral density

Fourier decomposition is a commonly used tool to analyse stationary stochastic processes in the frequency domain. It is used here as well to study the spectral content within the flow fields, especially the velocity time-series obtained from PIV. However, it is well know that the spectra computed using Fourier transform does not converge as the number of snapshots is increased which might affect the accuracy of the decomposition. To avoid this, the modified periodogram method proposed by Welch (1967) is employed to compute the power spectra. The time-series is broken down into multiple overlapping segments and discreet Fourier transform (DFT) is computed on each of them separately. The magnitude of DFT in each segment is then squared, and averaged to obtain the final power spectra. While this approach improves the accuracy of the spectral decomposition, it does this at the cost of frequency resolution.

In this thesis, whenever specified, the presented power spectra are normalized by equation 3.11,

$$\Phi_N^{q'} \cdot \Delta f = (\Phi^{q'} \cdot \Delta f) \frac{\frac{1}{ES-1} \sum_t (q')^2}{\sum_f (\Phi^{q'} \cdot \Delta f)}$$
(3.11)

where Φ is the computed power spectral density (PSD) of some fluctuating quantity q' and ES is the temporal ensemble size. The summations are performed with respect to the corresponding subscript, time (t) in the numerator and frequency (f) in the denominator. This equation is in agreement with Parseval's theorem.

3.3.3. Organisation of fluctuations through spectral proper orthogonal decomposition

In order to gain more insight into the different instability modes in the plane mixing layer flow field and eventually, the growth of the layer, the spatial organisation of the corresponding fluctuations is sought. For this purpose, proper orthogonal decomposition (POD) was applied on the entire PIV time-series. POD, first proposed by Lumley (1967), is a technique to objectively identify the coherent structures in a turbulent flow. A specific form of POD named spectral proper orthogonal decomposition (SPOD) (Delville et al., 1999; Towne et al., 2018) is utilized in this thesis. SPOD modes describe structures in a statistically stationary flow that are coherent in both space and time. Here, decomposition of the cross-spectral tensor (direct Fourier transform of the correlation tensor obtained from experimental data) is carried out. Thus, this method results in eigenmodes that each oscillate at a single frequency which is advantageous to study the current flow configuration.

The formulation of SPOD can be found in Delville et al. (1999) who also applied it on the plane mixing layer and observed the two known flow organisations in the mixing layer (i.e. two-dimensional spanwise vortices and streamwise aligned vortices) in the first mode. However, they applied SPOD on velocity data acquired from hot-wire measurements while here, it is being applied on spatially and temporally correlated PIV measurements for the first time. In a recent article, Towne et al. (2018) also show the formulation and procedure of SPOD and discuss its relationship to dynamic mode decomposition and resolvent analysis. In fact, a *MatLab* implementation provided by the authors is employed here to perform SPOD on the experimental

data.

A brief description of the method is as follow. Let us consider a stochastic variable $Q(\Omega, t)$ which in the current work is the fluctuations in the time-series acquired using PIV. Here, Ω represents the spatial domain of the measurement while t is the time. As a first step, the measured velocity vector fields are converted from the time domain to the frequency domain using discrete Fourier transform (DFT) as follows:

$$\widetilde{Q}(\Omega, f) = \int_{-\infty}^{\infty} \widetilde{Q}(\Omega, t) \cdot e^{-i2\pi f t} dt$$
(3.12)

The cross-spectral tensor at each frequency is then computed by the equation

$$S(\Omega, \Omega', f) = Q^*(\Omega, f)Q(\Omega, f)$$
(3.13)

where * represents the complex conjugate. The integral eigenvalue problem to be solved to decompose the flow field is then given by equation 3.14a

$$\int_{\Omega} S(\Omega, \Omega', f) W(\Omega') \Psi(\Omega', f) d\Omega = \Lambda(f) \Psi(\Omega, f)$$
(3.14a)

$$\Theta(f) = \int \widetilde{Q}(\Omega, f) \Psi^*(\Omega, f) \sqrt{\Lambda(f)} d\Omega \qquad (3.14b)$$

where, $\Psi(\Omega, f)$ is the frequency-dependent eigenfunction and $\Lambda(f)$ is the corresponding eigenvalue which is representative of the integrated turbulent kinetic energy (TKE) of fluctuations of that particular mode. The weight $W(\Omega)$ is a positive-definite Hermitian. The corresponding eigenvectors (Θ) are computed by equation 3.14b.

However, as mentioned in section 3.3.2, the DFT do not converge as the number of snapshots is increased. To obtain a converged spectra, the Welch (1967) method is utilized here as well. The time-series is broken down into a number of overlapping sub-sets or blocks each with equal number of frames (N_{fr}) and the DFT is computed on each of these blocks. The Fourier components of a particular frequency in all the blocks is stacked together and the SPOD is carried out. As a result, one receives a number of SPOD modes equal to the number of blocks at every frequency. Hereafter, the notation $\Psi(f, k)$ and $\Lambda(f, k)$ will be used to refer to the eigenfunction and eigenvalues respectively, of each mode (k) at a particular frequency (f).

The detection of coherent vortices and, the investigation of their dynamics and interactions due to the applied forcing is desired. For this purpose, it is necessary to reconstruct the velocity fields from the SPOD modes. This is done according to equation 3.15.

$$\widetilde{Q}_{R}(\Omega,t) = \sum_{f} \left[\sum_{k} \Theta(f,k) \Psi(f,k) \sqrt{\Lambda(f,k)} \right] \cdot e^{i2\pi ft}$$
(3.15)

It is good to note that due to application of Welch (1967) method, this generates a number of



Figure 3.14.: Vorticity $-\omega_z$ (a) and non-dimensional *Q*-criterion (b) computed on the near-wake PIV field presented in figure 3.10a.

realizations equal to the number of blocks of the SPOD computation. Using the block size and overlap of blocks selected for the SPOD computation, these blocks are stitched together in time to obtain the complete time-series.

3.3.4. Vortex identification through Q-criterion

Vorticity magnitude $|\omega|$ has been widely employed to detect vortex flow structures in different flow configurations. While this approach is fairly successful, it does have its limitations since it does not identify vortex cores and only represents the distribution of vorticity. In the current study, the *Q*-criterion propose by Hunt et al. (1988) (see also Jeong & Hussain, 1995) is employed for vortex identification wherever necessitated.

The variable Q is the second invariant of the velocity gradient. It represents the local equilibrium between the shear strain rate (S) and vorticity magnitude (ω) , computed as shown in equation 3.16a. These quantities are essentially the symmetric and anti-symmetric components of the velocity gradient. The value Q is then computed by equation 3.16b.

$$S = \frac{1}{2} \left(\frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \right) \qquad \qquad \Omega = \frac{1}{2} \omega = \frac{1}{2} \left(\frac{\partial u}{\partial y} - \frac{\partial v}{\partial x} \right) \qquad (3.16a)$$

$$Q = \frac{1}{2} \left(\Omega^2 - S^2 \right) \tag{3.16b}$$

Hunt et al. (1988) define a vortex core as the region where the second invariant Q is positive, with an additional condition that the pressure here be lower than ambient conditions. The vorticity $(-\omega_z)$ and non-dimensional Q-criterion computed on the near-wake PIV field in figure 3.10a is shown in figure 3.14. The vorticity distribution in figure 3.14a shows concentration of vorticity in certain regions and indicates the presence of vortices in this field. The location of vortex cores detected by the Q-criterion (figure 3.14b) corroborates with the regions of high vorticity. This supports the choice made in this thesis to use the Q-criterion to detect vortices and study their dynamics and interactions wherever necessary.

4

FABRICATION AND CHARACTERISATION OF DBD PLASMA ACTUATORS

DBD plasma actuators are part of the central theme of this thesis as they provide the required forcing amplitude to realize flow/instability control over the respective flow configurations. In order to achieve the objectives of this thesis which was introduced in section 1.1, a few innovations with respect to DBD plasma actuators is imperative. The first issue that arises is the thickness levels, especially of the air-exposed electrode. If this geometric parameter does not adhere to strict constraints, the introduction of DBD plasma actuators into the flow configuration, can have a detrimental effect and cause the destabilisation of the respective flow. This is especially the case with the swept-wing boundary layer being investigated in this thesis, which is highly receptive to surface roughness. Thus, one of the objectives of this thesis is to formulate a novel fabrication technique to generate plasma actuators with electrodes adhering to strict thickness and roughness constraints.

The second objective related to proposing this novel fabrication technique goes back to a phenomenon introduced in chapter 1. In the context of the manner in which they manifest in the flow field, instabilities can be temporal/frequency-dependent or spatially occurring. Considering the two flow configurations that is studied in this thesis, the stationary CF instability in the swept-wing boundary layer and secondary instabilities in the plane mixing layer are examples of spatial-occurring instabilities. While travelling CF instabilities and secondary instabilities in the swept-wing boundary layer (although not in the scope of the current thesis), and the primary instabilities of *Kelvin-Helmholtz* type in plane mixing layers are examples of temporal/frequency-dependent instabilities. Temporal instabilities can be controlled through the AC signal delivered to the actuator. Spatial instabilities on the other hand necessitate the application of modulated forcing which can be achieved through the geometry of the electrodes. Thus, the second objective is to formulate the fabrication technique such that it is capable of producing electrodes with intricate geometries.

The final objective that is related to this chapter pivots around the actuators designed to apply modulated forcing. A recurring theme in the current thesis is a comparison of the control effectiveness of the classical actuators applying two-dimensional spanwise-uniform forcing and the spanwise-modulated actuators that will have to realize a three-dimensional forcing in order to manipulate instabilities of spatial nature. To provide such a comparative study, proper quantification of these forcing amplitudes, through the characterisation of the performance of the corresponding actuators is of utmost importance. The research efforts, corresponding results and discussion in order to realize these goals are synthesised in this chapter. To start with, the geometric features of the two types of DBD plasma actuators that is employed in this thesis is first introduced in section 4.1. The formulation of the novel fabrication technique is presented in section 4.2. This task will be incomplete if the generated actuators are not validated. Thus, actuators with a classical design, that produce a two-dimensional EFD body force, were fabricated and their performance was quantified. Results pertaining to this is also included in section 4.2. In section 4.3, the three-dimensionality of the developing plasma layer and the flow field induced by spanwise-modulated geometric configuration is demonstrated, followed by a quantification of the performance of this actuator configuration by comparing it with a classical actuator that generates two-dimensional spanwise-uniform forcing.

Design parameters and performance characterisation of the DBD plasma actuators that will be employed further on in the thesis to realize instability/flow control in the swept-wing boundary layer and the plane mixing layer flow configurations are introduced in sections 4.4 and 4.5 respectively. A brief summary and the concluding remarks is given at the end in section 4.6.

4.1. GEOMETRIC CONFIGURATIONS

The two different DBD plasma actuator configurations that are employed extensively in this thesis are introduced here. The schematics of these two actuator designs are as shown in figure 4.1. These include actuators that exert a two-dimensional spanwise-uniform EFD body force and those that apply a three-dimensional spanwise-modulated EFD body force over the respective flow configurations.

The DBD plasma actuators fabricated to exert a spanwise-uniform EFD body force on the respective flow fields is as shown in figure 4.1a. These are composed of uniform strips of conductive metal as electrodes, w_E and w_C being the widths of the air-exposed and covered electrodes respectively. These electrodes are asymmetrically fabricated on either sides of the dielectric substrate. Here, plasma formation occurs at the interface of the two electrodes $(z_{PA} 4.1a)$, resulting in a two-dimensional EFD body force exertion. This is the classical design of DBD plasma actuators that has been extensively studied in the past (see reviews of Moreau, 2007; Corke, Enloe, et al., 2010; Benard & Moreau, 2014), a brief discussion of which is in section 2.1.

With DBD plasma actuators that apply a spanwise-modulated forcing on the flow, the design of the employed DBD actuators are as presented in figure 4.1b. The air-exposed electrode is a uniform metallic strip as before, while the encapsulated electrode is composed of a series of equi-spaced stems. The length and width of each steam is l_S and w_S respectively. These stems are projected from a uniform metallic strip towards the air-exposed electrode. In this geometry, individual stems have a relative overlap o with respect to the air-exposed electrode to ensure consistent plasma formation. Here, the flow ionisation region is expected to correspond only to the electrode-interface regions (i.e. where the electrodes overlap) along the span, thus presumably ensuring a spanwise modulation of the EFD body force. The spacing between adjacent stems defines the spanwise wavelength (λ_z) of the plasma forcing.

4.2. NOVEL TECHNIQUE TO FABRICATE MICRON-SCALE THICK ELECTRODES AND SUBSEQUENT VALIDATION



Figure 4.1.: Design of employed DBD plasma actuators (not for scale). (a) To apply spanwise-uniform EFD forcing. The dielectric substrate, air-exposed electrode (E.E., width w_E), covered or encapsulated electrode (C.E., width w_C) are shown. (b) To apply spanwise-modulated EFD forcing. Width and length of the stems (w_S and l_S respectively), spanwise wavelength (λ_z), and relative overlap between the electrodes (o) are represented. In both schematics, the spanwise extent on which the actuator applies flow control (z_{PA}) and distance between plasma-generating edge of the air-exposed electrode and the edge of the dielectric surface (x_{PA}) are shown as well.

4.2. NOVEL TECHNIQUE TO FABRICATE MICRON-SCALE THICK ELECTRODES AND SUBSEQUENT VALIDATION

The necessary flow control is achieved by DBD plasma actuators in a non-intrusive manner. However, a common predicament for all aerodynamic applications of these actuators is the thickness of the electrodes. If the thickness, especially of the air-exposed electrode is comparable to that of the shear layer being controlled, the mere presence of the actuator can potentially alter the frequency content and instability wavelengths in the shear layer, thus destabilizing the flow and causing early transition to turbulence. The story is not very different in the case of the swept-wing boundary layer. In order to control the primary CF instability and the eventual transition caused by it, flow control has to be applied close to the leading edge (Reibert et al., 1996; Radeztsky et al., 1999; Tempelmann et al., 2012; Serpieri, Yadala, et al., 2017). However, this flow configuration is highly receptive, especially close to the leading edge, to surface roughness (Bippes, 1999; Saric, Reed, et al., 2003; Downs & White, 2013; Rizzo et al., 2019). Thus, if the electrodes of the actuator do not adhere to strict thickness and roughness constraints, they can potentially induce additional stationary CF disturbances into the boundary layer, promoting transition.

The task of devising a novel technique to fabricate DBD plasma actuators with electrodes that adhere to strict thickness (few microns) and roughness (comparable to background surface roughness) constraints, to ensure minimum protuberance due to the actuator was undertaken. A few additional prerequisites were recognized and factored in as well while devising the technique, which are as follows. Many control strategies necessitates the application of spatially modulated forcing (for example, forcing at specific wavelengths for boundary-layer instability control Wassermann & Kloker (2002), Schuele et al. (2013), and Serpieri, Yadala, et al. (2017) or oscillating wall forcing for turbulent drag reduction). Hence fabrication of electrodes with intricate geometric features to exert such three-dimensional forcing must be attainable. The technique should facilitate fabrication of electrodes on a wide range of dielectric substrates. Thus, it must abstain from any chemical or thermal treatment of the dielectric sufface to preserve the chemical composition and dielectric properties of the material. The produced plasma actuators must have good performance characteristics, repeatability and lifetime. Lastly, the



Figure 4.2.: Different steps carried out to fabricate DBD plasma actuators.

technique must be cost effective and simple to be realized in typical laboratory conditions.

4.2.1. Fabrication procedure

With this as the motivation, a fabrication technique is envisioned which involves spraying of sub-micron conductive silver particles in the negative space left by a physical mask that bears the geometric features of the necessary electrode. To ensure repeatability and an even deposition of metal in the entire electrode, the whole process was carried out using a computer numerical control (CNC)-XY plotter *Tooli*. The different steps involved in this fabrication process are shown in figure 4.2 and explained here briefly.

Step 1: Electrode design

A vectorized mechanical drawing of the necessary electrode is generated using any commercially available computer-aided drawing and drafting (CADD) suite. This task is vital to articulate the required geometric features of the electrodes and to maintain close tolerances. These vectorized drawings are used to generate necessary sequences of machine control commands to govern the motion of different tool heads of the CNC-XY plotter.

Step 2: Generation of mask

The dielectric surface is first prepared for fabrication. A black, non-reflective, self-adhesive sheet $(DC \ Fix \ Decorative \ Foil \ 346-8091)$ is fixed on the dielectric substrate. It is placed in the working area of the CNC-XY plotter such that the electrode is fabricated at the required location. The geometric shape of the electrode is cut out of the self-adhesive sheet with a 2.5 W laser tool head (445 nm, 9 mm Laser diode with a three element glass lens), fixed to the plotter. The laser power and number of passes were chosen such that it cuts only the self-adhesive sheet and has no effect on the dielectric surface. Once the laser-cutting is complete, the self-adhesive with the shape of the electrode is removed from the dielectric substrate, thus generating an opportune mask to be used in $Step \ 3$. Generating masks in the manner briefed is advantageous as it improves the overall accuracy of the fabrication process, reduces chances of misalignment of electrodes and less cumbersome. However, there can be other means to generate the necessary masks. It must be stressed that the accurate placement and alignment of the two electrodes on either sides of the dielectric substrate is a high priority in order to avoid discontinuity in the electric field and thus guarantee a homogeneous EFD body force.

Step 3: Computer-controlled spraying

Once the opportune masks are generated, a tool head with a spraying nozzle is attached to the CNC-XY plotter. A conductive silver coating (3830, Holland Shielding Systems BV) is diluted with an industrial solvent to produce the metallic-ink which is fed into the spraying nozzle. The employed coating is composed of micron-sized silver particles (25-50% of volume) mixed with a thermoplastic binding resin (4-Methylpentan-2-one, less than 60% of volume). In the current work, one part of the coating is diluted with two parts of the industrial solvent, Methyl-Ethyl-Ketone (MEK). This metallic-ink is then sprayed in the negative space generated by the mask. When sprayed, the solvent evaporates, depositing the silver particles on the dielectric substrate. The coating is left to dry for a few minutes before the mask is removed. The dielectric substrate is immediately cleaned with ethyl-alcohol to remove any stray silver particles and patches of the adhesive from the mask. The thickness of the electrodes is controlled by the ink-flow rate and feed rate of the spraying-nozzle tool head, and number of passes performed.

4.2.2. Validation of geometric and electrical aspects of fabricated actuators

In order to validate the performance of the DBD plasma actuators fabricated through the novel technique, three actuators with a classical design that generate a spanwise-uniform EFD body force were constructed. These actuators were constructed on 3 mm thick PMMA plates which served as the dielectric substrate. The air-exposed electrodes of two of these actuators were fabricated using the technique briefed in section 4.2.1. These electrodes were $w_C = 10$ mm wide, 140 mm long straight electrodes made of conductive silver. To achieve different air-exposed electrode thickness, the number of passes carried out with the spraying nozzle was one and five respectively, for these two actuators. Both actuators had 40 μ m thick, 20 mm wide self-adhesive conductive copper tape as the grounded electrode. The two electrodes were asymmetrically mounted on either sides of the dielectric substrate with no relative overlap or gap between them.

The third actuator with 40 μ m thick, self-adhesive conductive copper tape as the air-exposed electrode (henceforth referred to as Cu_{40 μ m} actuator) was constructed to be used as a benchmark to validate the performance of the novel actuators. The rest of the dimensions of this actuator are similar to those described earlier. Finally, the grounded electrode of all three actuators were encapsulated with polymide tape.

The actuators were powered using a *Trek 30/20A* high-voltage amplifier (see section 3.2.3.1). For all test cases, the air-exposed electrode was supplied with the high-voltage AC signal. The voltage waveform was sinusoidal, supplied voltage amplitudes (V_{ac}) varying from 12 kV to 20 kV (frequency $f_{ac} = 2$ kHz). Then, the voltage amplitude was maintained at $V_{ac} = 16$ kV and AC frequencies $f_{ac} = 500$ Hz, 1 kHz and 4 kHz were tested.

4.2.2.1. Microscopic view of the fabricated electrodes

Surface features of the silver electrodes fabricated with the novel technique was studied using interferometry. As introduced in section 3.2.3.2, a *Talysurf CCI 6000* white light interferometer was employed for this purpose. A FOV of $3 \times 3 \text{ mm}^2$ in the $x_p - z_p$ plane centred about the edge of the air-exposed electrodes was captured with this device for both actuators with conductive-silver electrodes. The acquired surface profiles had a final spatial resolution of 1.8 μ m in both



Figure 4.3.: Surface features of the electrodes fabricated with the novel technique obtained through interferometry. Coloured contour levels (in μ m) in (a) and (b) represent height of the top-most surface from the dielectric substrate of the electrodes produced with one and five passes of the spraying nozzle. Surface features of the same in three-dimensions (colour scales in μ m) are presented in (c) and (d).

 x_p and z_p directions. These results are presented in figure 4.3, with 4.3a and c corresponding to the electrode fabricated with one pass of the spraying nozzle and 4.3b and d corresponding to the electrode fabricated with five passes of the spraying nozzle. It is worth noting that these actuators were operated before obtaining these surface features in order to confirm the generation of a plasma discharge when supplied with a high-voltage AC signal.

The colour scale in the contour maps in figure 4.3a and b represent the height of the top-most surface at each location in the measured FOV from the dielectric surface. While the fabricated electrodes seem continuous and smooth to the human-eye, these contour plots indicate that they posses irregularities although in micrometric scales. This becomes clearer in the three-



Figure 4.4.: (a) Microscopic view of the fabricated electrode's edge (1 pass of spraying nozzle). (b) Surface profile at the electrode edge obtained from interferometry.

dimensional surface plots presented in figures 4.3c and d. A downside of employing white-light interferometry for the current purposes must be highlighted before moving forward. Silver is a reflective material, especially with white light. As a consequence, a number of high-amplitude spikes are observed in both the surface plots. Thus, these spikes are not expected to be representative of any particular geometric feature but rather due to high-intensity light from the reflective electrode surface.

The first geometric feature that can be discerned from the surface plots in figures 4.3c and d is the difference of electrode thicknesses achieved by increasing the number of passes performed with the spraying nozzle. Secondly, the edge of the produced electrodes is not a sharp step as with the copper tape electrodes, but the thickness increases gradually and attains its maximum value. This feature is more evident in figure 4.3d where the electrode thickness increases until $x_p \approx -1$ mm before it achieves a quasi-constant level. After this level is achieved also, the top-most surface is not smooth but contains a number of micrometric irregularities. In this region, i.e. $x_p < -1.5$ mm along z_p , the roughness standard deviation is $\approx 2.5 \ \mu$ m for the electrode fabricated with one pass of the spraying nozzle and $\approx 4.5 \ \mu$ m for that fabricated with five passes. Indeed the actual value is expected to be lower as the spurious spiked due to light reflection increases the margin of error of these estimations.

Finally, the fabricated electrode's edge is uneven and not very perfect. This can also be observed in the microscopic view of the edge of the electrode fabricated with one pass of the spraying nozzle shown in figure 4.4a. The generated edge is not continuous as will be the case when a conductive metallic tape is used to fabricate the electrode, but has a number of microscopic undulations. Finally, indication of miniscule deposition of silver particles in the vicinity of the electrode's edge that are not physically connected to the electrode itself is also observed in the interferometry fields in figure 4.3. These minute depositions are also seen in the microscopic view in figure 4.4a. While operating the actuator might have eroded these out-ofplace depositions, they are still present in the vicinity of the electrode. Whether these depositions impacts the working of these actuators will be established through further characterisation that is to follow.

In order to estimate the thickness of the produced electrodes, profiles of the surface geometry at the fabricated electrodes' edge measured with interferometry, for both the actuators fabricated for this validation is shown in figure 4.4b. This figure reiterates the earlier observation that the



Figure 4.5.: Electrical power consumption versus voltage amplitude (a) and frequency (b) for the three DBD actuators.

the thickness of the electrodes increases gradually and attains a quasi-constant value at $x_p \approx -1$ mm. The thickness of the electrodes thus fabricated is found to be 6 μ m (one pass of spraying nozzle) and 30 μ m (five passes of spraying nozzle). Henceforth, these actuators will be referred to as Ag_{6µm} and Ag_{30µm} actuators respectively. Finally, the profiles presented here do not posses any spurious spikes due to light reflection. From these profiles, the roughness standard deviation of both electrodes in the region $x_p < -1$ mm is estimated to be $\approx 1.3 \ \mu$ m.

4.2.2.2. Power consumption

The electric power consumed by the three DBD actuators was estimated through integrated charge measurement technique described in section 3.2.3.3. To this effect, a monitor capacitor with $C_M = 33$ nF was connected between the encapsulated electrode and ground. For every tested case, 20 cycles were used to converge to a mean value of power. The power consumed by the two actuators with silver electrodes as a function of supplied AC voltage amplitude is presented in figure 4.5a. The onset voltage where the plasma discharge begins is $V_0 = 6$ kV for both actuators. The power consumed is observed to increase with according to $P_{avg} \propto$ $(V_{ac} - V_0)^2$ with the AC voltage amplitude as observed in previous studies (Pons et al., 2005; Forte et al., 2007). While with input AC frequency, it increases linearly as seen in figure 4.5b. This too is in line with previous research findings (see figure 2.3b, reproduced from Forte et al., 2007). The power consumed by both actuators with silver electrodes are very similar which suggests that electrode thickness does not influence power consumption. The power consumption of the $Cu_{40\mu m}$ actuator is also presented in figure 4.5 for comparison. The actuators with silver exposed electrode are observed to consume more power compared to the copper. However, the difference is very small (about 7 W m⁻¹ on average). Thus, even though the edge of the fabricated electrodes is not perfect and have a number of microscopic undulation, the performance of these actuators in terms of power consumption is not significantly affected.

4.2.3. Influence of air-exposed electrode thickness on the induced flow velocity

The measurement the fabricated electrodes' thickness with interferometry that was presented in section 4.2.2.1 indicated that the thickness of the two silver electrodes was different from one



Figure 4.6.: Mean-convergence of the u_p -velocity component at different x_p -locations ($y_p = 0.5 \text{ mm}$) in the DBD actuators' induced flow fields. $V_{ac} = 16 \text{ kV}$, $f_{ac} = 2 \text{ kHz}$.

another. The electrode fabricated with just one pass of the spraying nozzle was about 6 μ m thick (Ag_{6µm} actuator). While that fabricated with five passes was approximately 30 μ m thick (Ag_{30µm} actuator). Although the electrode edges are not perfect as already mentioned, these actuators provide a favourable circumstance to investigate the influence of air-exposed electrode thickness on the velocity within the induced flow field.

This exercise necessitates velocity field measurements with a high spatial resolution to capture intricate differences of velocity distribution within the induced flow. Therefore, velocity measurements were obtained through the low-speed planar PIV set-up described in section 3.2.2.1. Initially, a convergence test was carried out and it was found that an ensemble average of about 50 vector fields delivered a converged mean field (see figure 4.6). Thus, 300 PIV fields were acquired for each tested case to ensure full statistical convergence. Furthermore, Gaussian smoothing (window: 25×25 , $\sigma = 0.5$) was applied on the instantaneous fields before computing the mean. The obtained time-averaged velocity vector fields for both the actuators with silver air-exposed electrode are presented in figure 4.7.

The actuators produce a thin wall-jet approximately 0.75 mm thick (location of $0.9 \cdot u_{p,max}$ from the dielectric wall). A weak suction effect is detected upstream of the plasma-generating edge of the air-exposed electrode ($x_p = 0$). The u_p -velocity profiles in the ionic wind, perpendicular to the dielectric surface are extracted for both actuators at four different x_p -locations and presented in figure 4.7. The Ag_{6µm} actuator is observed to produce a thinner jet comparatively. The location of the maximum u-velocity ($u_{p,max}$) is also closer to the wall for this actuator. However, at all four x_p -locations, $u_{p,max}$ is lower for the Ag_{6µm} actuator. This is attributed to the effect of skin friction which is more prominent closer to the wall, thus reducing the magnitude of velocity.

To elucidate this further, the evolution of $u_{p,max}$ in the ionic wind along x_p for both actuators is presented in figure 4.8a (blue and red curves in figure 4.7a and b). The $u_{p,max}$ is seen to move towards the wall after the air-exposed electrode edge, until $x \approx 2$ mm. Then it starts to deviate away from the wall. It is clearly visible that the $u_{p,max}$ is closer to the wall in the ionic wind generated by the Ag_{6µm} actuator compared the Ag_{30µm} actuator in the entire mean field. For the Ag_{6µm} actuator, at $x_p = 2$ mm, the $u_{p,max} = 3.25$ m s⁻¹ and at $y_p \approx 0.13$ mm. At this location, it is lower than the $u_{p,max}$ of the Ag_{30µm} actuator by about 125 µm.

The change in the location of $u_{p,max}$ with air-exposed electrode thickness, at $x_p = 2$ mm for the five tested voltage amplitudes are presented in figure 4.8b. In all the tested cases, a similar



Figure 4.7.: Time-averaged velocity magnitude fields (in m s⁻¹). $V_{ac} = 16$ kV, $f_{ac} = 2$ kHz. The two grey bars below $y_p = 0$ represents the electrodes of the DBD actuator. The white solid (Ag_{6µm}) and dash-dotted (Ag_{30µm}) curves represent $u_{p,max}$ at each x_p location. The u_p -velocity profiles perpendicular to the wall at four different x_p locations are also presented.



Figure 4.8.: (a) Evolution of maximum u_p -velocity along x_p direction (from figure 4.7). (b) y_p -location of maximum u_p -velocity at $x_p = 2$ mm (dashed black line in (a)) versus voltage.

trend is observed. This confirms that the thickness of the air-exposed electrode has an influence on the topology of the produced ionic wind. It is expectable that a thinner electrode will generate a stronger local electric field close to the electrode edge. This results in the generation and acceleration of charged species closer to the wall. Thus, the thinner electrode is observed to impart more momentum closer to the wall.

It is well known that the maximum velocity in the ionic wind increases with the supplied voltage (Forte et al., 2007). The x_p -location of the maximum velocity also increases linearly with the applied voltage provided the discharge is not filamentary (Jolibois & Moreau, 2009).



Figure 4.9.: Thrust (T_x) generated versus voltage amplitude (a) and frequency (b) for the three DBD actuators.

The results acquired here show that the voltage amplitude has an effect on the y_p -location of $u_{p,max}$ as well. As observed in figure 4.8b, the location of $u_{p,max}$ moves closer to the wall with increasing voltage amplitude. For both actuators, $u_{p,max}$ does not move closer to the wall any further after $V_{ac} = 18$ kV. This is expected to be due to the onset of filamentary discharge.

The thrust generated by the DBD actuators is computed from the mean PIV vector fields using the momentum balance approach (see section 2.1.4). The chosen control volume is the entire mean vector field shown in figure 4.7 so as to minimize pressure effects. The thrust generated (T_x) as a function of applied voltage amplitude and frequency is presented in figure 4.9. The thrust generated by the Cu_{40µm} actuator increases with applied voltage as reported previously (Kotsonis et al., 2011; Debien, Benard, & Moreau, 2012). For the actuators with silver air-exposed electrode, the thrust increase is quasi-linear in the tested voltage range until $V_{ac} = 18$ kV. However, the curve slightly decreases at $V_{ac} = 20$ kV suggesting the onset of filamentary discharge. This confirms with the results of $u_{p,max}$ observed previously. The thrust is seen to increase linearly with frequency as observed by Kotsonis et al. (2011). The curve shows a small reduction for Ag_{6µm} actuator after $f_{ac} = 2$ kHz which again is attributed to the discharge becoming filamentary.

4.3. PERFORMANCE CHARACTERISATION OF SPANWISE-MODULATED DBD ACTUATORS

In this thesis, a comparison is made between the effect of spanwise-uniform and spanwisemodulated DBD plasma actuator forcing on the two flow configurations considered. While the influence of the input signal parameters (especially voltage amplitude) on the electrical and mechanical performance is well understood for spanwise-uniform actuators (see section 2.1 and reviews by Moreau, 2007; Corke, Enloe, et al., 2010; Benard & Moreau, 2014), little is known regarding spanwise-modulated devices. Previous research efforts (see for example Dörr & Kloker, 2017; Serpieri, Yadala, et al., 2017; Guo & Kloker, 2020) have considered the performance of the spanwise-modulated actuators to be similar to a corresponding spanwise-uniform actuator. As the applied EFD forcing by these spanwise-modulated devices is expected to be threedimensional, their performance cannot be deemed similar to a corresponding spanwise-uniform actuator without proper investigation. Furthermore, in the current study where a comparison of the two types of forcing is being carried out, the performance of spanwise-modulated actuators is not trivial, but it is absolutely necessary, to better quantify the efficiency of the applied control.

The first necessity in this regard is to demonstrate the three-dimensionality of the ensuing plasma layer and the corresponding induced flow when these spanwise-modulated actuators are operational. This is achieved through fast imaging of the plasma layer using an iCCD camera and three-dimensional velocity field measurements through tomographic PIV respectively. This is followed by a characterisation of the performance of these actuators by comparing it to a spanwise-uniform actuator.

4.3.1. Demonstration of three-dimensionality of the developing plasma layer and induced flow

In order to first demonstrate the three-dimensionality of the ensuing plasma layer and the corresponding induced flow, a spanwise-modulated plasma actuator was fabricated. Kapton polyimide film of 50 μ m thickness served as the dielectric substrate. The electrodes of this actuator were fabricated with the novel computer-controlled spraying technique introduced in section 4.2.1. The top electrode was a $w_E = 5$ mm wide uniform strip of conductive silver. The second electrode was composed of $w_S = 3$ mm wide and $l_S = 10$ mm long stems that were projected toward the air-exposed electrode from another 5 mm wide uniform strip of silver. The two electrodes were fabricated such that a relative overlap of o = 2 mm was maintained between the stems and the air-exposed electrode. The distance between adjacent stems was $\lambda_z = 9$ mm which results in a modulated actuator with spatial duty of $w_S/\lambda_z = 33.3\%$. The air-exposed electrode composed of stems was grounded.

4.3.1.1. Fast imaging of the developing plasma layer

Firstly, the discharge characteristics were investigated by measuring the current signal (see section 3.2.3.3) and capturing the discharge itself with an iCCD camera (see section 3.2.3.4). Discharge during the positive-going and negative-going halves of the AC cycle were captured separately. For each half-cycle, two fields with different opening gate-widths of the iCCD camera were acquired. The first two fields (A and B) were captured with an opening gate-width of 14 μ s while the second two (C and D) were acquired with an opening gate-width setting of 140 μ s. The measured current signal and images of the plasma discharge when the actuator was operated at $V_{ac} = 3$ kV and $f_{ac} = 2$ kHz is presented in figure 4.10.

The discharge current is seen to have two regimes, one during the positive-going half of the input AC signal and the second during the negative-going half. As observed in previous research efforts (see for example figure 2.2, reproduced from Benard & Moreau, 2012), the signal is composed of high-current positive peaks during the positive-going half of the AC cycle. During the negative-going half, while the peaks have a negative polarity as expected, the signal is composed of high-current peaks. This is a consequence of the grounded electrode not being properly encapsulated as demonstrated by Laurentie et al. (2009).

The structure of the discharge captured with the iCCD camera and shown in fields A and B (14 μ s opening gate-width) in figure 4.10 indicates the similarity of the plasma layer generated by this spanwise-modulated actuator to that observed in previous research efforts with spanwise-



Figure 4.10.: Discharge current and fast imaging of the developing plasma layer for the spanwisemodulated DBD actuator configuration. The employed actuator is operated at $V_{ac} = 3$ kV and $f_{ac} = 2kHz$. Dashed-green lines in A, B, C and D represent the stems of the grounded electrode. A & B: 14 μ s opening gate width; C & D: 140 μ s opening gate width.

uniform actuators (Hoskinson et al., 2008; Benard & Moreau, 2012; Debien, Benard, & Moreau, 2012). During the positive-going half of the AC cycle (field A), the plasma layer is composed of streamers (erratic conduction channels) extending from corona spots at the air-exposed electrode edge toward the dielectric substrate. During the negative-going half (field B) only the corona spots are observed, which is followed by a diffusive discharge resulting in a plume shape. This confirms that the high negative current peaks observed during this half-cycle is indeed due to the discharge occurring on the grounded-electrode side of the dielectric substrate and as such will be absent if this was encapsulated.

In both these regimes, the discharge is limited to the spanwise regions between the dashedgreen lines $w_S = 3$ mm apart from each other, which represent the interface of the air-exposed and the stems of the grounded electrode. However, a select few conductive channels are observed just outside these spanwise regions as well. The fields C and D, where the opening width of the iCCD camera is increased to 140 μ s, confirms the occurrence of this "secondary" discharge. This discharge is attributed mainly to two phenomenon. One is a weak electric field generated between the air-exposed electrode's plasma-generating edge and the length-wise edges of the stems (i.e. l_S in figure 4.1b). The second is due to charge deposition on the dielectric surface during subsequent AC cycles. The (average) spanwise length of this secondary discharge, compared to the width of each stem, is approximately $0.4 \cdot w_S$. What effect this will have on the induced flow (a) $V_{ac} = 3$ kV, $f_{ac} = 2$ kHz



Figure 4.11.: Time-averaged velocity magnitude field (m s⁻¹) of the flow induced by the spanwise-modulated plasma actuator captured with tomographic PIV.

fields must be investigated. Nevertheless, a good spanwise-modulation of the plasma layer is achieved with this geometric configuration. With this secondary discharge being accounted for, the spanwise duty of the plasma layer is about 46%. Finally, along the streamwise direction, the discharge spans to about $x_p = 2$ mm.

4.3.1.2. Organisation of the induced flow through tomographic PIV

Having established the spanwise-modulation of the plasma layer developing on the air-exposed side of this actuator, it is now envisaged to investigate whether the induced velocity field shows a similar organisation. This necessitates the use of a measurement system that can provide three-dimensional velocity data. To this effect, the tomographic PIV set-up introduced in section 3.2.2.2 was employed to measure the induced velocity fields. The imaged FOV was such that, it captured the flow organisation along two adjacent stems of the grounded electrode that are $\lambda_z = 9 \text{ mm}$ apart. The time-averaged velocity magnitude fields ($U_p = \sqrt{u_p^2 + v_p^2 + w_p^2}$) of the flow induced by the current spanwise-modulated plasma actuator are presented in figure 4.11. These time-average fields were computed after the starting vortices discussed in section 2.1.3 left the imaged FOV.

The time-averaged induced velocity field when the DBD plasma actuator is operated at $V_{ac} = 3$ kV and $f_{ac} = 2$ kHz is presented in figure 4.11a. This field clearly indicates the three-dimensional nature of the induced flow field. Two induced jets are observed in the imaged FOV, each originating at the interface of the air-exposed electrode and a stem of the grounded electrode. A low velocity is observed just upstream of this interface. This is the weak suction which feeds the fluid into the induced jets. The actuator is seen to accelerate the flow only at the spanwise regions where the electrodes overlap and leads to the formation of wall-jets along the length of the stems. Thus, it results in an induced flow field that is spanwise-modulated according to the distance between adjacent stems, which is $\lambda_z = 9$ mm in this case. This further indicates that the EFD body force generated by these actuators will also be spanwise-wise modulated (three-dimensional). Furthermore, the secondary discharge in the vicinity of the stems that was observed in the iCCD images of the plasma layer in figure 4.10 seems to not have a significant impact on the induced flow field as the generated jets are predominantly limited to the width of the stems. A weak flow is observed on either sides of the stems whose velocity magnitude $(U_p \approx 0.1 \text{ m s}^{-1})$ is only about 20% of the maximum velocity in the jet. While this weak flow can be an artefact of the secondary discharge, fluid entrainment between the shear layers at jets' edges cannot be ruled out either. Thus, even if the secondary discharge in the vicinity of the stems impacts the induced flow field, its effect seems to be rather weak.

The flow induced by the actuator when the frequency of the input AC signal is increased to $f_{ac} = 6$ kHz is shown in 4.11b. This increase in AC frequency will lead to higher power being delivered to the actuator as these parameters vary linearly with one another (see figure 2.3b, reproduced from Forte et al., 2007). The induced flow is still three-dimensional and is modulated according to the spacing between the stems ($\lambda_z = 9$ mm). At this operating condition with higher input power, a second smaller jet core is observed in the induced jet centred around $z_p = 9$ mm, at the stem's edge ($z_p \approx 11$ mm). The same however is not observed in the jet centred around $z_p = 0$. While the reason for this difference between the two jets is not known, it could be attributed to a stronger secondary discharge in the vicinity of the stem at $z_p = 9$ mm. Apart from this, the induced jet is predominantly limited by the width of the stem. The weak induced

flow outside the stem at $z_p = 0$ has the same velocity magnitude ($U_p \approx 0.1 \text{ m s}^{-1}$) as in the previous induced field shown in figure 4.11a. However, when compared to the maximum velocity in this jet, it amounts to less than 10%.

4.3.2. Quantification of electrical and mechanical performance

The discussion in section 4.3.1 clearly demonstrated the three-dimensionality of the developing plasma layer and the corresponding induced flow field generated by the spanwise-modulated DBD plasma actuator configuration. However, for a thorough characterisation of the performance of these actuators, especially to measure the generated force that will be utilized for flow control applications, the use of tomographic PIV measurements as done in section 4.3.1.2 is not very practical. The computational cost to obtain the velocity vector fields, the amount of data generated and FOV limitations renders the use of such techniques as less practical. Considering a large test matrix being desired here in order to thoroughly characterize the spanwise-modulated plasma actuators' performance, the use of tomographic PIV is unfeasible.

While a number of parameters can be measured to ascertain the performance of the actuator as discussed in section 2.1, two parameters, namely the power consumed by the actuator and the corresponding thrust generated, are chosen to quantify the spanwise-modulated actuators' performance. Five different DBD plasma actuators were fabricated to this effect. These actuators were all constructed on 2 mm thick PMMA plates which served as the dielectric. The choice of this dielectric is motivated by the spanwise-uniform actuators constructed on this dielectric being very robust and its performance being well documented in open literature (see for example Jolibois & Moreau, 2009). The electrodes were made with strips of aluminium with a non-reflective black coating on them (total thickness being 85 μ m).

The first actuator was fabricated such to generate a spanwise-uniform plasma discharge (see design schematic in figure 4.1a) and contained two strips of conductive aluminium tape with a non-reflective black coating on them (total thickness being 85 μ m) as the electrodes. The width of the air-exposed and covered electrodes were $w_E = 50$ mm and $w_C = 15$ mm respectively. They were asymmetrically mounted on either sides of the dielectric with a relative overlap of o = 5 mm. This actuator served as the benchmark to compare the performance characteristics of the four other actuators which were fabricated so as to generate a spanwise-modulated plasma discharge. The air-exposed electrodes of these actuators were $w_E = 50$ mm wide strips of aluminium tape, while the covered electrodes were composed of equi-spaced stems as shown in the schematic in figure 4.1b. Each of these stems were $l_S = 13$ mm long and were projected towards the air-exposed electrode from a 2 mm wide aluminium strip such that a relative overlap of o = 5 mm was realized. The distance between the centres of adjacent stems, which defines the wavelength of the spanwise-modulated forcing, was held constant at $\lambda_z = 37$ mm. The difference between the actuators arises through the variation of the stem width, which were $w_S = 30$ mm, 18.5 mm, 9 mm and 3 mm. These result in a spatial duty of $w_S/\lambda_z = 75\%$, 50%, 25% and 8% respectively. Henceforth, these actuators will be referred to by their spatial duty. The covered electrodes of all five actuators were encapsulated with an epoxy resin and were grounded. The air-exposed electrodes were supplied with the high-voltage AC signal from a Trek 30/20A amplifier (see section 3.2.3.1). For all the tested cases, the carrier frequency of the input AC signal was held constant at $f_{ac} = 2$ kHz while the voltage amplitude was varied.



Figure 4.12.: Images of the plasma discharge generated by the spanwise-modulated actuator with 50% spatial duty when operated at $V_{ac} = 10$ kV, $f_{ac} = 2$ kHz.

4.3.2.1. Estimation of discharge length

The results presented here are normalised with the length of the plasma discharge (l_p) . This is straightforward for the spanwise-uniform actuator as the discharge is limited to the interface of the two electrodes and as such, the extent of the discharge is equal to the length of this interface $(z_{PA} \text{ in figure 4.1a})$. The same is not the case with respect to the spanwise-modulated actuators. Ideally, the discharge must be limited to the electrode-interface regions along the span. In which case, the discharge length would amount to the width of the stems multiplied by the number of stems ($l_{CE} = l_p = w_S \times$ number of stems). In practice however, the plasma discharge is not limited here only to the electrode-interface regions, but a secondary discharge occurs in the vicinity of each stem as demonstrated by the iCCD images of the developing plasma layer in figure 4.10 (see also figure 4.12b). Thus, the actual discharge length must be measured for each spanwise-modulated actuator.

In order to do this, photographs of the discharge were captured with a camera. The actuator was placed in an upright position and the camera was positioned such that it imaged the discharge in the $x_p - z_p$ plane as shown in figure 4.12b (similar to the iCCD images in figure 4.10). As a first step, an image of the actuator when it is non-operational was captured (figure 4.12a). Given that the design parameters of the actuators are known, these images of the different DBD actuators were utilized for calibration of the images with a plasma discharge. The discharge length is then estimated by manually selecting the points corresponding to the extent of the plasma discharge around each stem (green '×' and '+' in figure 4.12b). Such length estimations were made for each actuator at different voltage amplitudes and averaged to be used for the normalisation. These estimates are tabulated in table 4.1 for the different spanwise-modulated DBD plasma actuators employed for this exercise

Table 4.1 Spanwise modulated actuator characterisation.				
Actuator	l_p/λ_z	$V_0 (\mathrm{kV})$	$A_P \; ({\rm W} \; {\rm V}^{-2})$	Effectiveness (mN W^{-1})
Spanwise-uniform	-	4	$6.82\cdot10^{-8}$	0.34
75% spatial duty	0.94	4	$7.74 \cdot 10^{-8}$	0.38
50% spatial duty	0.67	4	$5.35\cdot10^{-8}$	0.36
25% spatial duty	0.49	4.5	$3.78\cdot 10^{-8}$	0.4
8% spatial duty	0.33	5	$2.22\cdot 10^{-8}$	0.46

Table 4.1.: Spanwise-modulated actuator characterisation.

4.3.2.2. Power consumption and generated thrust

The power consumed was estimated through the integral charge measurement technique described in section 3.2.3.3. For this a monitor capacitor with capacitance $C_M = 47$ nF was employed. At every tested voltage amplitude, 100 AC cycles were used to converge to an average value of power (P_{avg}) which is presented in figure 4.13a. Considering first, the measurements concerning the spanwise-uniform actuator, the acquired results demonstrate that the variation of power consumption with input AC voltage amplitude is well approximated by a curve $P_{avg} = A_P (V_{ac} - V_0)^2$ (solid black curve in figure 4.13a). The plasma onset voltage $V_0 = 4$ kV and the constant coefficient is $A_P = 6.82 \cdot 10^{-8}$ W V⁻² (fitting done on measured data, before normalisation). It is worth noting that this fitting is only valid after the plasma ignition voltage (V_0). As such, the constant coefficient obtained here is comparable to that reported by Pons et al. (2005).

The power consumed by the spanwise-modulated actuators normalised with the corresponding discharge length (l_p) is also presented in figure 4.13a. It can be seen that the normalised power of these actuators is very similar to that of the spanwise-uniform actuator. Here too, the power consumed increases according to the same power law $P_{avg} = A_P (V_{ac} - V_0)^2$. The plasma onset voltages and the constant coefficients A_P (on measured power values, before normalisation) of these actuators are given in table 4.1. The constant coefficient decreases with the spatial duty which is a consequence of the reducing discharge length.

The time-averaged, wall-parallel thrust (T_x) produced by the spanwise-uniform DBD actuator is presented in figure 4.13b. It is worth noting once again that the measured values are not the total force generated by the DBD actuator, but the EFD force subtracted by frictional forces incurred due to the presence of the dielectric substrate. Similar to the power consumption, even the thrust generated by the spanwise-uniform actuator increases with applied voltage amplitude according to a power law (Thomas, Corke, et al., 2009; Kotsonis et al., 2011). Post the plasma ignition voltage, it is well approximated by the relation $T_x \propto (V_{ac} - V_0)^2$, as observed by Debien, Benard, & Moreau (2012). In the case of the spanwise-modulated actuators, the thrust generated comparable to that of the spanwise-uniform actuator when the measurement is normalised with the discharge length, as observed earlier with the power consumption.

A comment must be made regarding the measurements concerning the 8% spatial-duty actuator. In this case, the power measurements showed the plasma onset voltage to be $V_0 = 5$ kV while the balance measurements show thrust being generated by this actuator only after $V_{ac} = 10$ kV. This is an artefact of the force being measured by the laboratory balance, which (as has been mentioned already) is the integrated body force subtracted by the skin-friction losses due to the dielectric surface, i.e. $T_x - \tau_w$. Indeed, thrust will be generated by the 8% spatial-duty actuator from $V_0 = 5$ kV. But given its low spatial duty, the thrust value integrated along the span might be comparable to the skin-friction losses up to $V_{ac} = 10$ kV and as such is

4.3. PERFORMANCE CHARACTERISATION OF SPANWISE-MODULATED DBD ACTUATORS



Figure 4.13.: Quantification of the spanwise-modulated-actuator performance. (a) and (b) show the average power consumed (P_{avg}) and thrust generated (T_x) along the wall-parallel direction respectively, as a function of voltage amplitude V_{ac} ($f_{ac} = 2$ kHz). Wall-parallel thrust generated as a function of average power consumed are presented in (c). Solid black curves in (a) and (b) are the power-law fit on the corresponding data of the spanwise-uniform (2D) actuator.

not registered by the employed laboratory force balance.

The wall-parallel thrust generated by all the investigated actuators as a function of the corresponding average power consumption is presented in figure 4.13c. It is observed that the generated thrust is a quasi-linear function of the input power. This is expectable, as both these parameters are proportional to $(V_{ac} - V_0)^2$. The slope of a linear fit through the data points of the spanwise-uniform actuator is 0.34 mN W⁻¹ which is an approximation of the effectiveness of this DBD plasma actuator in converting the electrical input to a mechanical output. The effectiveness of the spanwise-modulated actuators are presented in table 4.1. A general trend can be observed here which is that with reducing spatial duty, the actuator's effectiveness increases. These results suggest that the spanwise-modulated actuators are inherently more efficient than their spanwise-uniform counterpart in performing the electro-mechanical conversion.

When analysing these measurements with the results pertaining to the demonstration of the three-dimensionality of the plasma layer in section 4.3.1, it can be concluded that the stems do not have any effect on the plasma regimes or the ionisation rate. In fact, the performance of these spanwise-modulated actuators scales according to the discharge length. This indicates that modifying the geometry in this manner has no significant effect on the plasma itself (in terms of plasma physics). Thus, the results demonstrate that the performance of the spanwise-modulated actuators can successfully be estimated just by knowing the performance of the corresponding two-dimensional spanwise-uniform actuator. This is particularly useful because fabricating and measuring the performance parameters of two-dimensional actuators are both simple and fast compared to the spanwise-modulated ones.

4.4. DESIGN AND CHARACTERISATION OF ACTUATORS FOR SWEPT-WING BOUNDARY LAYER CONTROL

4.4.1. Design and operation

DBD actuators are known to impart unsteady forcing at the AC frequency of the input waveform (Benard, Noté, et al., 2017), which can trigger early transition in the swept-wing boundary layer, as discussed previously in section 2.2.3. Therefore, forcing at $f_{ac} = 10$ kHz was employed in the current thesis to apply the necessary control over the swept-wing boundary layer. With this as one of the end-goals, 500 μ m thick polyethylene terephthalate (PET) foils was chosen to be the dielectric substrate. These foils have a root-mean-square roughness of 0.02 μ m (see also Rius Vidales et al., 2018), which is low enough for the investigated boundary layer that is highly receptive to surface roughness. A second motivation for this choice is the installation of these actuators over the swept-wing model. These foils have adequate flexibility to be wrapped around the leading edge of the model. As such, this allowed easy substitution of various actuators to investigate different control strategies.

DBD plasma actuators employed to exert the necessary forcing on the swept-wing boundary layer were fabricated with the technique introduced in section 4.2.1 in order the maintain minimum protuberance due to the air-exposed electrode close to the model's leading edge. For all actuators, the air-exposed electrode was a spanwise-uniform strip of conductive silver, while the encapsulated (or covered) electrode design varied according to the forcing necessitated by the control strategy being realized.

The first type of DBD plasma actuators were fabricated such that it generated a spanwiseuniform EFD body force (schematic in figure 4.1a). These actuators are employed to realize a pure base-flow modification (BFM) control strategy. They are composed of two $w_E = w_C = 5$ mm wide uniform strips of conductive silver as electrodes on either sides of the dielectric substrate, with no relative overlap or gap. The spanwise-modulated actuator design (schematic in figure 4.1b) was employed to realize the UFD control strategy. In these actuators, the encapsulated electrode was composed of a series of equi-spaced stems that are $l_S = 10$ mm long and $w_S = 1$ mm wide. These stems are projected from a 5 mm wide uniform silver strip towards the air-exposed electrode. Individual stems have a relative overlap of o = 2 mm with respect to the air-exposed electrode to ensure consistent plasma formation. As already mentioned, the spacing between adjacent stems defines the spanwise wavelength (λ_z) of the plasma forcing. Two such actuators were fabricated for this thesis, their spanwise wavelengths being $\lambda_z = 4$ mm and $\lambda_z = 5$ mm, respectively. The spatial duty of these actuators correspond to $w_S/\lambda_z = 25\%$ and 20% respectively.

The actuators were powered using the GBS Electronik Minipuls 4 high-voltage amplifier which was introduced in section 3.2.3.1. This choice of this power supply here motivated by the necessity of this investigation to apply forcing at a very high rate of $f_{ac} = 10$ kHz. For all test cases, the air-exposed electrode was supplied with a sinusoidal AC signal with voltage amplitudes (V_{ac}) between 2.5 kV and 4.5 kV. The encapsulated electrode was grounded.

4.4.2. Characterisation of electrical and mechanical aspects of employed actuators

This thesis includes the first experimental investigating that employs DBD plasma actuators fabricated with PET foils as their dielectric substrate. Thus, quantification of their performance is vital. Besides this, a key necessity to perform a comparative study between different control strategy as done here, is the characterisation of the employed forcing amplitude, in this case DBD plasma actuators. For this purpose, three DBD actuators, one producing a spanwiseuniform EFD forcing (referred to as '2D') and two generating spanwise-modulated EFD forcing with wavelengths $\lambda_z = 4$ mm and $\lambda_z = 5$ mm, were fabricated. They were all constructed such that, the plasma generating edge of the air-exposed electrode was $x_{PA} = 13$ mm away from the edge of the PET dielectric foil (length x_{PA} depicted figure 4.1). All other design parameters of these actuators are as defined in section 4.4.1. Finally, the covered electrode of all actuators were encapsulated with an epoxy resin to avoid the formation of a plasma discharge on this side of the dielectric substrate.

4.4.2.1. Electrical characterisation

The discharge current signal was obtained by measuring the voltage drop across a shunt resistor $R_s = 100 \ \Omega$ connected in series between the encapsulated electrode and the ground (see section 3.2.3.3). These signals, for the three different DBD plasma actuators being operated at $V_{ac} = 4$ kV and $f_{ac} = 10$ kHz is presented in figure 4.14. The discharge current of the spanwise-uniform actuator is similar to that observed in previous research efforts (see figure 2.2 reproduced from Benard & Moreau, 2012). The current signal during the positive-going half of the sinusoidal AC signal is composed of high-current positive peaks which reach a maximum value of about 70 mA as shown in figure 4.14a. During the negative-going half of the AC cycle, the discharge current is composed of low-amplitude negative peaks.

With the spanwise-modulated plasma actuators (figures 4.14b and c), the polarity of the current peaks during the two half cycles are similar to the previous case. However, the number of peaks (especially in the positive-going half) and their amplitudes are significantly lower. In fact, this low amplitude of the peaks makes the capacitive component of the discharge current, which manifests as a sinusoidal wave, more apparent. This significant difference in the current-peak density and amplitude of the spanwise-modulated actuators is a direct consequence of their spatial duty. As the cumulative length of the discharge of these actuators is lower than the spanwise-uniform actuator, the current-peak density also is reduced. It is worth noting that some spurious current peaks observed with all three actuators (for example, negative current peaks during positive-going half cycle) are a result of faulty encapsulation of the grounded electrode.

The power consumed was estimated through the integral charge measurement technique as described in section 3.2.3.3. The power consumed by the different actuators as a function of AC voltage amplitude is presented in figure 4.15. In this figure (and in figure 4.19), the measurements for the different actuators are normalized with the spanwise length of the encapsulated electrode that results in a plasma discharge l_{CE} and not the length of the plasma discharge as done in section 4.3.2. For the spanwise-uniform actuator this is just the length z_{PA} shown in figure 4.1a. While for the spanwise-modulated actuators, this amounts to the number of stems along the span multiplied by the width of each stem ($l_{CE} = w_S \times$ number of stems). Indeed, as demonstrated in section 4.3, the plasma discharge is not limited here only to the interface of the air-exposed electrodes and the stems of the encapsulated electrode. Estimation of the actual



Figure 4.14.: Discharge current signals (blue, vertical axis on right) of the employed DBD plasma actuators whe operated at $V_{ac} = 4$ kV, $f_{ac} = 10$ kHz. The input high-voltage signal is also presented (red, vertical axis on left).



Figure 4.15.: Average power (P_{avg}) consumed by the DBD plasma actuators versus applied voltage amplitude.

discharge length done in the same manner described in section 4.3.2.1 showed that with these actuators the average length of the secondary discharge is only about 15% of w_S . Compared to the two spanwise-modulated actuators investigated upon before, i.e. the one constructed on Kapton dielectric (section 4.3.1) and those constructed on the PMMA plates (section 4.3.2, see table 4.1), the length of the secondary discharge is lower. Given its weak nature, it is expected to have negligible influence on the boundary layer compared to the dominant component of the



Figure 4.16.: (a) Time-averaged velocity magnitude field (in m s⁻¹) of the wall-jet induced by the DBD plasma actuator when operated at $V_{ac} = 4.5$ kV, $f_{ac} = 10$ kHz. (b) u_p -velocity profile along at $x_p = 6$ mm and the corresponding Glauert (1956) fit for a laminar wall-jet.

generated force along x_p . Hence, its effect is neglected here.

The three actuators constructed for this investigation consume almost similar power at a given operating condition when normalised with the length scale l_{CE} , as shown in figure 4.15. The differences observed is mainly attributed to the aforementioned length-scaling used. This confirms the results discussed in section 4.3.2, that the power consumed by the spanwise-modulated actuator, when normalised with a length scaling that corroborates with the discharge length, is similar to the power consumed by the corresponding spanwise-uniform actuator. However, the power consumed by these actuators is significantly lower compared to those discussed in section 4.3.2. This is a direct consequence of the different dielectric substrates used.

4.4.2.2. Induced jet, generated thrust and EFD body force

The induced velocity fields of the spanwise-uniform actuator were measured using time-resolved PIV, the experimental set-up of which was introduced in section 3.2.2.3. These induced velocity measurements were further used to compute the integral thrust generated by the actuator and the spatial distribution of the EFD body force.

The time-averaged velocity field of the wall-jet induced by the spanwise-uniform DBD plasma actuator when operated at $V_{ac} = 4.5$ kV and $f_{ac} = 10$ kHz is presented in figure 4.16a. This timeaveraged was computed with 3000 instantaneous vector fields, after the initial vortex described in section 2.1.3 left the imaged field and a quasi-steady wall-parallel jet was established. The maximum uncertainty on the time-averaged velocity fields, computed using the linear error propagation techniques (see section 3.2.2.5) is $\epsilon_{\overline{u_p}} \approx 7\%$ and $\epsilon_{\overline{v_p}} \approx 0.6\%$.

The induced jet originates in the vicinity of the interface of the two electrodes and is aligned with the x_p direction. A wall-normal velocity component oriented downwards is observed close to the plasma-generating edge of the air-exposed electrode. This weak suction supplies the necessary mass-flow to the induced wall-jet. The orientation of the vectors within the induced jet are almost parallel to the dielectric surface. Thus, the wall-normal velocity component induced by the actuator is much lower than the wall-parallel velocity component.

The velocity profile along the wall-normal direction (y_p) at $x_p = 6$ mm is presented in figure 4.16b. The profiles bears good resemblance to those reported previously. The analytical solution

of a laminar plane wall-jet computed by Glauert (1956) is as shown below

$$\overline{u}_p = \left(\frac{5F_{EMF}}{2\nu x_p}\right)^{1/2} \cdot f' \tag{4.1}$$

Here, F_{EMF} is the external momentum flux, ν is the kinematic viscosity of air and f' represents the differential of the mass-flux (i.e. velocity distribution normal to the wall). A non-linear least-squares fit of this solution on to the current velocity distribution was carried out and is also shown in figure 4.16b. This further confirm that the induced flow is a wall-jet. The length parameter δ_p depicted in this figure represents the distance of $u_{p,max}$ along the wall-normal direction from the dielectric surface.

The integrated thrust generated by the actuator is computed by the momentum balance technique introduced in section 2.1.4, applied over the presented time-averaged velocity field (figure 4.16a). Pressure gradient due to the acceleration of the flow by the DBD actuator is not readily available from the velocity data obtained from PIV. Thus, the control volume boundaries were set far away from the bulk of the EFD body force which occurs just downstream of the plasma-generating edge of the air-exposed electrode (see figure 2.7b), so that the pressure can be considered to be uniform and constant. The wall-parallel component of the generated thrust is computed by equation 4.2a (adapted from equation 2.3a)

$$T_x = T_x^F + \tau_w = \rho \int_{LHS} \overline{u}_p^2 dy_p + \rho \int_{top} \overline{u}_p \overline{v}_p dx_p - \rho \int_{RHS} \overline{u}_p^2 dy_p$$
(4.2a)

$$\tau_w \approx \rho \nu \int_{x_p=0}^{6mm} \left[\frac{\partial \overline{u}_{p,R}}{\partial y_p} \right]_{y_p=0}^{\delta_p} dx_p \tag{4.2b}$$

where, ρ (= 1.225 Kg m⁻³) is the density of air. *LHS*, *RHS* and *top* represent the boundaries over which the respective line integrals are computed. τ_w is the skin-friction due to the dielectric surface.

It is worth noting that the thrust (T_x) along the wall-parallel direction computed in this manner contains both, the contribution of the EFD body force (T_x^F) and the frictional force (τ_w) due to the dielectric surface, as shown in equation 4.2a (also discussed in section 2.1.4). Given the inadequate spatial resolution close to the dielectric surface, it is not feasible to estimate the skin-friction from the acquired PIV mean fields. Thus, the analytical solution of Glauert (1956) was utilized for this purpose. Through the fitting of Glauert's laminar wall-jet profile on to the measured velocity profile at $x_p = 6$ mm as discussed earlier, the external momentum flux F_{EMF} was found (see equation 4.1). Assuming that this parameter remains constant along streamwise (x_p) direction (as shown by Glauert, 1956), the induced flow field was reconstructed using equation 4.1. The frictional force due to the dielectric surface is then estimated according to equation 4.2b. Here, the differential $\partial u_{p,R}/\partial y_p$ is the slope of the reconstructed velocity profile (see dashed red line in figure 4.16b) from the dielectric surface and up to δ_p (location of $u_{p,max}$).

The wall-parallel thrust generated (T_x) by the DBD plasma actuator as a function of voltage amplitude is shown in figure 4.17. The uncertainty bands in the figure represent the estimated



Figure 4.17.: Variation of wall-parallel thrust generated (T_x) as a function of voltage amplitude (V_{ac}) computed from the time-averaged PIV fields. The uncertainty bands represent the estimated skin-friction.



Figure 4.18.: EFD body force distribution (in kN m⁻³) when the DBD actuator is operated at $V_{ac} = 4.5$ kV, $f_{ac} = 10$ kHz. (a) Computed by solving the Navier-Stokes equation given in 2.4 (see Benard, Debien, et al., 2013). (b) Computed using the empirical model proposed by Maden et al. (2013).

skin-friction due to the dielectric surface. Given that the measured velocity field already accounts for the effect of the skin-friction, this parameter only needs to be negated from the computed thrust as shown in equation 4.2a. Thus, the uncertainty bands presented in the figure have a negative bias. The thrust increases with the voltage amplitude according to a power law as discussed in section 4.3 and shown in figure 4.13b (and in previous research efforts, see Thomas, Corke, et al., 2009; Kotsonis et al., 2011; Debien, Benard, & Moreau, 2012). The wall-normal component of the generated thrust (T_y) is known to be one order of magnitude lower than T_x and hence is neglected.

In order to analyse the influence of the applied forcing on boundary layer stability through the simplified numerical model introduced in section 3.3.1.2, the spatial distribution of the generated EFD body force is desired. The body force distribution estimated by solving the Navier-Stokes (NS) equation (2.4) with the acquired induced velocity fields is shown in figure 4.18a. The acceleration term of the NS equation has a significant influence initially but becomes negligible once a quasi-steady wall-parallel jet is established. The opposite is true for the convection and diffusion terms as can be observed in the figure 2.7a reproduced from Kotsonis et al. (2011). In the current computation, both the acceleration and pressure terms are neglected and the body force is estimated by solving the convection and diffusion terms with the time-averaged induced velocity field.

As demonstrated by previous research efforts, figure 4.18a shows that the maximum body force is concentrated close to the plasma-generating edge of the air-exposed electrode and spans only until $x_P \approx 2$ mm. However, the body force field obtained in this manner is noisy and as such, might not be effective to be used in the simplified numerical model of the plasma forced swept-wing boundary layer. Therefore, the empirical model proposed by Maden et al. (2013) (also used by Kloker et al., see Dörr & Kloker, 2015a) is employed to estimate the steady body force produced by the current DBD plasma actuators. This model assumes that the wall-normal component of the force is negligible ($F_y \approx 0$) and only the horizontal force F_x is computed and utilized. The parameters of the empirical model were calibrated such that, the spatially integrated body-force magnitude matches the horizontal thrust T_x computed from the experimental velocity field presented in figure 4.17. The EFD body force distribution computed through this empirical formulation is presented in figure 4.18b. It is worth noting that here too, the body force is concentrated between $x_p = 0$ and 2 mm as observed with the previous computation. However, compared to figure 4.18a, the computed body force field is much cleaner and is more suitable to be used for modelling purposes.

4.4.2.3. Thrust generated by the spanwise-modulated actuators

Time-averaged thrust generated by the different actuators along the wall-parallel (x_p direction) was measured using a laboratory force balance as described in section 3.2.3.5. These measurements, for the spanwise-uniform and the two spanwise-modulated actuators ($\lambda_z = 4 \text{ mm}$ and 5 mm) are presented in figure 4.19. A significant difference was observed in the measured thrust by the actuator generating a spanwise-uniform EFD body force as compared to that measured through PIV presented in figure 4.17. This is a direct consequence of how the skin-friction is resolved in these two measurements.

As velocity data in the induced flow is not available here (especially $x_p > 6$ mm), the skinfriction is estimated using the analytical solution of a laminar plane-wall jet formulated by Glauert (1956). The computation is performed as shown below.

$$\tau(x_p) = \frac{\rho}{9} \left(\frac{125F_{EMF}^3}{8\nu x_p^5}\right)^{1/4}$$
(4.3a)

$$\tau_w = \int_{x_p=0}^{x_{PA}} \tau(x_p) dx_p \tag{4.3b}$$

Here, F_{EMF} which is the exterior momentum flux, is found by fitting the Glauert (1956) profile on to the velocity profile at $x_p = 6$ mm (see figure 4.16b for example) obtained from the induced velocity fields acquired with PIV, as described previously in section 4.4.2.2. This is then used to compute the skin friction along the induced jet $(\tau(x_p))$ and until the edge of the dielectric surface $(x_p = x_{PA} = 13 \text{ mm})$ according to equation 4.3a. The total frictional force τ_w incurred is then estimated by integrating these values along x_p as shown in equation 4.3b. The quantity measured by the laboratory balance is $T_x - \tau_w$ and as such, the estimated skin-friction is added to the measured values to approximate the actual thrust generated by the actuators.

With the incurred skin-friction accounted for, it can be seen that the thrust generated by the spanwise-uniform actuator obtained here is almost equal to that computed using the momentum



Figure 4.19.: Wall-parallel component of thrust generated (T_x) by the actuator versus applied voltage amplitude. Measurements acquired with the laboratory balance are added with the corresponding estimate of skin-friction. Uncertainty bands (red) are skin-friction estimates for the spanwise-uniform (2D) actuator.



Figure 4.20.: Power consumed and thrust generated by different actuators normalized with the spanwise extent (z_{PA}) on which the respective DBD plasma actuators act upon. Dashed lines are the performance of the spanwise-modulated actuators approximated from that of the two-dimensional actuator. Blue: $\lambda_z = 4$ mm actuator (25% of 2D measurement); Green: $\lambda_z = 5$ mm actuator (20% of 2D measurement).

balance technique on the time-averaged PIV field. For the $\lambda_z = 4 \text{ mm}$ and $\lambda_z = 5 \text{ mm}$ spanwisemodulated actuators no force is measured by the force balance until $V_{ac} = 3.5 \text{ kV}$. Hence, only the measurements at $V_{ac} = 4 \text{ kV}$ are presented in figure 4.19. The measured values with these actuators are very similar to that of the spanwise-uniform actuator at this operating condition. The small differences are mainly attributed to the length-scaling not being the actual extent of the discharge (similar to the power consumed discussed in section 4.4.2.1).

4.4.2.4. Actuator performance parameters normalised with spanwise extent

The discussion of the power consumed and generated thrust in sections 4.4.2.1 and 4.4.2.3 respectively, have not directly accounted for the three-dimensionality of the forcing applied by the spanwise-modulated actuators. To do this, in figure 4.20, the power consumed and thrust generated by the different actuators are normalized with the spanwise extent of the actuator $(z_{PA} \text{ in figure 4.1})$ instead of the discharge length. This length scaling is representative of the performance of the actuators per unit spanwise-length over which they apply boundary-layer flow control. It is worth noting that the wall-parallel thrust shown in figure 4.20b is the estimate of


Figure 4.21.: Splitter-plate edge design (not to scale). The aluminium plate (black), the PMMA splitterplate edge (light gray), electrodes of the DBD plasma actuator (red lines), the epoxy resin encapsulating the grounded electrode (dark gray) are shown. The coordinate system used in this study is also depicted.

the actual thrust generated that accounts for the skin-friction (unfilled markers in figure 4.19).

As mentioned in section 4.4.1, the encapsulated electrode of spanwise-modulated actuators is composed of stems that are each $w_S = 1$ mm wide. This results in a spanwise-modulation with a duty (= w_S/λ_z) of 25% and 20% for the $\lambda_z = 4$ mm and 5 mm actuators respectively. Multiplying these with the measurement of the two-dimensional actuator (dashed lines in figure 4.20) results in a good estimate of the measurements obtained with the respective spanwisemodulated actuators. It is also important to note is that, for unit spanwise-extent, the spanwise-modulated actuators consume significantly less power compared to the two-dimensional actuators. Thus, if the applied control is effective in delaying laminar-to-turbulent transition, the beneficial effects would have been realized with significantly less power consumption.

4.5. DESIGN OF ACTUATORS FOR MIXING LAYER INSTABILITY CONTROL

Two DBD plasma actuators were designed and fabricated to investigate the effect of spanwiseuniform and spanwise-modulated EFD forcing on the mixing layer flow configuration. The DBD plasma actuators were constructed just before the trailing-edge of the splitter plate, on the PMMA attachment as shown in figure 4.21. As mentioned previously, an incision made on the bottom side from x = -42 mm to x = -2 mm (dark gray area in figure 4.21) resulted in a 2 mm thick PMMA plate on the top side which constituted the dielectric substrate for the actuator.

Ho & Huang (1982) show that the fundamental KH instability in the mixing layer is characterised by the high-velocity boundary layer. This is because, the higher shear contained within this boundary layer (compared to the low-velocity one) has a greater influence on turbulent energy production within the mixing layer and thus, linear instability growth and formation of coherent vortices. Thus, flow control is expected to be more effective if applied in this boundary layer. This was confirmed by Ely & Little (2013) who observed that the DBD plasma actuator has more control authority on the mixing layer when the EFD forcing is applied in the splitterplate boundary layer generated by the high-velocity stream. Thus, in the current experiment, the air-exposed electrode of the DBD actuator (at whose edge the EFD force is generated) was flush-mounted on the top (high-velocity) side of the PMMA attachment. The second electrode was installed inside the incision which was then filled with an epoxy resin to encapsulate the second electrode, thus avoiding any unwanted spark generation between the electrodes.

For the two forcing configurations being investigated here (i.e. spanwise-uniform and spanwise-modulated), two separate PMMA attachments were fabricated bearing the respective DBD plasma actuators. In both cases, the air-exposed electrodes were a uniform strip of aluminium with a non-reflective black coating. The total thickness (aluminium and coating) being 85 μ m. These electrodes were $w_E = 50$ mm wide. For the spanwise-uniform actuator, the encapsulated electrode was $w_C = 15$ mm wide uniform strip of aluminium. The electrodes were mounted asymmetrically, with a relative overlap of o = 5 mm. While for the spanwise-modulated actuator, the encapsulated electrode was composed of a series of equi-spaced stems (similar to that shown in figure 4.1b). The stems were $l_S = 13$ mm long and $w_S = 3$ mm wide. These stems were projected from a 2 mm wide uniform aluminium strip towards the air-exposed electrode. Here too, the electrodes were mounted asymmetrically, with a relative overlap of o = 5 mm. The distance between adjacent stems, which defines the spanwise-wavelength of the exerted EFD forcing (section 4.3.1), was $\lambda_z = 37$ mm. This results in a spanwise spatial duty of $w_S/\lambda_z = 8\%$.

These constructions resulted in DBD plasma actuators that exert the respective EFD body force into the boundary layer generated by the high-velocity stream along the free-stream direction at x = -12 mm (which is the location of the plasma-generating edge of the air-exposed electrode). The EFD forcing is applied along the streamwise direction here as the intention is to impart fluctuations into the mixing layer flow field in order to amplify certain instability modes. For both actuator constructions, the air-exposed electrode was supplied with the high-voltage AC while the encapsulated electrode was grounded. The current mixing layer flow configuration does not necessitate the application of forcing at very high frequencies as was the case with the swept-wing boundary layer. Thus, the actuators were powered using the more stable *Trek* 30/20A high-voltage amplifier that was introduced in section 3.2.3.1.1.

The reader is reminded now that the construction of the two actuators employed here are similar to those utilized to quantify the performance of the spanwise-modulated actuators in section 4.3.2. As with the actuators here, those actuators were constructed on 2 mm thick PMMA plates as well. As such, the geometric parameters of the actuators deployed are the exactly similar to the spanwise-uniform actuator and the modulated device with a spatial duty of 8% investigated in section 4.3.2. Thus, the performance characterisation of the actuators designed for mixing layer instability control has already been reported and is not repeated here once again. For the sake of completeness, the results pertaining to only these two actuator configurations are reproduced from figure 4.13 and presented in figure 4.22.

4.6. SUMMARY & CONCLUDING REMARKS

In this chapter, the discussion is centred around DBD plasma actuators, which is the active flow/instability control device employed in this thesis. While plasma actuators are known to exercise the necessary flow control in a non-intrusive manner, in certain situations, the thickness of the actuator's electrodes can have a detrimental effect on the flow field in consideration. Thus, a novel fabrication technique that produces very thin (micrometre scales) electrodes with intricate geometries was proposed. This technique involves computer-controlled spraying of conductive silver particle over the dielectric substrates through opportune masks that bear the geometric features of the desired electrodes. To validate the performance of the produced actuators, two DBD plasma actuators that generate a two-dimensional (spanwise-uniform) EFD body force



Figure 4.22.: Performance parameters of the two DBD plasma actuators employed for mixing layer flow control. (a) and (b) show the variation of average power consumed (P_{avg}) and the wall-parallel thrust generated (T_x) respectively, as a function of AC voltage amplitude V_{ac} ($f_{ac} = 2$ kHz). Thrust generated as a function of average power consumed is presented in (c).

were fabricated, the difference between them being the number of passes executed with the spraying nozzle to fabricate the air-exposed electrodes. The generated electrode edge was not uniform but featured micrometre-scale undulations. The thickness measurements revealed that the electrode thickness can be varied simply by varying the number of passes executed with the spraying nozzle.

Measurements of the power consumed by these novel actuators and the corresponding thrust generated show no significant difference at the tested voltages and frequencies. In terms of these two parameters, the performance of these actuators are found to be very similar to the standard actuator constructed with the air-exposed electrode made of conductive copper tape. The similarity of these measurements for the three different actuators demonstrates the insensitivity of the global performance of the actuators toward both the quality of the plasma-generating edge and thickness of the air-exposed electrodes. Studying the PIV vector fields of the flow induced by the two novel actuators reveals that a thinner air-exposed electrode features maximum induced velocity closer to the wall. Due to the concentration of the electric field in case of a thinner electrode, the actuator seems to impart momentum closer to the wall. This is a vital result for flow control applications of DBD actuators. As the thickness of boundary layers are dependent on a number of parameters, it is essential to apply the EFD body force close to the wall to ensure effectiveness of flow control for a wide range of conditions. The results obtained here suggests that such a forcing scenario can be achieved by reducing the thickness of the air-exposed electrode and increasing the applied voltage. Further research however is necessary to establish the underlying phenomenon that influences the topology of the ionic wind generated by DBD actuators. While the maximum velocity being closer to the wall suggests a stronger local electric field with a thinner air-exposed electrode, no increase in thrust generation is observed. This can be due to viscous effects due to skin-friction, which are more prominent closer to the wall. Thus, computation and analysis of these effects are vital as well.

A general description of the geometric features of the two types of actuators, i.e. those that apply spanwise-uniform and spanwise-modulated EFD forcing, is also given in this chapter. While the performance characteristics of the spanwise-uniform actuator is well known, the same is not the case with the spanwise-modulated actuator configuration. Thus, a thorough characterisation of these actuators was undertaken. Firstly, through fast imaging of the developing plasma layer and velocity field measurements obtained from tomographic PIV, the three-dimensionality of the flow induced by these modulated actuators is confirmed. While the plasma discharge was expected to occur only in the spanwise regions where the air-exposed electrode intersects with the stems of the encapsulated electrode, the images of plasma layer revealed the occurrence of a "secondary" discharge in the vicinity of the stems.

Quantification of the spanwise-modulated actuators' performance by measuring the power consumed and the wall-parallel thrust generated, and comparing them with a corresponding spanwise-uniform actuator reveals some important aspects. Firstly, when the acquired data is normalised with the cumulative length of the modulated discharge, the performance of these actuators is very similar to the spanwise-uniform actuator. This indicates that the current spanwise-modulated actuator geometry is very similar to the classical geometry in terms of the underlying plasma physics. Additionally, the performance of these can be estimated from those of a corresponding spanwise-uniform actuator that generates a two-dimensional forcing. Furthermore, the effectiveness of the actuator in converting the electrical power input into mechanical energy output (i.e. force) is observed to increase with decreasing spatial duty of the spanwise-modulated actuator. This is expected to be due to the modulation resulting in stronger electric fields and thus a better electro-mechanical conversion. Finally, accounting for the spatial duty of these modulated actuators results in a reduction of the magnitude of the performance parameters. This is especially important with respect to the power consumed by the actuator. The improvement of the electro-mechanical conversion effectiveness of these actuators combined with reduced overall power consumption suggests that application of these actuators can significantly improve the efficiency of the flow control system.

5 SWEPT-WING BOUNDARY LAYER CONTROL

The outcome of investigations pertaining to the swept-wing boundary layer is reported in this chapter. Modern jet aircraft feature wings with swept leading edges to increase the critical Mach number. As introduced in section 2.2, this geometry gives rise to a three-dimensional laminar boundary layer, featuring the so-called cross-flow (CF). The CF component inherently forms an inflection point, which results in the inviscid cross-flow instability which is characterised by co-rotating vortices of stationary (steady) or travelling (unsteady) nature, depending on external disturbances (Bippes, 1999; Downs & White, 2013). CF instability has a pivotal role in the process of laminar-to-turbulent transition. As such, suppression of these instabilities presents a promising route towards transition delay and aerodynamic drag reduction (Joslin, 1998).

In this chapter, DBD plasma actuators are deployed to realize two different strategies to control CF instabilities and the boundary layer's eventual transition to turbulence. These being upstream flow deformation (UFD, section 2.2.2.1) and base-flow modification (BFM, section 2.2.2.2). The UFD strategy will be realized with a spanwise-modulated plasma actuation while the BFM strategy will be attempted with a simple spanwise-uniform actuation. While the proficiency of plasma-based flow control to successfully realize these strategies have been demonstrated through numerical simulations (see Dörr & Kloker, 2015b, 2017), corresponding experimental proof is limited. Applying spanwise-modulated EFD forcing with DBD actuators to realize the UFD control strategy was experimentally investigated recently (Schuele et al., 2013; Serpieri, Yadala, et al., 2017). While the actuators were successful in conditioning the boundary layer with the forced sub-critical CF instability modes, no transition delay was reported in either of the studies. With regards to the plasma-based BFM control, the results presented in this chapter constitute the first attempt to realize this strategy experimentally. Finally, the receptivity of swept-wing boundary layers dominated by CF instabilities to the EFD forcing applied by the DBD actuator is poorly understood. Recent efforts (Dörr & Kloker, 2016; Serpieri, Yadala, et al., 2017; Guo & Kloker, 2020) have shown that the unsteady nature of the exerted forcing can have a detrimental effect on the boundary-layer stability by introducing growing travelling modes. These can lead to transition anticipation, which essentially curtails the effectiveness of the applied control.

The primary objective here is to investigate the effectiveness of high-frequency plasma actuation to delay laminar-to-turbulent transition in the swept-wing boundary layer through the UFD and BFM control strategies. To this effect, specially designed actuators were fabricated using the computer-controlled spray-on technique introduced in section 4.2.1. The actuator fabrication was carried out in this manner due to the sensitivity of this boundary layer to surface features (see section 2.2.1). Furthermore, the numerical simulations of Dörr & Kloker (2016) indicate that the spanwise-modulated body force being exerted to realize the UFD strategy can be expected to perform a BFM action locally. However, the effect of this action and how the two strategies contribute to controlling transition in such a situation is not well understood and as to date, not been explored experimentally in a systematic manner. Apart from being the first experimental investigation of the effectiveness of the BFM strategy, the first comparison of this control method with the UFD strategy under identical experimental flow conditions is also made here. While this facilitates the comparison of the two strategies in terms of effectiveness and efficiency, it also enables the examination of the collaboration (if there is any) between the UFD and BFM control strategies.

This chapter is organised as follows. The flow conditions at which the current experimental investigation will be carried out is introduced in section 5.1. Results of linear stability analysis in order to predict the critical stationary CF instability mode at the investigated flow conditions is presented in section 5.2. The conditioning performed on the baseline (unforced) boundary-layer configuration is introduced in section 5.3. Details regarding the installation of the deployed DBD plasma actuators and quantification of the forcing amplitude through the computation of a momentum coefficient is given in section 5.4. The experimental results pertaining to the application of spanwise-modulated EFD forcing to realize the UFD control strategy is reported in section 5.5. Results corresponding to the plasma-based BFM control method, both from the experimental domain and the theoretical domain through the simplified model introduced in section 3.3.1.2, is discussed in section 5.6. This study will be incomplete without a comparison between the control efficiency of the two strategies. This is done by estimating the control efficiency in terms of total drag reduction achieved with the two control methods. The formulation of this estimation procedure and the corresponding comparison of the UFD and BFM strategies is presented in section 5.7. The conclusions drawn from this investigation and some outlook for future research is in section 5.8.

5.1. INVESTIGATED FLOW CONDITIONS

Experimental investigations pertaining to controlling the CF instability prevalent in a sweptwing boundary layer was carried out in the low turbulence wind-tunnel at TU Delft, Netherlands. The model employed in the current investigation is a 45° swept-wing referred to as M3J. The streamwise chord of this model c = 1.27 m is used in this chapter to non-dimensionalise the spatial coordinates. The employed wind-tunnel and swept-wing model was detailed in section 3.1.1. For the entirety of this investigation, the wing's incidence angle was fixed at $\alpha = 3^{\circ}$ and the flow over the wing's pressure side will be investigated. For most of the tested cases, the free-stream velocity was $U_{\infty} = 25 \text{ m s}^{-1}$ ($M_{\infty} = 0.07$), which corresponds to a chord Reynolds number of $Re_c = 2.1 \cdot 10^6$. At these flow conditions, the pressure distribution along the streamwise direction is invariant along the span as shown in figure 3.3, indicating the realization of the infinite swept wing hypothesis, which is an absolute necessity for studying CF instability (Deyhle & Bippes, 1996; Reibert et al., 1996; Radeztsky et al., 1999). Certain cases were also acquired at $U_{\infty} = 27.5 \text{ m s}^{-1}$ ($Re_c = 2.3 \cdot 10^6$) and $U_{\infty} = 30$ m s⁻¹ ($Re_c = 2.5 \cdot 10^6$). The infinite swept-wing condition is satisfied at these Reynolds numbers as well. Finally, the coordinate/reference system used in this chapter is as described in section 3.1.3.2.



Figure 5.1.: N-factors of different stationary CF instability modes (f = 0 Hz) of the laminar boundary layer generated on the pressure side of the swept-wing at flow conditions $U_{\infty} = 25$ m s⁻¹ ($Re_c = 2.5 \cdot 10^6$) and $\alpha = 3^\circ$.

5.2. LINEAR STABILITY ANALYSIS TO DETERMINE CRITICAL STATIONARY CF INSTABILITY MODE

Linear stability analysis is utilized to determine the stability characteristics of the investigated boundary layer, especially the spanwise wavelength of the critical stationary CF instability mode at the investigated flow conditions of $Re_c = 2.1 \cdot 10^6$ ($U_{\infty} = 25 \text{ m s}^{-1}$). The stability variables were computed following the procedure introduced in section 3.3.1 (3.3.1.1). The laminar boundary layer solution on the pressure side of the swept-wing was computed from the corresponding experimentally-acquired pressure distribution (figure 3.3). The obtained boundary layer solution was then used to compute and analyse the stability characteristics of the boundary layer. The amplification N-factors along the chordwise direction of select stationary CF instability modes in the laminar boundary layer was computed by equation 3.6 and are presented in figure 5.1.

At the most downstream location (close to the pressure minima), the most dominant instability has a wavelength of $\lambda_z = 8$ mm. Thus, in the current boundary layer, the critical stationary CF instability has a wavelength of $\lambda_{z,crit} = 8$ mm. LST predicts the onset (neutral point) of this mode to be at x/c = 0.0068 and shows a continuous growth up to the most downstream chordwise station shown. The N-factors of two smaller-wavelength or subcritical stationary modes relevant to the UFD control strategy, $\lambda_z (= \lambda_{z,crit}/2) = 4$ mm and $\lambda_z (\approx \lambda_{z,crit} \cdot 2/3) = 5$ mm are also shown in the figure. These modes are dominant further upstream due to the lower local Reynolds number. However, they saturate and marginally decay along the streamwise direction, leading to the critical mode becoming dominant. The amplification N-factors of a super-critical mode with wavelength $\lambda_z = 13$ mm is also presented. This mode, while being unstable, is never the dominant stationary CF instability mode. As such, super-critical modes do not attenuate the critical mode and can potentially advance transition (Saric, Carrillo, et al., 1998).

Computations corresponding to the evolution of unsteady primary (travelling) CF instability modes are not presented here as it is outside the scope of the present study. For the current experimental conditions, this exercise is reported in Serpieri & Kotsonis (2016) who found (experimentally and through LST) the frequency band of the travelling CF instability modes to be 350 Hz $\leq f \leq 550$ Hz. It must be noted here that the results from LST analysis provides the evolution of each instability mode individually. As such, it does not account for any nonparallel and non-linear effects. The initial amplitude of the different instability modes, which



Figure 5.2.: IR thermography mean fields. The flow comes from left. (a) Boundary-layer transition topology of the clean wing. (b) Boundary-layer transition topology with an array of DREs installed at x/c = 0.017, forcing the $\lambda_{z,crit} = 8$ mm stationary CF instability mode. In both figures, PET foils are mounted on the swept-wing model but no DBD plasma actuator is fabricated on the top most foil.

is a product of flow receptivity, is not accounted for as well in these computations. Furthermore, the model curvature and the corresponding flow non-parallel effects are not considered. Nonetheless, the obtained results give valuable information about the stability characteristics of the boundary layer.

5.3. DRE ARRAY FOR CONDITIONING THE UNFORCED BOUNDARY LAYER

At the flow conditions being studied ($Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$), natural transition on the employed swept-wing model is downstream of the pressure minimum point of $x/c \approx 0.63$ (figure 3.3). As such, the transition process is not only dominated by the CF instability mechanism but also by laminar separation as shown in figure 5.2a. Thus, the critical stationary CF instability mode was forced using an array of discrete roughness elements (DRE).

The most unstable or critical stationary CF instability mode in the investigated boundary layer was identified, through the LST analysis presented in section 5.2, to have a spanwise wavelength of $\lambda_{z,crit} = 8$ mm. Furthermore, LST predicts the onset (neutral point) of this mode to be at x/c = 0.0068 and shows a linear growth up to the most downstream chordwise station analysed (see figure 5.1). This critical stationary CF instability mode was forced with the DRE array. The arrangement was composed of cylindrical DREs with a diameter of $d_{DRE} = 2$ mm and height $k_{DRE} = 60 \ \mu$ m such that, the distance between adjacent DREs was ($\lambda_{z,crit} =$) 8 mm. The DRE array was installed at x/c = 0.017, just downstream of the neutral point of critical CF instability mode (see figure 5.3a). The ratio of DRE height and the local boundary layer displacement thickness is $k_{DRE}/\delta^* = 0.47$.

The result of this DRE forcing on the swept-wing boundary layer and its eventual transition can be seen in figure 5.2b. The DRE forcing assists in bringing the transition upstream to about $x_{tr}/c \approx 0.32$ and achieving CF instability-dominated laminar-to-turbulent transition. Another factor that holds utmost importance is that amplification of the critical mode presents a "worst case scenario" for any given control scheme. Thus, application of the DRE array gives a maximum bound on the performance of the control. This swept-wing boundary layer condition, i.e. one being conditioned by the DRE array at a spanwise wavelength $\lambda_z = 8$ mm is considered to be the baseline case over which the effect of the applied plasma flow control will be investigated.



Figure 5.3.: (a) Schematic of the experimental set-up (not to scale). Electrodes of the DBD plasma actuator (black and grey bars), outline of the dielectric PET foil (black line, dashed part wrapped around the leading-edge), DRE (white dots) and chordwise location of the DREs and DBD plasma actuator are shown. IR FOVs (red rectangles) are also represented. (b) Image of the swept-wing model close to the leading edge showing the DREs and one of the employed DBD actuator while being operated.

5.4. PLASMA FORCING LOCATION AND MOMENTUM COEFFICIENT

5.4.1. Installation of DBD actuators on the swept-wing model

The installation of the various DBD plasma actuators on the swept-wing model is briefly described here. It is absolutely necessary to ensure that if there is a breakdown of the any of the used DBD actuators, it does not damage the surface of the wing model. Thus, two foils of the 500 μ m thick PET foils were first installed on the model surface. The upper foil was attached on to the lower foil (which was larger in size) such that their bottom edges were aligned with each other. The upper foil was geometrically and dimensionally similar to the PET foil on which the DBD plasma actuators were fabricated, so as to simplify the alignment of different actuators while being installed.

All the PET foils, including the one with the DBD actuator, were attached to the underlying surface using an aerosol glue. Additionally, a 50 μ m thick self-adhesive Kapton strip was pasted on the encapsulated electrode of the actuator before spraying the aerosol glue. This was done to preserve the encapsulated electrode and ensure that the DBD plasma actuator can be reused if necessary. The foils were wrapped around the leading-edge in order to avoid a forward-facing step in this region and also to conserve the wing's airfoil shape. Furthermore, these foils extended beyond the pressure minimum point x/c = 0.63 to form an uninterrupted smooth surface. The end of the PET foil that results in a backward-facing step is in fact noticeable in IR fields of figure 5.2 (also in figures 3.6b and c). The electrical connections of the electrodes to the power source were done on the suction side of the wing.

For all tested cases, the DBD plasma actuators were installed on the swept-wing model such that the plasma-generating edge of the air-exposed electrode was at x/c = 0.035. This is just downstream of the array of DREs at x/c = 0.017, forcing the critical CF instability mode $\lambda_{z,crit} = 8$ mm, as depicted in figure 5.3a. It is worth noting here that the DRE array was fabricated simultaneously with the DBD plasma actuator using the *Tooli* XY plotter. Multiple layers of the self-adhesive black sheet used to generate the mask for the DBD plasma actuator



Figure 5.4.: Momentum coefficient versus applied AC voltage amplitude, the AC frequency being $f_{ac} = 10$ kHz. (a) Spanwise-uniform EFD forcing at three different Reynolds number. (b) Spanwise-modulated EFD forcing at $Re_c = 2.1 \cdot 10^6$.

fabrication (see step 2 in section 4.2.1) were stuck one on top of the other to arrive at a height of $k_{DRE} = 60 \ \mu\text{m}$. The laser tool head of the XY plotter was employed to cut cylinders with a diameter of $d_{DRE} = 2 \ \text{mm}$, 8 mm apart from each other, resulting in the necessary DRE array. A photograph of the swept-wing model's leading edge portraying the installed DRE array, followed by the spanwise-modulated DBD actuator generating the corresponding discharge is shown in figure 5.3b.

5.4.2. Momentum coefficient

In order to compare the actuators' control authority for different forcing configurations and conditions, it is necessary to non-dimensionalise the applied force and to obtain an estimate of its magnitude relative to the experimental boundary layer. To this effect, the integrated thrust values T_x are used to compute the momentum coefficient c_{μ} according to equation 5.1,

$$c_{\mu} = \frac{T_x}{\frac{1}{2}\rho u_e^2 \theta_e} \cdot \frac{w_S}{\lambda_z} \tag{5.1}$$

where, u_e and θ_e are the boundary layer edge velocity and momentum thickness respectively, at the location where EFD forcing is applied by the actuator (x/c = 0.035), extracted from the laminar boundary layer solution at the investigated flow conditions. The generated thrust computed by applying the momentum balance technique over the time-averaged induced velocity fields captured using PIV (figure 4.17) is used for computing c_{μ} . This is chosen here as the EFD body-force distribution is estimated based on these measurements. An additional ratio w_S/λ_z , which represents the spanwise spatial duty of the spanwise-modulated actuators, is part of the equation here. This is done to account for the three-dimensionality of the applied forcing, following the discussions in sections 4.4.2.4. For spanwise-uniform forcing, this ratio is equal to 1 and the equation 5.1 reduces to the one defined and used by Serpieri, Yadala, et al. (2017) (also by Guo & Kloker, 2020).

The variation of c_{μ} for spanwise-uniform EFD forcing with input voltage amplitude at three different flow conditions ($Re_c = 2.1 \cdot 10^6$, $2.3 \cdot 10^6$ and $2.5 \cdot 10^6$) studied here is presented in figure 5.4a. Similar to the generated thrust shown in figure 4.17, the momentum coefficient too increases with input voltage amplitude. However, it is observed to reduce with increasing



Figure 5.5.: Schematic of the forcing configurations to realize the UFD control strategy. The DRE arrangement (•) and the DBD actuator with its encapsulated electrode (orange, with stems) and air-exposed electrode (grey bar) are depicted. Consecutive inviscid streamlines originating from the DRE are also shown. (a) UFD4_{in} configuration; (b) UFD4_{out} configuration; (c) UFD5 configuration.

free-stream velocity. This is expectable as c_{μ} is inversely proportional to the square of the boundary layer edge velocity (U_e) which increases with free-stream velocity. Variation of c_{μ} for the spanwise-modulated forcing configurations at $Re_c = 2.1 \cdot 10^6$ is presented in figure 5.4b. The momentum coefficient of these actuators are significantly lower compared to the spanwiseuniform actuator owing to the additional ratio w_S/λ_z introduced into the equation 5.1 to account for the spatial duty of these actuators. In this thesis, the multitude of cases being investigated provides a large range of c_{μ} values. Considering only the flow conditions at $Re_c = 2.1 \cdot 10^6$ the lowest tested value is $c_{\mu} = 0.01$ ($\lambda_z = 5$ mm actuator at $V_{ac} = 2.5$ kV) and the highest tested value is $c_{\mu} = 0.71$ (spanwise-uniform actuator $V_{ac} = 4.5$ kV). Henceforth, c_{μ} will be employed to discuss the different forcing conditions.

5.5. UPSTREAM FLOW DEFORMATION

Spanwise-modulated DBD plasma actuators, the geometric features of which was introduced in section 4.4.1, is now employed to force specific stationary CF instability modes, in order to realize the UFD control strategy. It is worth noting that all the test cases presented in this section correspond to the flow conditions at a chord Reynolds number $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$. This is because of the inherent dependence of the UFD control strategy on the wavelength of the critical stationary CF instability. Given that the critical mode varies with Reynolds number, if the flow conditions are changed, the 8 mm stationary mode forced by the DRE array will not remain the critical mode and the observed effects on the boundary layer will not be a consequence of only the UFD action of the employed actuators.

5.5.1. UFD strategy forcing configurations

In this study, the use of an array DREs not only conditions the boundary layer by locking the critical stationary CF instability mode $\lambda_{z,crit} = 8$ mm, but also fixes the spanwise position of the CF vortices along the inviscid streamline. In doing so, it provides an ideal experimental

condition to investigate how the UFD and BFM strategies realized with DBD plasma actuator forcing collaborate with or counteract one another. To this effect, two DBD actuators (see section 4.4.1) that perform UFD control by generating spanwise-modulated jets to force a single monochromatic stationary CF instability mode were designed and fabricated. The first actuator was designed to force the $\lambda_z (= \lambda_{z,crit}/2) = 4$ mm sub-critical CF instability mode and had a fixed (spatial) phase relation with the oncoming CF vortices due to the DRE array. The second actuator forced the $\lambda_z (\approx \lambda_{z,crit} \cdot 2/3) = 5$ mm CF instability mode and hence had a spanwise-varying phase relation with the oncoming CF vortices.

With the DBD actuator forcing the sub-critical stationary CF instability mode with wavelength $\lambda_z = 4$ mm, by manipulating the relative spanwise position of the DBD actuator with respect to the DREs, two different forcing configurations were realized. In the first, the DBD actuator was installed such that, the incoming inviscid streamline passes over the station x/c = 0.035, at the electrodes' interface where the EFD force is being exerted (see figure 5.5a). The inviscid streamline was computed from the experimentally acquired pressure distribution presented in figure 3.3 (also used to compute the laminar boundary layer) assuming an invariant CF velocity component w. As the stationary CF vortices are roughly aligned with the inviscid streamline, the plasma jets are "in-phase" with the incoming CF vortices (henceforth referred to as UFD4_{in} configuration). For the second configuration, the DBD actuator is displaced along the span (zdirection) by 2 mm resulting in the plasma jets being 'out-of-phase' with the incoming CF vortices (henceforth referred to as UFD4_{out} configuration). In this case, the stationary CF vortices pass between two adjacent stems of the encapsulated electrode (or plasma jets) as shown in figure 5.5b.

The second DBD actuator forcing the $\lambda_z = 5$ mm sub-critical stationary CF instability modes provides the third forcing configuration (figure 5.5c, henceforth referred to as UFD5 configuration). This is such that every fifth CF vortex passes directly over the electrodes' interface at x/c = 0.035 (as in UFD4_{in} configuration). While the other four vortices pass the x/c = 0.035chordwise station at variable distances from plasma sources, thus resulting in a spanwise varying phase relation between the two.

5.5.2. Boundary layer and transition topology through IR thermography

The topology of the swept-wing boundary layer and the eventual laminar-to-turbulent transition is analysed through IR thermography. The set-up of this measurement system is as described in section 3.2.1. Two different FOVs were captured for each test case (see figures 3.5b and 5.3a). The first one captured the thermal signature of the flow over the entire swept-wing model (global IR-FOV) and was employed mainly to study the transition topology. The second was a zoomed-in view (zoomed-in IR-FOV) in the region of the wing model with a laminar flow. These fields were acquired to study the wavelength of the different stationary CF instability modes prevalent in the boundary layer. It should be stressed here once again that the presented fields are raw images captured by the IR camera and as such, no thermal calibration was performed on the recorded IR images. However, prior to any analysis, dewarping/deskewing of the acquired images is applied using predefined calibration images, thus shedding the wing surface curvature and parallax (viewing angle) errors as discussed in section 3.2.1.

The flow and transition topology at the investigated flow conditions ($Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25$ m s⁻¹), visualized using IR thermography, is presented in figure 5.6. The upper left-corner of these images is masked (grey triangle) as it falls upstream of the swept-wing leading edge. In the



Figure 5.6.: IR thermography mean fields. DBD plasma actuators were operated at $V_{ac} = 4 \text{ kV}$, $f_{ac} = 10 \text{ kHz}$. (a) UFD4_{in} unforced flow; (b) $c_{\mu} = -0.068$; (c) Subtraction of (a) from (b); (d) UFD4_{out} unforced flow; (e) $c_{\mu} = -0.068$; (f) Subtraction of (d) from (e); (g) UFD5 unforced flow; (h) $c_{\mu} = -0.055$; (i) Subtraction of (g) from (h). The solid white lines represent constant chord positions.

baseline cases (no EFD forcing) shown in figures 5.6a, d, and g, the DBD actuators corresponding to UFD4_{in}, UFD4_{out}, and UFD5 forcing configurations respectively, are mounted on the sweptwing model but not operated. The bright areas in these figures correspond to the regions on the swept-wing where the flow is in laminar regime while the darker areas correspond to a turbulent flow. In all baseline IR fields, the laminar breakdown occurs along the characteristic jagged pattern reported for stationary CF instability-induced transition (Saric, Reed, et al., 2003; Saric, Carpenter, et al., 2011; Downs & White, 2013; Serpieri & Kotsonis, 2016). Furthermore, the transition for all baselines occurs at approximately the same chordwise position as that with just the DRE array conditioning the boundary layer with the critical CF instability mode ($x_{tr}/c \approx 0.32$). This indicates that there is no influence from the air-exposed electrodes of actuators fabricated through the novel technique presented in section 4.2.1, over the boundary layer and its eventual transition. This indicates the success of the proposed fabrication techniques in achieving one of its core goals: producing electrodes with thickness of a few micrometres which does not influence the boundary layer in any manner.

As shown previously by Serpieri, Yadala, et al., 2017, stationary nature of the CF vortices provides the opportunity to quantify their spatial wavelength. Wavelength power spectral density analysis along constant-chord lines in the regions enclosed by dashed (white) polygons in the IR fields of the three baselines are presented in figure 5.7a, c and e. These computations were performed over the *zoomed-in IR-FOV* captured in this region. The dominance of the DRE-triggered $\lambda_{z,crit} = 8$ mm mode is evident in all three baseline cases. Given that no thermal calibration was performed on the recorded IR images, this spectral analysis only aides in identifying the spectral components and not the amplitude of CF instability modes.



Figure 5.7.: Wavelength power spectral density of IR pixel intensity along the chord. The analysis is performed in the region enclosed by the dashed (white) polygons in figure 5.6. (a) UFD4_{in} unforced flow; (b) $c_{\mu} = -0.068$; (c) UFD4_{out} unforced flow; (d) $c_{\mu} = -0.068$; (e) UFD5 unforced flow; (f) $c_{\mu} = -0.055$.

The flow topology as a consequence of the EFD forcing applied by the respective DBD plasma actuators operated at $V_{ac} = 4$ kV and $f_{ac} = 10$ kHz is shown in figures 5.6b, e and h. In all instances, a bright band is observed in the top-left corner. This is the thermal signature of the dielectric barrier discharge, ensuing from the actuator operation. Furthermore, the topology of the transition front allows further evaluation of the characteristics of the underlying instabilities. More specifically, a spanwise uniform laminar-to-turbulent transition front would suggest that the transition is dominated by travelling CF instability modes. As these unsteady modes propagate in a direction perpendicular to the local inviscid streamline, they distribute or "smooth-out" the spanwise velocity modulation caused by stationary CF instability modes in the boundary layer and hence, also erode the characteristic jagged transition front still has a robust and persistent jagged pattern when EFD forcing is applied, which suggests that the transition to turbulence is still dominated by the stationary CF instability modes.

The results of wavelength power spectral density analysis carried out in the regions within the dashed (white) polygons of the IR fields with EFD forcing are shown in figures 5.7b, d and f (through the zoomed- IR FOV measurement). Although the spanwise-modulated DBD plasma actuators are configured to trigger sub-critical CF instability modes, these modes are not observed in the power spectra and the dominant CF instability mode still has the wavelength of $\lambda_{z,crit} = 8$ mm in all three cases. This is unlike the observations of Serpieri, Yadala, et al. (2017) who observed the sub-critical mode forced by their DBD plasma actuators in the IR flow visualisation of the boundary layer.

It is reasonable to expect that the actuators did condition the boundary layer with the respective sub-critical CF instability modes. However, the critical $\lambda_{z,crit} = 8$ mm modes introduced by the DRE array as well as the mean-flow deformation caused by them could result in the forced sub-critical CF instability modes saturating at a very early chordwise station. Unfortunately, due to FOV limitations, it is not possible at this stage, to asses particularities pertaining to the



Figure 5.8.: Experimentally determined laminar-to-turbulent transition-front displacement $(\Delta x_{tr}/c)$ versus momentum coefficient c_{μ} for the three forcing configurations investigated here.

forced sub-critical CF instability modes further upstream and close to the forcing location. In addition, for the case of UFD5 forcing configuration, spectral energy is observed in the mode with wavelength $\lambda_z \approx 13$ mm. The amplification of this super-critical CF instability mode is expected to be due to the constructive interaction of the critical and sub-critical modes being forced by the DRE array and DBD actuator respectively, however, no definitive confirmation can be made with the available data.

5.5.3. Location of laminar-to-turbulent transition

To illustrate the movement of the transition front due to the EFD forcing better, figures 5.6c, f and i show the subtraction of the mean IR fields of the baseline cases (figure 5.6a, d and g) from those corresponding to the respective cases with EFD forcing (figure 5.6b, e and h). In these figures, the bright regions along the transition front represent a downstream shift of the laminar-to-turbulent transition. In all three forcing configurations, the laminar-to-turbulent transition shifts downstream with EFD forcing along many parts of the transition front. Among the three presented spanwise-modulated forcing cases, the UFD4_{in} configuration shows the most transition delay.

Notwithstanding the visual confirmation of transition delay, evident in figure 5.6, a quantitative estimation of transition location is sought. Employing the technique introduced in section 3.2.1, experimentally determined transition-front displacements ($\Delta x_{tr}/c$) are extracted from IR fields for the tested plasma forcing configurations and momentum coefficients and presented in figure 5.8. The spanwise-averaged transition front location is observed to move downstream with c_{μ} for all three investigated configurations. For the UFD4_{in} configuration, the largest achieved transition delay is about $\Delta x_{tr}/c = 2.9\%$ of chord, occurring at the highest tested momentum input of $c_{\mu} = -0.068$. For the UFD4_{out} forcing configuration, the largest achieved transition delay is $\Delta x_{tr}/c = 1.7\%$ of chord which is significantly less compared to that achieved with the previous configuration. The largest displacement of the transition front with the UFD5 forcing configuration is even less and only about $\Delta x_{tr}/c = 0.8\%$ ($c_{\mu} = -0.055$).

5.5.4. Implications of applied spanwise-modulated EFD forcing

DBD plasma actuators have been successful in conditioning the swept-wing boundary layer by forcing monochromatic stationary CF instability modes in previous experimental efforts (Schuele et al., 2013; Serpieri, Yadala, et al., 2017). However, the investigated boundary layers were not conditioned with DRE forcing as done here. Therefore, in the current investigation, although the DBD actuators apply sub-critical forcing at $\lambda_z = 4$ mm and $\lambda_z = 5$ mm on the boundary layer, the spanwise wavelength of the dominant CF instability modes is still $\lambda_{z,crit} = 8$ mm as observed in figure 5.7.

Given that the forced subcritical CF instability modes are not observed in the wavelength power spectra (figure 5.7), it is probable that the amplification of the critical CF instability mode by the DRE array renders the momentum imparted by the spanwise-modulated DBD actuators as insufficient. This can result in the actuators failing to force the intended subcritical CF instability modes at sufficient amplitudes and hence, the UFD control strategy not being realized in the current study. Nevertheless, with the available data, the materialisation of the UFD control strategy due to the spanwise-modulated EFD forcing cannot be completely ruled out. This argument (UFD strategy not being realized) is further supported by the obtained results and specifically the measurable difference in the achieved transition delay, especially comparing UFD4_{in} and UFD4_{out} forcing configurations. If the observed shift in the transition front should result entirely from the UFD action by the DBD plasma actuators, the magnitude of delay achieved with these two configurations would be expected to be comparable, if not identical. Given that this is not the case, it is equitable to deduce that there are other mechanisms that collaborate with the UFD strategy to deliver the observed delays in laminar-to-turbulent transition.

It is well established that the EFD force generated by the DBD actuator is distributed such that it is highest close to the plasma-generating edge of the air-exposed electrode (electrodeinterface regions in the current actuators) owing to the strong electric field (see figure 2.7b, reproduced from Benard, Debien, et al., 2013), and diminishes as one moves away from it. This is also evident in the EFD body force distribution generated by the current DBD actuators, presented in figure 4.18. In the UFD4_{in} forcing configuration, the incoming CF vortices pass over the electrode-interface regions and hence encounter the maximum force at any operating c_{μ} . Therefore, these plasma induced jets (every second electrode-interface, see figure 5.5a) tend to directly attenuate the CF vortices as shown in the DNS by Dörr & Kloker (2016). Concurrently, the neighbouring plasma jets along the span reduce the CF velocity component in the boundary layer, thus performing an additional localised base-flow modification (BFM) action. The combined effect of these two actions stabilizes the boundary layer and hence delays laminar-to-turbulent transition.

With the UFD4_{out} configuration (see figure 5.5b), as none of the oncoming CF vortices encounter the maximum force, attenuation of the CF vortices is not expected. Thus, the produced transition delay can be attributed to just the additional BFM action of the EFD body force. The UFD5 forcing configuration, as discussed in section 5.5.1, has a variable phase relation with the incoming CF vortices (see figure 5.5c). Every fifth CF vortex passes over the electrodes' interface and hence will be attenuated. Additionally, localised BFM action will also be performed by the DBD actuator as in the previous configurations. Nevertheless, the eventual transition delay is least here among the three tested forcing configurations which is expected to be a consequence of the variable phase relation. As demonstrated by Dörr & Kloker (2016), in certain scenarios of location and orientation of the plasma body force with respect to the incoming CF vortices,



Figure 5.9.: Schematic of the forcing configurations to realize the BFM control strategy (not to scale). The swept-wing leading edge (thick black line), inviscid streamline (thin black line) and local velocity components (u_{ISL}, w_{ISL}) are shown. The DRE arrangement (•) and the DBD actuator's air-exposed electrode (grey bar) are also depicted. (a) $-F_x$ EFD plasma forcing (red arrows); (b) F_x EFD plasma forcing (blue arrows).

the applied forcing can have a detrimental effect on the boundary layer. The current variable phase relation can potentially have such an effect. A footprint of this is in fact observed in figure 5.6i where the transition to turbulence is delayed at certain spanwise regions and promoted in others.

5.6. BASE-FLOW MODIFICATION

The discussion in section 5.5 shows a delay in laminar-to-turbulent transition through the application of spanwise-modulated forcing on the swept-wing boundary layer. However, the results also indicate that with the employed spanwise-modulated actuator configuration, there is no certainty regarding the realization of instability control. Furthermore, the results also indicate that the achieved transition delay is not solely due to the UFD action of the applied EFD forcing but includes the actuator performing base-flow modification (BFM) locally. This suggests that a pure UFD action cannot be realized with the current spanwisse-modulated DBD plasma actuator geometry. On the other hand, a purely BFM control strategy can be realized with an actuator that imparts spanwise-uniform EFD forcing. The results pertaining to this, obtained from both the theoretical and experimental domains are presented in this section. It is worth noting that, unless mentioned otherwise, the presented results pertain to the flow conditions at $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1}).$

5.6.1. BFM strategy forcing configurations

In the vicinity of the leading-edge attachment line of a swept wing, the inviscid streamlines are almost aligned with the z direction. Hence, the CF component (w_{ISL}) of the base flow in this region is approximately aligned with the x direction as shown in figure 5.9. Thus, application of forcing along the x direction will largely impact the CF component in the base flow and aids in realizing the BFM control strategy. This is capitalized in the current study, through the application of a two-dimensional, spanwise uniform EFD body force in the x direction, i.e. oriented perpendicular to the leading edge. This design is adopted to avoid unwanted forcing of CF instability modes as well as to ensure a simple and robust forcing configuration. To this effect, two DBD plasma actuators were fabricated, one that applied $-F_x$ forcing (figure 5.9a) and the second to apply F_x forcing (figure 5.9b) on the boundary layer.



Figure 5.10.: Computed boundary layer velocity profiles for $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$, parallel $(u_{ISL}: \text{ dashed line})$ and perpendicular $(w_{ISL}: \text{ solid line}, \text{ magnified by factor 3})$ to the local inviscid streamline at three chord locations. Three cases shown are: baseline (black), $c_{\mu} = -0.71$ (red), $c_{\mu} = 0.71$ (blue).

It should be noted that the relative chordwise position of the DREs to the plasma actuator is not significant as long as they are positioned downstream of the first neutral point (to ensure the CF instability is formed and growing). However, the position of the plasma actuator will affect its orientation with respect to the external streamline. Aligning the plasma body force to operate normal to the inviscid streamline is expected to enhance the effectiveness of the control, as this technique is based on attenuating the CF component.

5.6.2. Laminar boundary layer solutions with EFD body force

The boundary layer solution is computed using the simplified numerical model introduced in section 3.3.1.2 in order to study the effect of the applied EFD forcing on the layer's stability. The EFD body force distribution when the actuator is operated at $V_{ac} = 4.5$ kV, $f_{ac} = 10$ kHz ($c_{\mu} = 0.71$), computed with the empirical model proposed by Maden et al. (2013) shown in figure 4.18b, is included into the x-momentum equation given in equation 3.8, to compute the respective plasma-forced laminar boundary layers. Profiles of the streamwise and spanwise velocity components (u_{ISL} and w_{ISL} , respectively) within the boundary layer at three different chordwise stations for the unforced, $-F_x$ and F_x forcing cases are shown in figure 5.10. Here, the velocity components are non-dimensionalised with the free-stream velocity. The y-coordinate is non-dimensionalised with the local boundary layer thickness δ_{99}^{uISL} , which is the distance from the surface where the streamwise velocity within the boundary layer becomes $u_{ISL}/U_{\infty} = 0.99$.

When no EFD forcing is exerted on the boundary layer, the external inviscid streamline angle is about 55° with respect to the x direction at the actuator's location of x/c = 0.035. As the applied forcing is perpendicular to the leading edge, it only partially affects CF component. F_x forcing imparts momentum partially along the cross-flow and has an enhancing effect on both the streamwise and spanwise velocity components. Contrary to this, $-F_x$ forcing opposes both velocity components, thus diminishing their respective magnitudes. This effect is clearly evident in the profiles at x/c = 0.05 shown in figure 5.10a, which is just downstream of the EFD forcing location.

Considering the low magnitude of the applied EFD body force and its localised spatial extent as shown in figure 4.18, it influences the boundary layer only locally. As such, the difference in the magnitude of the velocity components of the forced cases with respect to the unforced flow is observed to diminish downstream as seen in the profiles at x/c = 0.1 in figure 5.10b. The



Figure 5.11.: Streamwise amplification rate $(-\alpha_i \cdot c)$ (spanwise wavenumber $\beta_r = 2\pi/\lambda_z$) of stationary CF instability modes. Contour lines show the amplification factor (N = 1 (light grey), 3, 5 and 7 (black)). Dashed black line indicates the $\lambda_{z,crit} = 8 \text{ mm}$ mode. Dash-dotted vertical black line represents the chordwise station where the EFD body force is being applied. Three cases shown are: (a) Unforced flow; (b) $c_{\mu} = -0.71$; (c) $c_{\mu} = 0.71$.

profiles of the forced cases tend to collapse on the baseline flow case close to $x/c \approx 0.3$ (figure 5.10c).

5.6.3. Stability characteristics of the forced laminar boundary layers

These boundary layer solutions are now input into the LST solver introduced in sections 3.3.1 and 3.3.1.1, to assess the effect of the EFD force on the stability characteristics of the boundary layer. Stability diagrams illustrating the non-dimensionalised spatial amplification rates $(-\alpha_i \cdot c)$ and amplification N-factors along the streamwise direction of different stationary CF instability modes (represented by the spanwise wavenumbers $\beta_r = 2\pi/\lambda_z$) of the cases being considered here are presented in figure 5.11.

Stability diagram of the unforced boundary layer shown in figure 5.11a confirms that the $\lambda_{z,crit} = 8 \text{ mm}$ mode (dashed line in figure), chosen as the mode forced by DRE in the experiment, is one of the most unstable stationary CF instability modes at these flow conditions. Comparing the stability plots of the the forced cases with that of the unforced flow confirms that $-F_x$ forcing (figure 5.11b) hinders the growth of stationary CF instability modes, thus stabilising the boundary layer. As expected, F_x forcing (figure 5.11c) augments the growth rate of all stationary CF instability modes and has a destabilising effect on the boundary layer. In both cases, the effect of forcing is rather localized as observed in the boundary layer velocity profiles in figure 5.10. Indeed, upstream of $x/c \approx 0.3$, the growth rates show insignificant disparity, as suggested by the almost identical velocity profiles at that station (figure 5.10c).

The amplification N-factors, represented by the contour lines in figure 5.11, reveal the integral effect of the exerted forcing on the stability characteristics of the boundary layer. Although the spatial extent of the applied force and its modification of the boundary layer is localized, the N-factors show that its effect on flow stability is evident well downstream of the forcing location. Assuming that the critical N factor does not significantly change due to actuation, the transition location can be expected to shift downstream with $-F_x$ forcing and upstream with F_x forcing.



Figure 5.12.: IR thermography mean fields $(Re_c = 2.1 \cdot 10^6, U_{\infty} = 25 \text{ m s}^{-1})$. (a) Unforced flow. (b) $-F_x$ $(c_{\mu} = -0.71)$. (c) Subtraction of (a) from (b). (d) Unforced flow. (e) F_x $(c_{\mu} = 0.71)$. (f) Subtraction of (d) from (e). The solid white lines represent constant chord positions.

5.6.4. Boundary layer and transition topology through IR thermography

The flow topology, visualised through IR thermography imaging, is shown in figure 5.12 for $Re_c = 2.1 \cdot 10^6 \ (U_{\infty} = 25 \text{ m s}^{-1})$. IR images of the unforced cases are shown in figures 5.12a and d. As in section 5.5.2, the DBD plasma actuators configured to apply the necessary forcing are installed on the wing but not operative here as well. As seen in figure 5.6 (and in figure 5.2b), the laminar breakdown occurs along the characteristic jagged pattern reported for stationary CF instability induced transition (Saric, Reed, et al., 2003; Saric, Carpenter, et al., 2011; Downs & White, 2013; Serpieri & Kotsonis, 2016), at approximately the same chordwise position ($x/c \approx 0.32$). This once again confirms the effectiveness of the proposed fabrication technique in generating micron-thick electrodes which do not influence/interfere with the evolution of the boundary layer.

The wavelength spectra along constant chord lines within the regions enclosed by the dashed (white) polygons in the IR images of the two unforced flow cases, shown in figures 5.13a and c, confirm that the $\lambda_{z,crit} = 8$ mm mode forced by the DREs is indeed the spanwise dominant mode in this boundary layer as well.

The flow arrangement when $-F_x$ and F_x EFD forcing is applied on the boundary layer is



Figure 5.13.: Wavelength power spectral density of IR pixel intensity along constant chord lines. The analysis is performed in the regions enclosed by the dashed (white) polygons in figure 5.12. (a) Unforced flow. (b) $-F_x$ forcing ($c_\mu = -0.71$). (c) Unforced flow. (d) F_x forcing ($c_\mu = 0.71$).

shown in figures 5.12b and e respectively. A bright region is noticed in these fields close to the swept-wing leading edge here as well. This is the consequence of the thermal radiation due to the plasma discharge from the respective plasma actuators being operated. A streaky flow arrangement is observed in these fields as well. Spectral analysis again confirms that the spanwise dominant mode in the boundary layer is the $\lambda_{z,crit} = 8$ mm mode forced by the DREs (figures 5.13b and d). This suggests that, as expected, the applied spanwise-uniform EFD forcing has no effect on the dominant CF instability modes within the boundary layer. In addition, to facilitate visualisation of the movement of the transition front, figure 5.12c and f show the subtraction of the IR field with EFD forcing from the corresponding unforced case. F_x forcing evidently shifts the transition front upstream, compared to the unforced flow. Instead, $-F_x$ forcing, at the same forcing and flow conditions, shifts the transition front downstream.

5.6.5. Location of laminar-to-turbulent transition

Experimentally measured transition locations extracted from the time-averaged IR fields according to the procedure introduced in section 3.2.1 are presented in figure 5.14, for both $-F_x$ and F_x plasma forcing. Transition delay or advancement is proportional to the magnitude of c_{μ} . Within the tested cases, the largest transition delay is approximately 4.5%, occurring at the lowest Reynolds number and highest momentum coefficient tested ($Re_c = 2.1 \cdot 10^6$, $c_{\mu} = -0.71$).

At higher Reynolds numbers, the displacement of the transition fronts with c_{μ} shows similar trends. For $Re_c = 2.3 \cdot 10^6$ ($U_{\infty} = 27.5 \text{ m s}^{-1}$, \triangle), $c_{\mu} = -0.62$ and $Re_c = 2.5 \cdot 10^6$ ($U_{\infty} = 30 \text{ m s}^{-1}$, \Box), $c_{\mu} = -0.55$ the transition delay is 3.7% and 1.6%, respectively. The continuity in increasing the transition delay with $-c_{\mu}$ is well-defined, however, the percentage of delay decreases. Effectively, the current control technique is a direct approach against the weak CF velocity component. This component scales inviscidly with the free-stream velocity and thus, the authority of the applied EFD body force on the boundary layer reduces as Reynolds number increases (see variation of c_{μ} in figure 5.4a). Furthermore, the critical mode in the boundary layer changes with the Reynolds number, however the applied forcing is observed to still have an effect. This further demonstrates the inherent insensitivity of the base-flow modification



Figure 5.14.: Experimentally measured transition locations (x_{tr}/c) versus momentum coefficient c_{μ} . \bigcirc : $Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$; \triangle : $Re_c = 2.3 \cdot 10^6$, $U_{\infty} = 27.5 \text{ m s}^{-1}$; \Box : $Re_c = 2.5 \cdot 10^6$, $U_{\infty} = 30 \text{ m}$ s⁻¹; •: transition location predicted by simplified numerical model for $Re_c = 2.1 \cdot 10^6$, $U_{\infty} = 25 \text{ m s}^{-1}$.

strategy to the wavenumber of the stationary CF instability modes.

The uncertainty bands in figure 5.14 represent the width of the transition-front, estimated by the standard error of the linear fit through the jagged transition-front (see figure 3.6d). The transition-front width observed here appears rather insensitive to the plasma forcing. A possible reduction in the transition-front width would suggest the amplification of travelling CF instabilities in the boundary layer, as these unsteady modes erode the characteristic jagged transition front and result in a more spanwise uniform transition-front (Downs & White, 2013). The transition front maintaining the jagged shape suggests that no fluctuations in the frequency band of travelling CF instabilities were introduced. This, along with the discussions in section 5.5, serve as a preliminary indication that applying plasma EFD forcing at such high frequencies will not lead to the amplification of unsteady fluctuations in such boundary layers.

5.6.6. Estimation of transition location through simplified numerical model

The evolution of N factors for the $\lambda_{z,crit} = 8$ mm stationary CF instability mode along the chord for the unforced and the two forced cases, computed by the simplified model, is presented in figure 5.15. While the existence of other stationary modes cannot be excluded, the 8 mm mode is the most dominant and observed to drive transition (figure 5.13), as it is conditioned by the DREs. Hence, the N factor of only this mode is presented and analysed here. Using the transition locations obtained from IR thermography measurements, the critical N factor at which transition occurs for the unforced flow case is observed to be $N_{crit} \approx 5.6$ (dash-dotted line in figure 5.15).

In compliance with the stability diagrams presented in figure 5.11, although the applied forcing affects the amplification of instabilities locally, its integral effect results in a global response on the N factors. The reduction of instability mode growth rates with $-F_x$ forcing leads to a global decrease of N-factors compared to the unforced flow. Thus, the critical value N_{crit} at which transition naturally occurs in the unforced case is reached downstream, thus delaying transition. Contrarily, F_x forcing increases the amplification of instabilities and results in higher



Figure 5.15.: N-factors of the $\lambda_{z,crit} = 8$ mm stationary CF instability mode in unforced and forced numerical boundary layers. Chordwise locations of the DRE array (dotted line) and DBD plasma actuator (dashed line) are depicted. Experimentally obtained transition locations (\Box) are also shown. Unforced boundary layer $N_{crit} = 5.6$ (dash-dotted line).

N-factors when compared to the unforced flow. This increase in *N* factors results in arriving at N_{crit} upstream of the unforced case, thus promoting laminar-to-turbulent transition. The experimentally detected transition locations of the forced flow cases are also shown in figure 5.15, indicating the ability of the simplified modelling approach in capturing the overall trends of transition manipulation with forcing.

The transition location for different c_{μ} for the flow with Reynolds number $Re_c = 2.1 \cdot 10^6$, predicted by the simplified numerical model in the aforementioned manner is also shown in figure 5.14. It is evident that the model is able to capture the experimental trend, however fails to predict the transition location accurately. This is expected, considering the range of assumptions and simplifications entering the chain of modelling steps (see section 3.3.1.2). Notwithstanding the caveats brought forth, the simplified modelling approach presents a robust, fast and efficient tool, ideally suited for augmenting the experimental design process and provide basic, back-ofenvelope scaling of the control performance.

5.7. COMPARISON OF SPANWISE-MODULATED AND UNIFORM FORCING THROUGH CONTROL EFFICIENCY

The discussion until now has demonstrated successful delay in laminar-to-turbulent transition with both spanwise-modulated and spanwise-uniform forcing, as long as the EFD body force is oriented such that it acts against the incoming flow. This corroborates well with the observations of Dörr & Kloker (2016). Having tested both forcing methods under similar flow and boundary layer conditions, a comparison between the two can be performed, based on the achieved displacement of laminar-to-turbulent transition. The test cases pertaining to the flow conditions at $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$ are considered here. The transition-front displacement realized through all three forcing configurations involving spanwise-modulated actuators (UFD4_{in}, UFD4_{out} and UFD5) and the corresponding spanwise-uniform forcing cases that perform baseflow modification through the application of a EFD body force along the -x direction ($-F_x$ forcing case, $c_{\mu} = -0.05$, -0.07, -0.13 and -0.28) are considered for this comparison.

The achieved transition-front displacement as a function of the average power delivered to the different DBD plasma actuators is presented in figure 5.16. Here, the delivered power is



Figure 5.16.: Experimentally determined laminar-to-turbulent transition-front displacement $(\Delta x_{tr}/c)$ versus normalized average power delivered (P_{avg}/z_{PA}) to the corresponding DBD plasma actuators. The investigated flow conditions correspond to $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$.

normalized with z_{PA} (see figure 4.1), the spanwise length over which the corresponding DBD actuator exercises boundary-layer control, as done in figure 4.20a. This figure demonstrates that **spanwise-modulated forcing has the potential to realize a comparable delay** in transition as spanwise-uniform forcing, while requiring significantly less input power. This becomes more apparent when two specific cases corresponding to the UFD4_{in} and BFM $(-F_x)$ configurations are compared. For the test case of UFD4_{in} configuration with the highest input power of $P_{avg}/z_{PA} = 2.63$ W m⁻¹ ($c_{\mu} = -0.068$), the average transition delay achieved is $\Delta x_{tr}/c = 2.9\%$. A comparable transition delay is achieved through the spanwise-uniform configuration, with the delivered power in this case being $P_{avg}/z_{PA} = 6.67$ W m⁻¹ ($c_{\mu} = -0.13$). This is about 2.5 times that delivered to the spanwise-modulated actuator for the UFD4_{in} forcing configuration. Thus, spanwise-modulated forcing can substantially improve the efficiency of the control system.

The achieved transition delay will lead to a reduction of skin-friction drag, which is significantly lower for a surface with laminar flow as compared to one with turbulent flow. Comparing the skin-friction drag reduction with the power delivered to the control system can provide the net-gain realized through the applied boundary-layer control, an estimate of which is made here. To this effect, a flat-plate boundary layer with zero-pressure-gradient is considered here as a first-order approximation, the chord-length of this flat-plate being equated to the streamwise chord of the swept-wing model c = 1.27 m.

The skin-friction coefficient depends on the Reynolds number, Mach number and surface roughness. For the surface with a laminar boundary layer, the skin-friction coefficient ($C_{f,laminar}$) is computed by equation 5.2a, which is in fact integrated form of the local skin-friction coefficient from the Blasius solution of a laminar boundary layer. For a turbulent boundary layer ($C_{f,turbulent}$) is computed through equation 5.2b (Raymer, 2018, equation 12.27).

$$C_{f,lamimar} = \frac{1.328}{\sqrt{Re_c}} \tag{5.2a}$$

$$C_{f,turbulent} = \frac{0.455}{(\log_{10} Re_c)^{2.58} \cdot (1 + 0.144M_{\infty}^2)^{0.65}}$$
(5.2b)

It is worth noting that the Mach number correction (second) term in the denominator of equation 5.2b goes to 1 for low sub-sonic flow conditions as is in the current experiments. Thus, the skin-friction coefficient computed here is dependent only on the chord Reynolds number Re_c . For the investigated flow conditions at $Re_c = 2.1 \cdot 10^6$, these parameters are $C_{f,laminar} = 9.16 \cdot 10^{-4}$ and $C_{f,turbulent} = 3.9 \cdot 10^{-3}$. Thus, the skin-friction drag due to turbulent flow is about 4 times that due to laminar flow.

The overall skin-friction coefficient C_f can then be computed through a weighted-sum according to equation 5.3,

$$C_f = \left(\frac{x_{tr}}{c} + \frac{\Delta x_{tr}}{c}\right) \cdot C_{f,laminar} + \left(1 - \left(\frac{x_{tr}}{c} + \frac{\Delta x_{tr}}{c}\right)\right) \cdot C_{f,turbulent}$$
(5.3)

where, x_{tr}/c is the transition location of the unforced boundary layer and $\Delta x_{tr}/c$ is the achieved transition displacement when boundary-layer control is applied. The change in skin-friction drag is obtained by subtracting the corresponding coefficient of the forced case from the unforced one, and is as expressed in equation 5.4.

$$\Delta C_f = \frac{\Delta x_{tr}}{c} \cdot C_{f,laminar} - \frac{\Delta x_{tr}}{c} \cdot C_{f,turbulent}$$
(5.4)

Thus, through this method, the change in the overall skin-friction coefficient can be estimated by just knowing the displacement of the transition front.

The net gain η realized through the applied boundary-layer flow control is then evaluated as the ratio between the power saved due to the reduction in skin-friction drag and the power delivered to the DBD plasma actuator normalized with the spanwise-extent of the actuator (P_{avg}/z_{PA}) . The power saved per unit span is equated to the product of change in skin-friction coefficient ΔC_f , the dynamic pressure $\rho U_{\infty}^2/2$, the streamwise chord c and the free-stream velocity U_{∞} , as shown in equation 5.5. It is worth noting that a negation of ΔC_f is done here since a reduction in the skin-friction drag represents a gain in power.

$$\eta = \frac{(-\Delta C_f) \cdot \frac{\rho U_{\infty}^2}{2} \cdot c \cdot U_{\infty}}{(P_{avg}/z_{PA})}$$
(5.5)

The estimate of the net-gain realized through the different forcing configurations investigated here is presented in figure 5.17. The spanwise-uniform forcing configuration realizing the BFM control strategy, while being robust in achieving a transition delay, is not very efficient. The maximum gain for this configuration is only $\eta = 0.4$ ($P_{avg}/z_{PA} = 1.28$ W m⁻¹, $c_{\mu} = -0.07$). In fact, for the case with maximum input power of $P_{avg}/z_{PA} = 10.5$ W m⁻¹ ($c_{\mu} = -0.27$), which resulted in the highest transition delay among all test cases considered for this comparison ($\Delta x_{tr}/c = 4.5\%$), the net gain is a mere $\eta = 0.16$. Thus, the power consumed to realize the necessary control is more than that gained through the skin-friction drag reduction due to the achieved transition delay.

The spanwise-modulated forcing configurations demonstrate a significant improvement in the overall system efficiency. Especially with the UFD4_{in} configurations, a maximum gain of $\eta = 6.9$ is realized, when the actuator is operated at a low input power. The low



Figure 5.17.: Net-gain estimation due to the achieved transition delay. The dotted (black) line represents a gain of $\eta = 1$. The investigated flow conditions correspond to $Re_c = 2.1 \cdot 10^6 (U_{\infty} = 25 \text{ m s}^{-1})$.

input-power cases of the UFD5 configuration too result in a gain greater than $\eta = 1$. However, with both these configurations the gain decreases with increasing input power and becomes less than $\eta = 1$, even if the achieved transition delay continues to increase. The UFD4_{out} configuration, while is not as efficient as the other spanwise-modulated actuators, it is more efficient than the spanwise-uniform forcing case. This outcome of the current estimation suggests that even a small delay in transition can result in a considerable improvement in the efficiency of the overall system. This is especially useful at higher Reynolds number, where the authority of the control system while being marginal, can result in improvements to the system efficiency.

One must bear in mind that this net-gain estimation is rather qualitative and extension to the swept-wing boundary layer is not straightforward. As such, the current estimation is based on the assumption of a zero-pressure-gradient, which is not the case in the swept-wing boundary layer. Moreover, the weight, complexity and robustness of the plasma control system as well as the influence of the applied boundary-layer control on other components of the total drag is not accounted for here. Indeed, thorough experimental investigations are necessary to obtain a quantitative measure of the total drag.

5.8. CONCLUDING REMARKS AND PERSPECTIVES FOR THE FUTURE

The present investigation provides the first experimental proof of transition delay in threedimensional swept-wing boundary layers that is dominated by the cross-flow (CF) instability using DBD plasma actuators. The upstream flow deformation (UFD) and base-flow modification (BFM) control strategies were realized by designing actuators that exert spanwise-modulated and spanwise-uniform EFD forcing respectively, on the boundary layer. A comparison is made between these two control strategies and finally the collaboration of these two control methods are investigated upon. Finally, an array of discrete roughness elements (DRE) installed just upstream of the DBD actuator conditioned the boundary layer with the critical stationary CF instability mode which provided a "worst case scenario" to investigate the effectiveness of the applied control.

The UFD control strategy shows favourable results. By manipulating the relative position of these spanwise-modulated DBD actuators with respect to the DREs, three forcing configurations

are investigated. The differences in the achieved transition delay with these three configurations suggests that the relative position of the plasma forcing with the incoming CF vortices affects the performance of the control technique. While the UFD action of these spanwise-modulated actuators are not clearly observed, a combination of direct attenuation of stationary CF vortices and localised BFM action seem to be controlling the boundary layer and thus its transition to turbulence.

The presented results also provide the first experimental proof-of-concept of plasma-based base-flow modification for CF instability control. The experimental results and the numerical simulation using a simplified numerical model based on LST show that application of forcing against the CF component results in the stabilisation of the boundary layer and leads to transition delay. Nevertheless, several key-points need to be still addressed for the successful establishment of BFM as a viable technique for the control of swept-wing boundary layers. Furthermore, the simplified numerical model, although captured the trends and order of magnitude for transition delay/advancement satisfactorily, ample room for improvement remains for such models, mainly in incorporating non-parallel and non-linear theories. Notwithstanding the level of approximation, such simple, fast and robust models are indispensable for the design of this type if investigations.

Cross-flow-dominated transition is known to be extremely sensitive to surface roughness effects (Saric, Reed, et al., 2003; Saric, Carpenter, et al., 2011; Saric, West, et al., 2015). The DRE/UFD approach has shown promising results for control of CF-dominated transition, however also revealed the strong sensitivity of this technique to a number of parameters such as free-stream turbulence and background roughness (Saric, West, et al., 2015; Lovig et al., 2014). The base-flow modification strategy can be expected to be more robust to these effects, especially considering its effectiveness at the three different flow conditions tested. However, the acquired results show that spanwise-uniform forcing that performs BFM, requires more energy to realize control of transition rendering the control system as less efficient. Comparatively, spanwise-modulated forcing consumes significantly less energy and has the potential to improve the overall efficiency of the aerodynamic body.

Perspectives for the future

In the experiments of Serpieri, Yadala, et al. (2017) plasma actuators were found to introduce velocity fluctuations in the band of travelling cross-flow instabilities. In the present study, the issue was circumvented by high AC driving frequency ($f_{ac} = 10$ kHz). However, upscaling this technique to higher Reynolds numbers, will scale the frequency of instabilities as well. This will effectively require a further increase of the AC frequency, which will in turn challenge the current plasma generation approach. In addition, possible receptivity mechanisms between the plasma formation and incoming disturbances (i.e. free-stream turbulence) could give rise to generation of secondary instabilities either through non-linear interactions or transient growth. Dedicated studies of these mechanisms should be pursued to better understand receptivity of the boundary layer to the control being applied.

Finally, a key point is the future upscaling of the plasma forcing to higher Reynolds numbers. In the present study, the employed actuator was relatively weak. Higher overall momentum coefficients can be achieved by increasing the dielectric thickness, applied voltage and number of successive actuators.

6

PLANE MIXING LAYER

The plane mixing layer generated by two flow streams with different free-stream velocities interacting with one another downstream of the splitter-plate trailing edge is investigated in this chapter. Previous research efforts which were briefly documented in section 2.3 have demonstrated that this flow configuration is representative of a free shear flow regime and is dominated by two instability mechanisms. The primary instability of *Kelvin-Helmholtz* type governs the evolution of large-scale spanwise vortical flow structures and the control of this instability is frequency-dependent. The growth of the mixing layer is a product of the interaction of these large-scale vortices. Based on the presence of sub-harmonic instabilities and its relative phase with respect to the fundamental KH instability, these vortices amalgamate into one vortex, resulting in an increase in the width of the layer. The secondary instability manifests itself as streamwise vortices spread along the span. These vortices loop around the primary vortex cores and when the latter undergo pairing, these streamwise vortices are engulfed into the the newly formed primary vortex. This eventually causes the onset of small-scale turbulence which aids in mixing enhancement. Given that these vortices are spread along the span of the mixing layer, these instabilities can be manipulated through spatial forcing applied along the span.

In this thesis, the primary instabilities are controlled with a DBD plasma actuator exerting a spanwise-uniform EFD body force. The frequency response of these actuators is taken advantage of to excite these instabilities and manipulate the stability characteristics of the mixing layer flow field. The uniqueness of this forcing needs to be highlighted here. The EFD forcing exerted does not have a strong vertical component and the imparted perturbation is confined near the dielectric surface. In the current study, this feature is exploited to investigate the impact of exciting different primary KH instabilities in the mixing layer, by imparting wellcontrolled streamwise perturbations into the splitter-plate boundary layer ensuing on the highvelocity side. Thus the necessary instability control is realized here in an indirect manner, by manipulating the shear and spectral content within the incoming boundary layer. Largely horizontal forcing emulates the effect of boundary layer instabilities (i.e. Tollmien-Schlichting waves) developing upstream of the splitter-plate trailing edge. Thus the chosen forcing technique provides a more "realistic" forcing environment than other techniques and replicates the features of perturbations in the natural case, albeit at a well-controlled frequency and amplitude. This is followed by exciting secondary instability modes with a spanwise-modulated actuator design. The deployment of DBD actuators for this purpose provides another unique scenario. While the spanwise-modulated forcing is achieved through the actuator's geometry, the electrical signal input will be used to impart fluctuations corresponding to primary KH instabilities. The effect of imparting these fluctuations in a three-dimensional manner on the mixing layer will also be

investigated.

This chapter is organised as follows. The experimental flow conditions are introduced in section 6.1 and information regarding the flow velocity measurement is presented in section 6.2. Given that the first objective of this thesis corresponding to the plane mixing layer was to assemble a set-up to generate this flow configuration, thorough characterisation of the baseline (unforced) flow holds utmost importance. This firstly aids in comparing the current results with past experimental efforts. Secondly, to study the impact of the applied instability control effectively, a thorough understanding of the unforced case is key. Results pertaining to this is reported in section 6.3. The analysis of the unforced flow includes the features of the mean and fluctuating quantities, the layer's growth, linear stability analysis and power spectral density to identify the fundamental KH instability and finally, the spatial organisation of different fluctuations through spectral proper orthogonal decomposition (SPOD).

Section 6.4 is comprised of the results and discussions of the mixing layer forced with the spanwise-uniform actuator to manipulate the primary KH instability modes. The methodology followed to investigate theses forced mixing layer are as follows. Firstly, the operating and forcing conditions of the employed actuators are reported. The mean-flow features and growth of the mixing layer is investigated and compared to the unforced mixing layer and other forced mixing layers. The effect of the applied forcing on the spectral content of the ensuing layer is then investigated which is followed by a detailed discussion on the spatial organisation of the different fluctuations obtained through SPOD. Finally, select SPOD modes are used to reconstruct the velocity time-series and the dynamics of the large-scale vortices are investigated. Results pertaining to the excitation of secondary instability modes through spanwise-modulated actuators are presented in section 6.5. The conclusions drawn and the perspectives for future research are summarised in section 6.6.

6.1. EXPERIMENTAL FLOW CONDITIONS

The experimental investigations concerning instability control in a plane mixing layer flow configuration was performed at Institut PPrime, Université de Poitiers, France. The closed-loop wind-tunnel employed for this investigation is as described in section 3.1.2. The plane mixing layer configuration was generated with two air-flow streams at different velocities separated by a splitter-plate, interacting with each other at its trailing edge. The velocity of the bottom stream was $U_1 = 5 \text{ m s}^{-1}$ and that of the top stream was held constant at $U_2 = 10 \text{ m s}^{-1}$. This resulted in a velocity ratio of $r = U_1/U_2 = 0.5$. Two other velocity parameters are utilized in this article which are introduced here. The first is the velocity difference between the two streams $\Delta U = U_2 - U_1 = 5 \text{ m s}^{-1}$. The second parameter is the average of the two free stream velocities $\overline{U} = (U_2 + U_1)/2 = 7.5 \text{ m s}^{-1}$ which is expected to be the convection speed of the spanwise vortices in the ensuing mixing layer (Oster, Wygnanski, & Fiedler, 1977). With these velocity parameters, a second velocity ratio is generally defined as $R = \Delta U/(2\overline{U})$ which in the current experiment is R = 0.33. At these flow conditions, the turbulence intensity measured using hot-wire anemometry (single-wire probe, bandpass filtered between 10 and 1000 Hz) in the low and high-velocity streams are $Tu_1/\overline{U} \approx 0.9\%$ and $Tu_2/\overline{U} \approx 0.3\%$ respectively. Finally, it is worth recalling that both the splitter-plate boundary layers are not tripped during all the reported test cases.

Recalling the coordinate system used in this study (see section 3.1.3.3), it is represented by



Figure 6.1.: Statistical convergence of different components of velocity in the acquired PIV vector fields of the plane mixing layer, at different streamwise stations along y = 0.

(x, y, z) where x is along the free-stream direction, y is perpendicular to the plane of the mixing layer and z is along the span. The corresponding velocity components are represented by u vand w respectively. The geometrical origin of the mixing layer is fixed at the trailing-edge of the splitter plate (see figure 3.4b), with z = 0 being at the centre of the splitter plate along the span (15 cm from both walls).

6.2. FLOW VELOCITY MEASUREMENT THROUGH TIME-RESOLVED PLANAR PIV

Previous experimental studies on the mixing layer flow configuration have mainly employed hotwire anemometry (HWA), which is a point-field measurement, to obtain the flow velocity. Some more recent studies have made use of planar PIV, however these were mostly not time-resolved (see for example Mathis et al., 2007; Ely & Little, 2013; Parezanovic et al., 2014; Singh & Little, 2020). While these studies have contributed a great deal to the current physical understanding of the stability and dynamics of the mixing layer, pertinent features of the flow field remain unresolved due to the inherent limitations of these measurement techniques. These are essentially the functional relationship between the mean flow and development of instabilities in both space and time. The importance of such spatio-temporal information cannot be understated as it has the potential of providing a closure between the topology of the mean flow, its ensuing stability and the development of instabilities such as KH rollers, which in turn affect the mean flow through non-linear interactions such as vortex merging. To this end, flow velocity was measured using time-resolved planar PIV, the instrumentation and set-up of which is as described in section 3.2.2.4. The inspected plane was along the x - y directions. The acquisition rate for all testes cases was fixed at $f_s = 9.5$ kHz (in single frame mode) and the cross-correlation resulted in a final vector spacing of 1 mm in both the x and y directions.

Firstly, statistical convergence of different velocity components was checked. To this effect, convergence tests were carried out at different points in the PIV FOV, results of which are presented in figure 6.1. For the horizontal component of velocity u, both the mean and root-mean-square values converge within 2000 instantaneous fields (210 ms) in most of the imaged spatial domain as shown in figures 6.1a and b. Far downstream however, statistical convergence is achieved with approximately 4000 frames. With respect to the vertical velocity component v, both the mean and root-mean square values converge in about 1000 instantaneous fields (105 ms) in the entire measured domain as shown in figures 6.1c and d. Thus, a total of 20,000 frames were acquired to ensure that both the mean and fluctuating quantities are converged, and so are the different derived quantities. The resulting measurement time is 2.1 s per test case.

An estimate of the error in measurements based on the methods introduced in section 3.2.2.5 is now reported. In the current experiment, typical particle displacement was approximately 6-8pixels as reported previously in section 3.2.2.4. Therefore, the relative error of the measured velocity (instantaneous and/or fluctuating fields) is approximated to be $\epsilon_{u(t),v(t)} (= 0.1/6) \approx$ 2%. The uncertainty on the time-averaged velocity fields is computed using the linear error propagation technique (Sciacchitano & Wieneke, 2016) according to equation 3.2. The maximum uncertainty on the time-average fields in the current experiment is $\epsilon_{\overline{u}} \approx 4\%$ and $\epsilon_{\overline{v}} \approx 3\%$.

6.3. UNFORCED MIXING LAYER

Designing and modifying an existing wind tunnel to accommodate mixing layer studies was carried out during the course of this thesis. Therefore, before the application of any flow control and analysis of the corresponding effects, a thorough characterisation of the baseline (no EFD forcing) flow configuration holds utmost importance. This was in fact one of the objectives of this thesis and results corresponding to this exercise is reported first. The velocity fields of this flow case were measured using PIV at mid-span (z = 0). It is worth noting that while acquiring the PIV fields of this baseline, the PMMA attachment housing the DBD actuator designed to apply spanwise-uniform EFD forcing (see section 4.5) formed the trailing edge of the splitter-plate. Therefore, this DBD plasma actuator (especially its air-exposed electrode) is mounted on the splitter plate. Thus, this case is henceforth referred to as unforced mixing layer (UML).

Initially, the constancy of the free-stream velocities in the measured spatial domain was checked using the streamwise velocity component at the top and bottom edges of the PIV FOV. The velocity of the bottom stream was $U_1 = 5 \text{ m s}^{-1}$ and varied by a maximum of 3%. The maximum variation in the velocity of the top stream was 1.7% compared to the expected value of $U_2 = 10 \text{ m s}^{-1}$. These values of variations from the expected velocity in the free-streams are lower than the uncertainty of the current PIV measurement reported in section 6.2.

6.3.1. Incoming flow conditions through boundary layer measurements

As indicated previously, both the splitter-plate boundary layers, on the high-velocity and low-velocity sides, were not tripped. The mean boundary layer profiles just upstream of the splitter-



Figure 6.2.: Splitter-plate boundary layer profiles at x = -1 mm. (a) High-velocity side boundary layer; (b) Low-velocity side boundary layer. Here, y = 0 is the splitter-plate surface. The δ_{99} location shown is that of the measured profile.

plate trailing edge (at x = -1 mm) acquired through PIV, are presented in figure 6.2. Also shown in this figure is the Blasius profile fit on the experimentally acquired profiles and boundary layer thickness (δ_{99}) computed through this profile. Both high-velocity (6.2a) and low-velocity (6.2b) side profiles are well fit by the Blasius velocity profile suggesting that both boundary layers are laminar. The shape factors computed with the velocity data from the splitter-plate surface until $y = 2 \cdot \delta_{99}$ are $H_1 \approx 2.8$ and $H_2 \approx 2.9$ confirming that both boundary layers are in laminar conditions. It must however be indicated here that due to the laser sheet being introduced from the top of the wind-tunnel test section (see figure 3.9), a shadow is produced by the splitter plate on its bottom side. Thus, the computation of vectors on the low-velocity (bottom) side of the splitter plate and the reported boundary-layer parameters are more prone to errors.

The momentum thickness of the splitter-plate boundary layer ensuing on the high-velocity side, at this location (x = -1 mm) was computed to be $\theta_i = 0.55$ mm. This value is utilized henceforth to non-dimensionalise the spatial coordinates and other length parameters wherever necessary. The Reynolds number based on the initial momentum thickness and kinematic viscosity ν is computed by $Re_{\theta_i} = \overline{U}\theta_i/\nu$. For the investigated flow conditions where the average velocity between the two streams is $\overline{U} = 7.5$ m s⁻¹ and assuming the kinematic viscosity to be $\nu = 1.47 \cdot 10^{-5}$ m² s⁻¹ (of air at 15° C), the Reynolds number is found to be $Re_{\theta_i} = 280$.

6.3.2. Time-averaged flow topology

The normalised, time-averaged u-velocity component (\overline{u}_N) is presented in figure 6.3. The normalisation of the mean field is performed according to equation 6.1 given below.

$$\overline{u}_N = \frac{\overline{u} - U_1}{\Delta U} \tag{6.1}$$

It is worth noting here that all velocities presented henceforth are non-dimensionalised with ΔU . The \overline{u}_N velocity presented in figure 6.3 contains both the near-wake and far-wake PIV FOVs. The mixing layer developed due to the interaction of the high and low-velocity flow streams, downstream of the splitter-plate trailing edge, is clearly visible. Also shown in this figure is the



Figure 6.3.: Time-averaged, normalised \overline{u} -velocity (\overline{u}_N) of the unforced mixing layer. The dashed (black) line represents $y_{0.5}$.



Figure 6.4.: Thickness of the unforced mixing layer along the streamwise direction.

 $y_{0.5}$ trace (dashed black line) which represents the y-coordinate where $\overline{u}_N = 0.5$.

The splitter plate employed for this study has a 3 mm thick trailing edge as discussed in section 3.1.2 (see figure 3.4b). This results in a wake effect that manifests as the formation of a recirculation region just downstream of the splitter-plate trailing edge as also observed by Laizet et al. (2010) who performed direct numerical simulation (DNS) to study the effect of such splitter-plate trailing edge shape (among others) on the spatial development of the mixing layer. However, it is good to note that the thickness of the current trailing edge is much lower than the case with a similar splitter-plate shape in their DNS study. In the current unforced mixing layer the recirculation region is observed to last until about $x/\theta_i = 44$ ($x \approx 24$ mm). Beyond this streamwise location, the occurrence of the mixing layer is observed.

The vorticity thickness, computed according to equation 2.6, of the unforced mixing layer along the streamwise direction is presented in figure 6.4. Until $x/\theta_i \approx 30$ no growth is observed due to the occurrence of the recirculation region. In the streamwise region, the vorticity thickness measures $\delta_{\omega} \approx 3$, which in dimensional terms ($\delta_{\omega} = 1.65$ mm) is approximately half the thickness of the splitter-plate trailing edge. An increase in the vorticity thickness is observed just downstream of this location which becomes linear after $x/\theta_i \approx 40$. This pertains to the laminar region where the linear instability causes the roll-up of the shear layer at the fundamental (or naturally preferred) frequency, resulting in the formation of spanwise-oriented vortices (Michalke, 1965a; Winant & Browand, 1974). This is followed by a linear growth from $x/\theta_i \approx 100$ which can be attributed to vortex pairing occurring stochastically in space and time (Winant & Browand, 1974; Oster & Wygnanski, 1982; Ho & Huang, 1982). It is also noteworthy



Figure 6.5.: Profiles of \overline{u}_N -velocity at different streamwise locations and the corresponding tanh fit.

that the growth pattern of the current mixing layer resembles that of the mixing layer ensuing from a thin trailing edge in the numerical work of Laizet et al. (2010).

The spreading rate of the mixing layer can be approximated by the Abramovich (1963)-Sabin (1965) rule given by equation 2.7. In the linear growth region of the current mixing layer, the overall spreading rate is well approximated by $d\delta_{\omega}/dx = 0.25 \cdot R$ (dashed red line in figure 6.4). The value of the fitting parameter c = 0.25 is greater compared to the generally accepted range of c = 0.16 to 0.18 (Brown & Roshko, 1974; Browand & Latigo, 1979; Mehta, 1991). The increase in spreading rate can be attributed to the higher level of free-stream turbulence in the low velocity stream, which results in an increased entrainment and hence, growth rate (Pui & Gartshore, 1979). However, the complex interaction of the splitter-plate wake with the mixing layer cannot be neglected either. The value of the fitting parameter c is lower than that reported by Laizet et al. (2010) (c = 0.34) but similar to that observed in the experiments of Huang & Ho (1990) (approximated to be 0.27).

Also shown in figure 6.4 is the mixing layer integral thickness (often referred to as momentum thickness θ) computed by equation 2.5. In the turbulent region of the current unforced mixing layer, the integral thickness increases linearly along the streamwise direction as well. The growth rate measured corresponding to the integral thickness is $d\theta/dx \approx 0.017$ (= $0.05 \cdot R$). This growth rate is comparable to that observed by Ho & Huang (1982) ($d\theta/dx \approx 0.013$) in their mixing layer with a velocity ratio (R = 0.31) similar to the current study.

Profiles of the normalised mean streamwise velocity (\overline{u}_N) at different streamwise locations are shown in figure 6.5. The hyperbolic tangent profile is given by the equation 3.9. The *tanh* fit on the velocity profiles is also presented in figure 6.5 (red line). It is worth noting that the *y*-coordinate in this plot is shifted with the location of $y_{0.5}$ at the particular streamwise location and is non-dimensionalised with the local vorticity thickness $(\delta_{\omega}(x))$.

The wake effect of the thick splitter-plate trailing edge manifests itself as a velocity defect in these profiles (Mehta, 1991; Braud et al., 2004). As mentioned earlier, the recirculation region lasts until $x/\theta_i \approx 44$. However, a velocity defect is observed downstream of this location (see profile at $x/\theta_i = 50$ in figure 6.5a). Thus, the wake effect extends beyond the recirculation region, shifting the location where self-similarity is achieved further downstream (Dziomba & Fiedler, 1985; Braud et al., 2004). The mean velocity profiles downstream of $x/\theta_i = 130$ (see figures 6.5b to d) collapse on to the *tanh* profile, suggesting the absence of the splitter-plate wake


Figure 6.6.: Profiles of fluctuating quantities at different streamwise locations non-dimensionalised with $(\Delta U)^2$. (a) Streamwise component of TKE $(R_{uu} = \overline{u'u'})$; (b) Transverse component of TKE $(R_{vv} = \overline{v'v'})$; (c) Reynolds shear stress $(R_{xy} = -\overline{u'v'})$.

effect. Furthermore, as observed in figure 6.4, the mixing layer has a linear growth downstream of this location as well, which suggests that with respect to the mean flow quantities, the current mixing layer achieves self-similarity at $x/\theta_i = 130$.

6.3.3. Fluctuating statistics

However, as observed in previous works (Bell & Mehta, 1990; Mehta, 1991), the turbulent properties become self-similar further downstream. Profiles of the streamwise component ($R_{uu} = u'u'$) of the turbulent kinetic energy (TKE) at selected streamwise locations is shown in figure 6.6a. Just after the roll-up of vortices, the profiles contain three extrema (see $x/\theta_i = 80$), a big one in the high-velocity side, a small one in the low-velocity side and one in the centre close to $y/\delta_0 = 0$. This central extremum is expected to be due to the extended effect of the recirculation region (Laizet et al., 2010). Eventually they evolve into a bell-shaped profile with just one extremum within the mixing layer as observed in previous research works. Profiles of the transverse component of the TKE ($R_{vv} = \overline{v'v'}$, figure 6.6b) have a similar bell-shaped profile in the entire measured domain. Both these parameters show that the fluctuations are contained within the mixing layer. Furthermore, once the wake effect disappears, the maxima in the profiles of both these parameters occur in the high-velocity side. Profiles of the Reynolds shear stress R_{xy} (= $-\overline{u'v'}$) are also presented at different streamwise locations (figure 6.6c). Similar to the observations of Braud et al. (2004) with a thick splitterplate trailing edge, in the region where the wake effect is significant, the profiles of the Reynolds stress contain two extrema. The high-velocity side has a positive extremum while the low-velocity side contains a comparatively smaller negative extremum as seen in the profile at $x/\theta_i = 50$ (more pronounced further upstream). The negative extremum decreases along the stream and eventually disappears completely indicating the decrease in the influence of the splitter-plate wake. Once the wake effect is overcome by the mixing layer, the profiles of the Reynolds stress have a classical bell-shape with just one extremum within the mixing layer as seen in previous works (Oster & Wygnanski, 1982; Bell & Mehta, 1990; Huang & Ho, 1990; Braud et al., 2004).

However, the Reynolds stress profiles (and also the profiles of the R_{uu} and R_{vv}) do not collapse onto each other in the measured domain. Following these observations, it is estimated that the current unforced mixing layer does not attain self-similarity with respect to fluctuating parameters within the measured domain. This is expectable as previous research efforts such as Bradshaw (1966) and Huang & Ho (1990) have shown that in mixing layers generated by two un-tripped boundary layers (as is the case in the present study), the fluctuating quantities attain self-similarity beyond $x/\theta_i \approx 1000$, which is beyond the currently measured spatial domain.

6.3.4. Identification of instabilities through LST and power spectra

It is well known that the mixing layer flow field is unstable to the perturbations arising from the inviscid Kelvin-Helmholtz (KH) instability. This is the primary instability mechanism that controls the dynamics of large-scale vortical flow structures in this flow configuration (see section 2.3.1.1). Sub-harmonic frequencies of the most-amplified or fundamental KH instability, forming immediately downstream of the splitter plate, act as catalyst to the pairing process and hence are observed to control the growth of the mixing layer (Ho & Huang, 1982; Mansour et al., 1988; Husain & Hussain, 1995). Thus, knowing the frequency (or wavelength) of these two-dimensional waves forming initially are absolutely necessary. In order to determine the fundamental instability, linear stability theory (LST) is applied on the velocity profiles acquired through PIV.

As observed in figure 6.5a, a significant velocity defect is seen in the u_N -velocity profiles initially, due to the splitter-plate wake effect. By comparing the experimentally determined vortex passage frequency initially (which is a result of the fundamental KH instability) with that computed through LST based on *tanh* profile, Ho & Huang (1982) showed that the presence of the wake does not change the frequency of the fundamental instability. However, it might affect the amplification rate of the low-frequency instabilities (as shown by Miksad, 1972). Hence, linear stability analysis is carried out on the *tanh* fit to this velocity profile (red line in figure 6.5a). The spatial growth rate (or wave number) is then obtained by solving the Orr-Sommerfeld equation with the Chebyshev collocation method as described in section 3.3.1 (and 3.3.1.3).

The non-dimensional spatial amplification rates $(-\alpha_i \cdot \delta_\omega(x))$ of KH instability waves with different frequencies computed using LST is shown in figure 6.7a. Here, the vertical access denotes the frequency of the KH instability represented in non-dimensional form as the Strouhal number $St_{\theta_i} = f\theta_i/\overline{U}$, where θ_i and \overline{U} are the momentum thickness of the high-velocity boundary layer and the average velocity of the two streams respectively. These parameters are chosen to compute the Strouhal number following the works of Monkewitz & Huerre (1982) and Ho &



Figure 6.7.: (a) Non-dimensional amplification rate $(-\alpha_i \cdot \delta_\omega(x))$ of primary instabilities in the unforced mixing layer at different frequencies computed using LST. The black curve represents the most amplified instability mode at each streamwise station. (b) Amplification N-factors computed from $-\alpha_i$ according to equation 3.6. (c) Zoomed-in view of non-dimensional amplification rate $(-\alpha_i \cdot \delta_\omega(x))$ from the dashedgreen rectangle in (a).

Huang (1982).

The frequency of the most unstable fluctuations is seen to increase until $x/\theta_i \approx 14$ and then starts to decrease, as demonstrated by the black curve in figure 6.7a (and c) which represents the most amplified instability mode at each streamwise station. Until $x/\theta_i \approx 40$, fluctuations greater than $St_{\theta_i} = 0.044$ are most amplified. However, due to the occurrence of the splitterplate wake, the fluctuations in this region are discounted. Beyond this streamwise station, frequency of the most unstable fluctuations reduces along the streamwise direction as expected. This feature becomes clearer in the amplification N-factors of the different instability modes computed by equation 3.6 presented in figure 6.7b. This shows that lower frequency fluctuations are more amplified further downstream. However, it is important to note that once the



Figure 6.8.: Non-dimensional, normalised power spectra $(\Phi_N^{v'} \cdot \Delta f/\overline{U}^2 \text{ (dB)}, \Delta St_{\theta_i} = 3.67 \cdot 10^{-4}, \Delta f = 5 \text{ Hz})$ of the v'-velocity component sampled along the $y_{0.5}$ -trace in figure 6.3.

fundamental instability becomes neutrally stable, sub-harmonic instability growth is influenced by their non-linear interaction with the fundamental (Monkewitz, 1988). This renders the current LST computation as inaccurate. Thus, only trends can be observed and no comments can be made regarding the computed values.

The vorticity thickness of this mixing layer along the streamwise direction presented in figure 6.4 shows a sudden increase between $x/\theta_i = 40$ and until 75 which signifies the roll-up of spanwise vortices due to the disturbance waves generated by the fundamental KH instability. In this streamwise region, LST predicts the most unstable modes to be between $St_{\theta_i} = 0.034$ and 0.013 as can be seen in the zoomed-in view of the amplification rate in figure 6.7c. Thus, one among these modes must be the fundamental KH instability of the current mixing layer configuration.

Another method to determine the fundamental KH instability is to measure the passage frequency of the spanwise vortices initially, which will correspond to the frequency of the fundamental instability (Ho & Huang, 1982). To this effect, time-series of the fluctuations of the transversal velocity component (v') sampled along the $y_{0.5}$ line in the unforced mixing layer (dashed (black) line in figure 6.3) are analysed in the Fourier domain. The power spectra are computed using Welch (1967) method as described in section 3.3.2. To this effect, the acquired time-series was broken down into overlapping segments, each containing 1900 snapshots and overlapping one another by 50%. This resulted in a final frequency resolution of $\Delta St_{\theta_i} = 3.67 \cdot 10^{-4}$ ($\Delta f = 5$ Hz). The spectra are normalized by equation 3.11 and non-dimensionalised with \overline{U}^2 . The computed power spectra is shown in figure 6.8. Here, a discontinuity is observed in the power spectra between $x/\theta_i \approx 500$ and 550 which is an artefact of the stitching of the near and far-wake PIV FOVs (similar to figure 6.5c). A similar discontinuity is observed at higher frequencies around $x/\theta_i = 800$ which is expected to be due to the stitching of the frames of the two cameras when capturing the far-wake field, done on *LaVision Davis* suite.

Fluctuations with highest energy correspond to $St_{\theta_i} = 0.022$ (frequency band of 300 ± 10 Hz). These are observed to be strongest close to the splitter plate and reduces downstream, becoming comparatively low for $x/\theta_i > 350$. Thus, the fundamental KH instability in the current flow configuration has a Strouhal number of $St_0 = 0.022$ ($f_0 = 300$ Hz). The corresponding wavelength of the fundamental instability wave is λ_0/θ_i (= $\overline{U}/f_0\theta_i = 1/St_0$) = 45.5 ($\lambda_0 = 25$ mm). This instability is in the band of frequencies that the LST

computation predicted could be the fundamental instability. This is a demonstration of the feedback mechanism. As highlighted by Ho & Huang (1982), stability computations provides a band of possible frequencies that can be the fundamental KH instability. But the feedback from vortex pairing events occurring downstream fine-tunes the development of the mixing layer into a specific frequency.

The Strouhal number computed with the local integral momentum thickness at $x/\theta_i = 91$ (solid blue marker in figure 6.4) is $f_0\theta(x)/\overline{U} = 0.03$ which corroborates well with past works involving linear stability theory (see Ho & Huerre, 1984). The choice of this location is motivated by the suggestions of Ho & Huerre (1984), who recommend that comparison between measured and computed values of Strouhal number must be made one instability wavelength downstream of the splitter-plate trailing edge and before the occurrence of any non-linear growth and interactions of sub-harmonic instabilities. Given the manifestation of the recirculation region up to $x/\theta_i \approx 44$ in the current set-up, this comparison is carried out approximately one instability wavelength after the disappearance of this recirculation region. Finally, following the work of Michalke (1965a), if the free-stream velocity of the high-velocity stream (U_2) is used as the characteristic velocity to compute the Strouhal number of the fundamental KH instability, it is found to be ($f_0\theta_i/U_2 =$) 0.017 which corroborates well with their computation.

Fluctuations at the harmonic of the fundamental instability band $(St_{\theta_i} = 2 \cdot St_0 = 0.044)$ are also observed. However they contain lesser energy and occur in a smaller region along the streamwise direction. Furthermore, these instabilities are not unstable as indicated by the LST computations in figure 6.7a. Fluctuations are also observed around $St_{\theta_i} = 0.011$ (f = 150 Hz)which is in the frequency band about the first sub-harmonic of the fundamental KH instability $(St_0/2)$. The onset of these fluctuations are downstream compared to the fundamental instability, as expected. Additionally they become strongest at $x/\theta_i \approx 350$ which is approximately the location where the fluctuations of the fundamental instability are observed to become very weak. Fluctuations at $St_{\theta_i} = 0.016$ (f = 220 Hz) are seen to co-exist with the fundamental instability, although with less energy. However, the energy of these fluctuations maximise further downstream at $x/\theta_i \approx 200$.

6.3.5. Spatial organisation of fluctuations

Spectral POD (SPOD) is now employed to investigate the spatial organisation of the fluctuations corresponding to the different KH instabilities within the unforced mixing layer. The formulation of SPOD was introduced in section 3.3.3.

6.3.5.1. Parameters for SPOD computation

For all the test cases reported in this thesis, including the unforced mixing layer currently being discussed, SPOD was performed on 19,000 snapshots captured through time-resolved PIV. A block size of $N_{fr} = 1900$ snapshots with 50% overlap was used, leading to 20 (number of blocks) realizations of the plane mixing layer. Similar to the power spectra reported in the previous section, this resulted in a frequency resolution of $\Delta S t_{\theta_i} = 3.67 \cdot 10^{-4}$ ($\Delta f = 5$ Hz) and at each frequency there were $N_k = 20$ SPOD modes. Hereafter, the notation $\Psi(St_{\theta_i}, k)$ and $\Lambda(St_{\theta_i}, k)$ will be used to refer to the eigenfunction and eigenvalue respectively, of each mode (k) at a particular non-dimensional frequency (St_{θ_i}) .



Figure 6.9.: SPOD decomposition of the unforced mixing layer. (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) Convergence of SPOD modes.

The integrated TKE ($\Lambda(St_{\theta_i}, k)$) of the first two modes at each frequency, normalised with the total TKE of the flow field ($\sum \Lambda(St_{\theta_i}, k)$) is presented in figure 6.9a. Also shown in the figure is the normalised TKE of all modes at a particular frequency ($\sum_k \Lambda(St_{\theta_i}, k)$, black \Box). The convergence of POD is the ratio between the cumulative energy contained in a certain number of modes and the total energy of the flow (Delville, 1994). For SPOD, the convergence of the modes is computed according to equation 6.2

$$c_k = \frac{\sum_{i=1}^k \left[\int \Lambda(St_{\theta_i}, i) \cdot df \right]}{\sum_k \sum_f \Lambda(St_{\theta_i}, k)}$$
(6.2)

The energy contained in the whole field in a particular mode is found by integrating over all the frequencies (integral in equation 6.2). This criterion for the SPOD computation over the unforced mixing layer is shown in figure 6.9b. The first mode contains about 35% of the total turbulent kinetic energy of the flow field. While 99% of the energy was captured by about 15 modes.

6.3.5.2. Analysis of SPOD eigenfunctions

From figure 6.9a, it is observed that the fluctuations at $St_{\theta_i} = 0.022$ contains the highest energy. This agrees with the discussion in section 6.3.4 where the frequency of the fundamental instability was found to be $St_0 = 0.022$. The first eigenmode at this frequency contains 89% of the TKE of all modes at this frequency. The vertical component of the corresponding eigenfunction $(\Psi_v(0.022, 1))$ is presented in figure 6.10a. The colour scale in the figure shows the distribution



Figure 6.10.: Vertical component of the eigenfunction of the first mode $(\Psi_v(St_{\theta_i}, 1))$ at different frequencies (grey - negative; black - positive). The colour-scale in (a) represents the normalised Reynolds stress component of the fluctuations corresponding to this mode $(R_{xy}(St_{\theta_i}, 1)/(\Delta U)^2)$.

of the Reynolds shear stress of the corresponding fluctuating field. For this purpose, the flow field was reconstructed with just this particular SPOD mode using equation 3.15 and then the Reynolds stress was computed.

Coherent structures are clearly visible until $x/\theta_i \approx 250$ in this eigenfunction. It is important to note that these structures only portray the spatial organisation of fluctuations in the flow field and do not represent any physical "vortical" flow structures. Initially, these coherent structures attain an obtuse angle with the streamwise direction. This corresponds to the spatial domain where the Reynolds stress is positive. As discussed by Rajaee et al. (1994) who performed 'snapshot' POD (Sirovich, 1987) on their hot-wire measurements of the plane mixing layer, this positive Reynolds stress working against the positive shear strain rate of the mean flow field represents transfer of energy from the mean flow to the fluctuating field. The structures eventually attain a vertical alignment at $x/\theta_i \approx 80$ which represents the mode attaining its



Figure 6.11.: Non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of the SPOD eigenfunction (vertical component) presented in figure 6.10 along the $y_{0.5}(x)$ trace.

maximum amplitude (negligible Reynolds stress). Beyond this point, the structures have an acute angle with the streamwise direction and the Reynolds stress in this region is observed to be negative. The action of this negative Reynolds stress against the positive rate of strain of the mean flow results in the transfer of energy from the corresponding fluctuating field to the mean flow. This criterion of the orientation of the structures in the SPOD eigenfunctions will be utilized hereafter to analyse the different instability waves in the mixing layer.

An estimate of the amplitude of different SPOD modes along the streamwise direction is necessitated to better analyse the amplification (and/or decay) of fluctuations corresponding to a particular SPOD mode and also compare different modes. To this effect, the parameter A is introduced here. This estimate is obtained using only the transversal component of the SPOD eigenfunctions ($\Psi_v(St_{\theta_i}, k)$), as shown in the system of equations 6.3.

$$\widetilde{Q}_{R,v}(St_{\theta_i},k) = \Theta(St_{\theta_i},k)\Psi_v(St_{\theta_i},k)\sqrt{\Lambda(St_{\theta_i},k)}$$
(6.3a)

$$\Phi_{R,v}(St_{\theta_i}, k) = \frac{\sum_{1}^{N_k} \left[|Q_{R,v}(St_{\theta_i}, k)|^2 \right]}{N_k}$$
(6.3b)

$$A(St_{\theta_i}, k) = \sqrt{\Phi_{R,v}(St_{\theta_i}, k)}$$
(6.3c)

Firstly, the contribution of the desired SPOD mode to the component of DFT at that frequency $(\tilde{Q}_{R,v}(St_{\theta_i},k))$ is computed by equation 6.3a. This operation is also used to reconstruct the time-series from the desired SPOD modes as shown in equation 3.15. Given the application of Welch (1967) method here, this operation produces N_k (= 20) reconstructions of the DFT, corresponding to the number of blocks used to perform SPOD. In order to estimate the amplitude of a mode for the entire time-series, a periodogram ($\Phi_{R,v}(St_{\theta_i},k)$) is computed using these DFT estimates by equation 6.3b. The parameter $A(St_{\theta_i},k)$ is then estimated by finding the square-root of this periodogram as shown in equation 6.3c. A validation of estimating the amplitude of different fluctuations in this manner is presented in appendix A.

This computation pertaining to all the SPOD modes being considered here, along the $y_{0.5}$ trace is shown in figure 6.11. It can be seen that the fluctuations with the highest amplitude correspond to $St_{\theta_i} = St_0 = 0.022$. It is clear that these fluctuations amplify initially, when the structures in the eigenfunction (figure 6.10a) are oriented with an obtuse angle with respect

to the streamwise direction. The maximum amplitude is attained at $x/\theta_i \approx 75$, which is close to the streamwise station where the structures attained a vertical alignment. Beyond this, the amplitude diminishes and reaches a low asymptotic level at about $x/\theta_i \approx 350$. This once again agrees with the power spectra in figure 6.8 where fluctuations at $St_{\theta_i} = 0.022$ were observed to be very weak downstream of this streamwise station.

The eigenfunction $\Psi_v(0.011, 1)$ containing fluctuations at the first sub-harmonic of the fundamental instability is shown in figure 6.10b. From the structures here and the amplitude A(0.011, 1) in figure 6.11, these fluctuations are observed to grow until $x/\theta_i \approx 340$ where they attain maximum amplitude and decay further downstream. Structures pertaining to fluctuations at the second sub-harmonic of the fundamental instability are represented by the eigenfunction $\Psi_v(0.0055, 1)$ in figure 6.10c. Compared to previous fluctuations, initial structures are only observed here further downstream at $x/\theta_i \approx 210$ and thus, these fluctuations amplify only beyond this streamwise station as seen in figure 6.11. These fluctuations amplify until $x/\theta_i \approx 700$ beyond which they start to decay. Finally, the eigenfunction of fluctuations at $St_{\theta_i} = 0.0117$ which is close to the first sub-harmonic instability is presented in figure 6.10d. The organisation of these fluctuations is shown here as the instability corresponding to these fluctuations will be studied upon later on, to investigate the effect of detuning the first sub-harmonic on the mixing layer. The structures in this eigenfunction resemble those at the first sub-harmonic but these fluctuations reach their maximum amplitude marginally upstream at $x/\theta_i \approx 320$, beyond which they start to decay.

6.3.6. Brief summary

The results pertaining to the characterisation of the baseline or unforced mixing layer is reported in this section. The analysis of the acquired results indicate that the boundary layers on either sides of the splitter plate is laminar at the trailing edge. Owing to the thickness of this trailing edge, a wake effect is observed until $x/\theta_i \approx 40$, at which station, the onset of the mixing layer flow configuration is observed. Although, investigating the velocity profiles revealed the existence of a velocity defect due to this wake further downstream, until $x/\theta_i \approx 130$. Beyond this streamwise station, the profiles are well approximated by a hyperbolic tangent profile. Through this observation, combined with the linear growth pattern, it can be concluded that the current mixing layer is self-similar after $x/\theta_i = 130$ based on mean-flow quantities. However, with respect to the the fluctuating quantities, the current layer does not attain self-similarity in the measured spatial domain. This is expected as previous research efforts have shown (Browand & Latigo, 1979; Huang & Ho, 1990) that mixing layers originating from laminar initial conditions attain self-similarity beyond $x/\theta_i = 1000$, which is the approximate extent of the current measurement domain.

In order to investigate the influence or effectiveness of instability control, which forms the rest of this chapter, it is imperative to determine the frequency of the fundamental KH instability in this mixing layer. To this effect, two analytical tools were deployed on the acquired velocity time-series. Namely, linear stability theory and power spectra. Through these computations, it was ascertained that the fundamental instability in the current mixing layer configuration has a frequency of $f_0 = 300$ Hz. This corroborates well with those reported in previous research efforts (Michalke, 1965a; Ho & Huerre, 1984), thus indicating the realisation of a canonical mixing layer. The corresponding Strouhal number, computed based on average velocity of two streams \overline{U} and high-velocity boundary layer momentum thickness at the splitter-plate trailing edge θ_i , is $St_0 = 0.022$. Finally, the organisation of fluctuations corresponding to select primary



Figure 6.12.: Sample input AC waveform (blue) and square-wave burst modulation (red). $V_{ac} = 14$ kV, $f_{ac} = 2$ kHz, $f_b = 150$ Hz (50% duty cycle). Thick green line indicates trigger to PIV acquisition system.

instability modes that will be investigated in the rest of this chapter was desired. This was obtained by applying SPOD on the acquired time-series, the results of which was also reported. These computations once again confirmed that the fundamental KH instability corresponds to a Strouhal number of $St_0 = 0.022$.

6.4. MANIPULATION OF PRIMARY KH INSTABILITIES THROUGH SPANWISE-UNIFORM EFD FORCING

The effect of forcing the primary *Kelvin-Helmholtz* (KH) instability in the mixing layer through spanwise-uniform EFD forcing applied by the installed DBD plasma actuator is discussed here. The PMMA attachment that houses a DBD plasma actuator which exerts a spanwise-uniform EFD body force over the flow field is fixed to the aluminium (splitter) plate for this purpose.

6.4.1. DBD actuator operating conditions, forced instabilities and momentum coefficient

Design of the employed PMMA attachment that houses the spanwise-uniform DBD plasma actuator is as reported in section 4.5. The construction is such that, the actuator exerts a spanwise-uniform EFD body force into the splitter-plate boundary layer on high-velocity side, at $x/\theta_i = -22$ (x = -12 mm). The actuator was operated with a burst-modulated sinusoidal waveform. As reported previously, the fundamental instability of the current mixing layer was determined to be $f_0 = 300$ Hz in dimensional terms. Thus, a frequency of $f_{ac} = 2$ kHz, which is more than six times that of the fundamental KH instability, was chosen for the sinusoidal waveform. This AC frequency is responsible for the ionisation process and hence the production of the EFD body force. A low-frequency (f_b) square-wave burst modulation was applied on the input sinusoidal waveform which resulted in a final input signal as shown in figure 6.12. The operating parameter that is varied to investigate the effect of exciting different KH instability modes in the mixing layer is the frequency of the square-wave burst modulation f_b . This will be represented with the corresponding Strouhal number computed as $St_f = f_b \theta_i/\overline{U}$. The air-exposed electrode of the actuator was supplied with the high-voltage signal while the encapsulated electrode was grounded. Given the operational frequency band necessary for these investigations (the maximum being $f_{ac} = 2$ kHz), the DBD actuator was powered using the *Trek 30/20A* high-voltage amplifier introduced in section 3.2.3.1.1. The input signal to the amplifier was supplied by the *Teledyne LeCroy WaveStation 3082* generator. Furthermore, for these test cases, in order to achieve a PIV acquisition synchronised with the operation of the DBD plasma actuator, the trigger to initiate the acquisition was also provided by the same waveform generator. The trigger to the PIV system was synchronized with the beginning of a burst modulation cycle (green vertical line at t = 0 in figure 6.12). It is worth noting that for all tested cases, the DBD plasma actuator was operational for a few seconds before the start of the PIV acquisition. This was to ensure that the mixing layer has reorganised itself according to the applied EFD forcing before the PIV acquisition.

The LST analysis, power spectra and spectral POD performed on the unforced mixing layer previously (sections 6.3.4 and 6.3.5) showed that the fundamental KH instability of the current mixing layer flow configuration corresponds to fluctuations at Strouhal number $St_0 = 0.022$ $(f_0 = 300 \text{ Hz})$. Thus, with the DBD plasma actuator, perturbations are imparted into the mixing layer here to force the fundamental instability $St_f = St_0 = 0.022$ $(f_b = 300 \text{ Hz})$, its first subharmonic $St_f = St_0/2 = 0.011$ $(f_b = 150 \text{ Hz})$ and its second sub-harmonic $St_f = St_0/4 = 0.0055$ $(f_b = 75 \text{ Hz})$. This is to investigate the effect of sub-harmonic instability forcing on the dynamics of large-scale vortices within the mixing layer.

Finally, pairing in the unforced mixing layer does not occur periodically. Due to the inability of the sub-harmonic to phase-lock with the fundamental instability, a spatial and temporal jitter is common. This manifests as sideband frequencies and amplitude modulation of the sub-harmonic instability wave. These sidebands are known to occur in the frequency band $St_0/2 \pm 0.08 \cdot St_0$ (Monkewitz, 1988; Thomas & Chu, 1989). Excitation of a detuned sub-harmonic instability will result in the amplification of sideband frequencies and amplitude-modulation of the sub-harmonic as demonstrated by Husain & Hussain (1995). In fact, such a forcing condition is more "realistic" as sub-harmonic instability forcing that is phase-locked with the fundamental instability cannot always be guaranteed in practical applications. Thus, the impact of imparting fluctuations marginally detuned from the first sub-harmonic instability is also investigated here. To this effect, forcing the instability at $St_f = 0.0117$ ($f_b = f_b/2 + 10$ Hz = 160 Hz), which corresponds to detuning the first sub-harmonic by about 3.3% of the fundamental instability is investigated ($St_0/2 + 0.033 \cdot St_0 = 0.0117$).

In order to represent the applied force in non-dimensional terms and to give an estimate of its magnitude relative to the experimental flow field, the momentum coefficient (c_{μ}) is utilized here as well. This parameter is computed here in a similar manner as done in section 5.4.2 for the swept-wing boundary layer control investigation.

$$c_{\mu} = \frac{T_x}{\frac{1}{2}\rho U_2^2 \cdot \theta_i} \cdot \frac{w_S}{\lambda_z} \tag{6.4}$$

Here, ρ (= 1.225 kg m⁻³) is the density of air at ambient conditions, θ_i (= 0.55 mm) is the momentum thickness of the high-velocity side boundary layer close to the splitter-plate trailing edge and U_2 is the high-velocity free-stream. The ratio w_S/λ_z is to account for the spanwise duty when three-dimensional spanwise-modulated forcing is applied on the flow field (see section 6.5). In the case of spanwise-uniform EFD forcing, this ratio is equal to 1. The variation of



Figure 6.13.: Variation of momentum coefficient c_{μ} with average input power P_{avg} . Solid red marker is an estimate of c_{μ} when the spanwise-uniform DBD plasma actuator is operated at $V_{ac} = 14$ kV, $f_{ac} = 2$ kHz with a burst-modulation duty of 50%.

the momentum coefficient computed this way with the average power consumed is presented in figure 6.13. Corresponding to the generated wall-parallel thrust shown in figure 4.22c, the momentum coefficient too increases linearly with the power consumed by the actuator.

For most cases presented in this section, the DBD plasma actuator is operated at voltage amplitude $V_{ac} = 14$ kV and frequency $f_{ac} = 2$ kHz. The duty-cycle of burst-modulation cycle is also held constant at 50%. A brief discussion on the effect of two operating parameters on the mixing layer, i.e. burst-modulation duty cycle and AC voltage amplitude, is reported at the end in section 6.4.7. From figure 4.22a, the average power consumed by the DBD plasma actuator when operated at $V_{ac} = 14$ kV, $f_{ac} = 2$ kHz is found to be $P_{avg} = 52.2$ W m⁻¹ (solid black marker in figure 4.22a). Considering the application of a burst modulation with 50% duty cycle on the input AC signal to manipulate the primary instabilities in the mixing layer, this value of power can be halved (see figure 2.5a, reproduced from Benard & Moreau, 2009). Thus, the average power consumed by the actuator while applying the necessary EFD force on the mixing layer is $P_{avg,f} = 26.1$ W m⁻¹.

The same cannot be done to estimate the thrust generated, due to inertial and viscous effects of the flow. However, given the proportional variation of the generated thrust with the consumed power as seen in figure 4.22c, the thrust value can be estimated from this plot. For the current operating conditions (i.e. $P_{avg,f} = 26.1 \text{ W m}^{-1}$), the wall-parallel component of thrust is found to be $T_{x,f} = 7.26 \text{ mN m}^{-1}$. The momentum coefficient has to be approximated in a similar manner from figure 6.13. At the present flow conditions, the corresponding momentum coefficient is $c_{\mu} = 0.22$ when the average input power to apply the necessary forcing is $P_{avg,f} = 26.1 \text{ W}$ m⁻¹. This is represented by the solid red marker in figure 6.13. The momentum coefficients at different operating conditions being investigated in this section is tabulated in 6.1.

Finally, Jolibois & Moreau (2009) reported the maximum velocity induced by a DBD actuator with a 2 mm thick PMMA as the dielectric substrate, at different input electric power. Using the measured power consumption of the DBD actuator at the same operating conditions ($P_{avg,f} =$ 26.1 W m⁻¹), the maximum time-averaged velocity in the induced jet is estimated to be $U_f \approx 2.5$ m s⁻¹. The velocity ratio between this velocity of the induced jet and the high-velocity freestream is then $U_f/U_2 \approx 0.25$. The momentum coefficient of $c_{\mu} = 0.22$ for this input power is comparable to the estimated velocity ratio. Since the momentum of the free-stream is not changing, the computed momentum coefficient is representative of this velocity ratio. The ratio

		<i>,</i>	1	v
	V_{ac} (kV)	f_{ac} (kHz)	f_b duty (%)	c_{μ}
-	14	2	50	0.22
	14	2	20	0.06
	14	2	80	0.37
	12	2	50	0.11
	16	2	50	0.35

Table 6.1.: Momentum coefficient at different operating condition. The first row of values (bold) represents the operating parameters used predominantly in this section. The bold values in other rows represent the operating parameter being varied for the parametric study in section 6.4.7.

of the induced velocity with average velocity of the two streams (convective speed of large-scale vortices Oster, Wygnanski, & Fiedler, 1977) is $U_f/\overline{U} \approx 0.33$.

As reported in table 6.1, the different test cases investigated in this here correspond to a range of c_{μ} between 0.06 and 0.37. This range partially coincides with the range reported by Parezanovic et al. (2014), who tested input amplitudes corresponding to momentum coefficients $0.25 < c_{\mu} < 0.5$. Apart from two of the current test cases, the rest of them (including the primary test case with $c_{\mu} = 0.22$) correspond to input amplitudes lower than that reported by Parezanovic et al. (2014). When the induced flow velocity is considered, the micro-jets used to control the mixing layer in the experiments of Parezanovic et al. (2014) seems to be twice the convection speed of their mixing layer. While in the experiments of Parezanović et al. (2016), the employed micro-jet velocity of 3 m s⁻¹ is comparable to the convection speeds in the experiment. Mathis et al. (2007) report the velocity ratio between the employed steady-blowing jets and the high-velocity free-stream to be 0.8. Comparing these to the primary forcing condition tested here at $c_{\mu} = 0.22$ ($U_f/U_2 \approx 0.25$, $U_f/\overline{U} \approx 0.33$), the current actuation is clearly significantly weaker.

6.4.2. Mean streamwise-velocity and Reynolds shear stress profiles

The \overline{u}_N -velocity profiles of the forced mixing layers at different streamwise stations are presented in figure 6.14. As in figure 6.5, here too, the profiles are shifted with the $y_{0.5}$ -location but nondimensionalised with the momentum thickness of the high-velocity boundary layer close to the splitter-plate trailing edge θ_i . As can be seen from the velocity profiles at $x/\theta_i = 50$ (figure 6.14a), the velocity defect in the forced mixing layers is lower compared to the unforced flow. This indicates a reduction of the splitter-plate wake effect due to the applied forcing.

In the work of Mehta (1991), it is shown that the effect of the splitter-plate wake strength reduces with the velocity ratio r of the mixing layer. In the current scenario, the applied EFD forcing imparts momentum into the high-velocity boundary layer and accelerates the flow (also observed by Singh & Little, 2020). This will result in a local reduction of the velocity ratio just downstream of the splitter plate, which ultimately leads to the reduction of the strength of the splitter-plate wake. An estimate of the local change of velocity ratio is made here. As reported in section 6.4.1, the maximum time-averaged velocity in the induced jet at the current operating conditions ($P_{avg,f} = 26.1 \text{ W m}^{-1}$) is estimated, through the experimental work of Jolibois & Moreau (2009), to be $U_f \approx 2.5 \text{ m s}^{-1}$. This is comparable to the estimate of the induced velocity due to momentum input by the DBD actuator obtained from the PIV measurement. The difference of the \overline{u}_N velocity profiles of the forced and unforced mixing layers ($\overline{u}_{N,St_f} - \overline{u}_{N,UML}$) at $x/\theta_i = 2.5$ is shown in figure 6.15. The maximum difference in velocity



Figure 6.14.: Profiles of \overline{u}_N velocity of the forced mixing layers ($c_\mu = 0.22$) at different streamwise stations.

due to the applied forcing is about 0.4, which in dimensional form is $0.4 \cdot \Delta U = 2 \text{ m s}^{-1}$. This estimate being lesser than that from the consumed power can be attributed to frictional losses incurred due to the presence of the splitter-plate surface downstream of the DBD actuator's plasma-generating edge. With this estimate from PIV (2 m s⁻¹), the velocity of the top stream close to the splitter-plate trailing edge can be approximated to be $\approx 12 \text{ m s}^{-1}$. Thus, just downstream of the trailing edge, the velocity ratio will be $r \approx 0.42$, instead of r = 0.5 in the unforced mixing layer, hence resulting in the weakening of the splitter-plate wake.

The u_N -velocity profiles at $x/\theta_i = 130$ (figure 6.14b) indicate the profiles of the forced mixing layers collapsing on to that of the unforced flow. This is in fact the streamwise station where the unforced mixing layer attained self-similarity with respect to the mean flow quantities as discussed in section 6.3.2. This implies that the applied forcing is not strong enough to modify the base-flow profile initially, but probably alters the spectral content within the flow field. Differences in the profiles do develop further downstream as seen in figures 6.14c and d. These are expected to be a result of forcing different instabilities which will potentially control the dynamics and interactions of vortical flow structures and consequentially, lead to different growth patterns of the forced mixing layers.

Profiles of the Reynolds shear stress $(R_{xy} = -\overline{u'v'})$ of the four forced mixing layer at four different spanwise stations is presented in figure 6.16. At $x/\theta_i = 50$, the maximum Reynolds stress of all forced mixing layers is lower than the unforced flow which is once again an artefact



Figure 6.15.: Difference between the \overline{u}_N -profiles at $x/\theta_i = 2.5$ of the forced ($c_\mu = 0.22$) and unforced mixing layers for three forcing cases.



Figure 6.16.: Profiles of Reynolds shear stress $(R_{xy} = -\overline{u'v'})$, non-dimensionalised with $(\Delta U)^2$ of the forced mixing layers $(c_{\mu} = 0.22)$ at different streamwise stations.

of the reduced strength of the splitter-plate wake. While the \overline{u}_N -velocity profiles at the same streamwise stations as here (figure 6.14) were only marginally different for the four forced mixing layers and also compared to the unforced flow, the Reynolds shear stress profiles indicate significant differences. This is a preliminary indication of the applied forcing manipulating the fluctuating fields while it is not strong enough to modify the mean-flow profile itself. Far down-



Figure 6.17.: Non-dimensional vorticity thickness of all forced mixing layers ($c_{\mu} = 0.22$) close to the splitter-plate trailing edge.

stream at $x/\theta_i = 800$, the differences between the different cases is much lower compared to the other stations. This could be an indication that the four forced modes are not in control of the dynamics of the mixing layer any more.

6.4.3. Forced mixing layer growth

The evolution of the vorticity thickness δ_{ω} (equation 2.6) due to the exerted spanwise-uniform EFD forcing at $c_{\mu} = 0.22$ to excite different KH instabilities within the mixing layer, is discussed here. The weakening of the splitter-plate wake effect and the subsequent reduction in the velocity defect in the forced mixing layers discussed in section 6.4.2, manifests here as well. The EFD forcing results in the initial development of the mixing layer occurring upstream compared to the unforced case as observed in the vorticity thickness evolution close to the splitter-plate trailing edge shown in figure 6.17. Regardless of the forced instability mode, the initial development of the unforced mixing layer is at $x/\theta_i \approx 15$ which is indeed upstream where the initial development of the unforced mixing layer is observed. This is essentially attributed to the action of the time-averaged momentum input by the DBD actuator against the splitter-plate wake as discussed previously in section 6.4.2.

The growth of the different forced mixing layers in the measured domain is shown in figure 6.18. Each forcing condition is presented separately and compared to the unforced mixing layer. Furthermore, linear fits to different segments of the growth curves, that represent the spreading rate of the mixing layer according to the Abramovich (1963)-Sabin (1965) rule (equation 2.7), are also shown in these figures. The initial growth of all the forced mixing layers, indicating the roll-up of two-dimensional spanwise vortices, lasts until $x/\theta_i \approx 60$ as seen also in figure 6.17. Further downstream, forcing the different instabilities results in different growth patterns of the ensuing mixing layers.

Forcing the fundamental instability $(St_f = St_0 = 0.022)$ (figure 6.18a) reduces mixing layer growth significantly ($\approx 0.07 \cdot R$) until $x/\theta_i = 200$ compared to the unforced mixing layer. This is expected to be due to the attenuation of sub-harmonic instabilities (see case mode I of Ho & Huang, 1982) which play a pivotal role in mixing layer growth. The mixing layer then grows rapidly ($\approx 0.7 \cdot R$) until $x/\theta_i = 500$. Further downstream, the growth curve collapses on to that of the unforced mixing layer.



Figure 6.18.: Growth of the mixing layers when different primary KH instabilities are forced using the DBD plasma actuator ($c_{\mu} = 0.22$). Dashed (black) lines represent linear fits to different segments of the growth curves. Black (thick) curve represents the growth of the unforced mixing layer and the dashed (red) line represents the corresponding $0.25 \cdot R$ linear fit (see figure 6.4).

Application of forcing at the first sub-harmonic of the fundamental instability $(St_f = St_0/2 = 0.011)$ leads to a different growth curve (figure 6.18b). After the initial roll-up, the mixing layer starts to grow from $x/\theta_i \approx 100$, eventually attaining a linear growth at a rate of $\approx 0.43 \cdot R$, which is higher than the unforced mixing layer. Post this, the growth halts temporarily from $x/\theta_i \approx 300$ to about $600 \ (\approx 0.089 \cdot R)$. It is re-initiated further downstream and the layer grows linearly thereafter, until the end of the measured spatial domain, at a similar rate ($\approx 0.26 \cdot R$) as the unforced mixing layer. This growth curve bears qualitative similarity to those observed in the experimental efforts of Oster & Wygnanski (1982), Ho & Huang (1982), and Huang & Ho (1990), when sub-harmonic forcing was applied on the mixing layer (see for example figure 2.22, reproduced from Ho & Huang, 1982).

A qualitatively similar growth curve is observed when the second sub-harmonic of the fundamental instability ($St_f = St_0/4 = 0.0055$, figure 6.18c) is forced, although with some differences. After the initial growth, the mixing layer grows linearly further downstream compared to that forced at $St_f = 0.011$, until $x/\theta_i \approx 550$. In this region, the mixing layer spreads at a rate of $0.28 \cdot R$. Beyond this streamwise station, the spreading rate reduces to $0.085 \cdot R$ similar to the previous case. This reduced growth rate last almost until the end of the measured domain.

Forcing the mixing layer at $St_f = 0.0117$ (figure 6.18d) results in a linear growth of the layer from $x/\theta_i = 100$ at a rate of $\approx 0.41 \cdot R$. This linear growth, at a higher rate than the unforced mixing layer, lasts until $x/\theta_i \approx 250$. Beyond this streamwise station and until about $x/\theta_i = 500$ the spreading rate reduces to about $0.13 \cdot R$. The mixing layer then shows rapid growth for a small streamwise distance and collapses on to the growth curve of the unforced mixing layer at $x/\theta_i = 550$. The growth of this mixing layer bears some similarity to that forced at the first subharmonic of the fundamental instability. Indeed, this instability is in the frequency band of the sub-harmonic sidebands $(St_0/2 \pm 0.08 \cdot St_0 \text{ Monkewitz}, 1988; \text{Thomas & Chu}, 1989)$. However, given that the frequency of this instability is marginally higher than the first sub-harmonic, the higher linear growth rate (although comparable in magnitude) ends further upstream in the current case.

When the growth curves of the four forced mixing layers are compared with each other, one can observe that the streamwise location where the applied forcing has a significant influence varies. The effect of the fundamental KH instability forcing on the growth of the mixing layer is detected further upstream compared to when the first sub-harmonic instability is forced which is followed by the second sub-harmonic instability. This particular phenomenon is more evident when comparing the growth curves of the mixing layer forced at the first sub-harmonic instability and at $St_f = 0.0117$ as just discussed. Additionally, when these results are analysed in congruence with the Reynolds stress profiles shown in figure 6.16, the variation of the Reynolds stress correlates with the growth pattern of the forced mixing layer. For example, the maximum Reynolds stress of the mixing layer forced at the fundamental is minimum at $x/\theta_i = 130$, where the growth rate of the layer is marginal and it is significantly higher at $x/\theta_i = 400$, where the also the growth rate of the layer is very high. These results point towards the observed growth patterns being a consequence of the applied forcing manipulating the fluctuating fields. This is expected to be a consequence of the stability characteristics of the flow field, as perturbations pertaining to different instabilities are most amplified, and hence control the dynamics and growth of the layer, at different streamwise locations (Ho & Huang, 1982). Furthermore, in these test cases (especially with sub-harmonic forcing), the applied EFD forcing does increase the growth rate locally but it has failed to realize a permanent increase in the spread/thickness of the mixing layer compared to the unforced case.

6.4.4. Spectral distribution

As done on the unforced mixing layer (figure 6.8), the power spectra of fluctuations of the transverse velocity component (v'), sampled along the $y_{o.5}$ trace of the forced mixing layers is presented in figure 6.19. This power spectra too is normalized according to equation 3.11 and non-dimensionalised with \overline{U}^2 .

The power spectra of the mixing layer forced at the fundamental KH instability (figure 6.19a) shows significant fluctuations at the forced instability $St_{\theta_i} = St_f = 0.022$. These appear just downstream of the splitter-plate trailing edge and span further downstream compared to the unforced mixing layer. Furthermore, energy of fluctuations only at the forced instability is observed to be high here, unlike the unforced case where frequencies in a band around the fundamental instability were amplified. However, energy in the frequency band of the sub-harmonic instabilities are significantly reduced compared to the unforced mixing layer, especially until $x/\theta_i \approx 200$. Strong fluctuations are also observed at the harmonic of the forced fundamental instability close to the splitter-plate trailing edge. These can be attributed to the applied forcing. As demonstrated by Benard & Moreau (2010, discussed in section 2.1.3), fluctuations at the harmonic of the burst frequency are observed in the plasma induced flow. However, these fluctuations will become stable before the fundamental instability (see figure 6.7) and are not expected to affect the dynamics of the mixing layer.

When the mixing layer is forced at $St_f = St_0/2 = 0.011$, fluctuations at the forced subharmonic instability and the fundamental instability are observed in the power spectra (figure



Figure 6.19.: Non-dimensional, normalised power spectra $(\Phi_N^{v'} \cdot \Delta f/\overline{U}^2 \text{ (dB)}), \Delta S t_{\theta_i} = 3.67 \cdot 10^{-4}, \Delta f = 5$ Hz) of the v'-velocity component sampled along the $y_{0.5}$ trace of the mixing layer forced with different primary KH instabilities ($c_{\mu} = 0.22$).

6.19b). Both these fluctuations are seen just downstream of the splitter-plate trailing edge. Strong fluctuations are observed in the frequency bands around $St_{\theta_i} = 0.015$ and $St_{\theta_i} = 0.026$ as well. Spectral distribution in the mixing layer forced with the second sub-harmonic of the fundamental KH instability ($St_f = St_0 = 0.0055$) is presented in figure 6.19d. Similar to the previous cases, fluctuations corresponding to the forced instability is observed just downstream of the splitter-plate trailing edge. While fluctuations in the frequency band around $St_{\theta_i} = 0.015$ are not strong here, those around $St_{\theta_i} = 0.026$ are amplified just downstream of the splitter-plate trailing edge as seen in the previous case.

Exciting fluctuations at $St_f = 0.0117$ which is an instability marginally detuned from the first sub-harmonic (by 3.3% of the St_0) results in the spectral distribution presented in figure 6.19c. As in the previous cases, fluctuations corresponding to the forced mode is observed just downstream of the splitter-plate trailing edge here. Additionally, fluctuations in the frequency bands around $St_{\theta_i} = 0.0154$ and $St_{\theta_i} = 0.0271$ are strong as well. This is similar to the case with first sub-harmonic forcing, although these frequency bands here are centred around

marginally higher frequencies. However, unlike forcing the first sub-harmonic, no fluctuations at the fundamental instability are observed. But considerable energy is seen at $St_{\theta_i} = 0.0235$ which is the harmonic of the forced instability that is in the frequency band of the fundamental instability (see figure 6.8). This does not however, corroborate well with the results of Husain & Hussain (1995).

In their experiments, Husain & Hussain (1995) applied a detuning of 2.5% of their fundamental instability to the first sub-harmonic and excited the shear layer with this signal. They observed the amplification of fluctuations at these sub-harmonic sidebands, but also the sub-harmonic itself due to the feedback from vortex pairing. It is worth noting however, that they also excite the fundamental instability simultaneously. They go so far to suggest that the same might occur with single-frequency excitation as well. In the current mixing layer, this does not seem to be the case. As seen in figure 6.19c, fluctuations at the first sub-harmonic instability are not amplified at all initially. Thus, the imparted fluctuations do not behave like a sub-harmonic sideband. Rather it acts as a sub-harmonic itself, with the new fundamental instability being at $St_{\theta_i} = 0.0235$ (its harmonic).

This behaviour can be attributed to the location and manner in which the fluctuations are being imparted in the current study. In the experiments of Husain & Husain (1995), the fundamental and detuned sub-harmonic shear layer instabilities were excited simultaneously by imparting acoustic fluctuations at the lip of a jet (which corroborates to the splitter-plate trailing edge here). On the other hand, in the current experimental investigation, the fluctuations are being imparted into the high-velocity side boundary layer over the splitter-plate at $x/\theta_i = -22$, rather than directly into the mixing layer flow configuration. As demonstrated by Ho & Huang (1982, discussed previously in section 4.5), higher shear contained within this boundary layerhas a greater influence on the turbulent energy production and thus, instability growth. When forcing is applied at $St_f = 0.0117$, its harmonic also is expected to be energised (Benard & Moreau, 2010). This harmonic being in the frequency band of the fundamental instability seems to project itself as the fundamental, making fluctuations imparted at $St_f = 0.0117$ the first sub-harmonic instability. Although, the action of the pairing feedback mechanism cannot be discounted either. In any case, this scenario (of forcing modifying the fundamental instability) is expected to occur only when the harmonic of the forced fluctuations is within the frequency band of the unforced flow's fundamental instability.

6.4.5. Spatial organisation of fluctuations through SPOD

The results reported until now have demonstrated that the EFD forcing exerted by the DBD plasma actuator, specifically the burst-modulation, influences the growth of the mixing layer and its spectral characteristics. In order to gain more insight into the effect of this unsteady low-frequency forcing, the spatial organisation of the fluctuating fields is sought. For this purpose, SPOD was applied on the PIV time-series. The parameters chosen for these computations are as introduced in section 6.3.5.1. These results are presented and discussed below for the different forced mixing layers being studied here.



Figure 6.20.: SPOD decomposition of the mixing layer forced at $St_f = St_0 = 0.022$ ($c_{\mu} = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of selected frequencies along the $y_{0.5}(x)$ trace. (c), (d) and (e) show the vertical component of the eigenfunction of the first mode $(\Psi_v(St_{\theta_i}, 1))$ of these different frequencies (grey - negative; black - positive).

6.4.5.1. Fundamental instability forcing

The results of SPOD applied on the mixing layer forced at the fundamental instability is presented in figure 6.20. The component of TKE at different frequencies (black \Box in figure 6.20a) shows that the modes with the highest energy pertain to the forced fundamental instability $St_f = St_0 = 0.022$. Concentration of TKE is also observed at the harmonics of the forced mode as expected, however, sub-harmonic instabilities have been attenuated. These are also evident in the amplitude of the corresponding fluctuations $(A_v(St_{\theta_i}, 1))$, computed according to the system of equations 6.3, shown in figure 6.20b. These observations agree well with the power spectra of this case shown in figure 6.19a.

The vertical component of the eigenfunction of the first SPOD mode containing fluctuations of the forced instability ($\Psi_v(0.022, 1)$) is shown in figure 6.20c. When compared to the same mode in the unforced mixing layer (figure 6.10a), clear structures are observed further upstream. This is a direct consequence of the upstream shift of the mixing layer development due to the applied EFD forcing as discussed in section 6.4.3. The fluctuations pertaining to this mode amplifies at a rate of $-\alpha_i \cdot \theta_i \approx 0.0071$ and reaches its highest amplitude at about $x/\theta_i \approx 80$ (similar to the unforced flow case), as observed in figure 6.20b and the vertical alignment of the structures. However, the maximum amplitude of these fluctuations is significantly higher ($A_{v,max}(0.022, 1)/\overline{U} = 0.326$) when compared to the unforced mixing layer (see figure 6.11). Additionally, the structures eventually had an acute angle orientation with the streamwise direction in the unforced mixing layer, the structures remain vertically aligned and are present until about $x/\theta_i \approx 320$. Correspondingly, the amplitude of these fluctuations too attain a low level at $x/\theta_i \approx 550$, which is further downstream compared to the unforced case.

The vertical component of the eigenfunction $\Psi_v(0.011, 1)$ of the first SPOD mode containing fluctuations pertaining to the first sub-harmonic of the forced instability is shown in figure 6.20d. Initial structures are observed further downstream of those in the corresponding eigenfunction in the unforced mixing layer (figure 6.10b), at about $x/\theta_i \approx 250$. Consequently, these fluctuations amplify only after this streamwise station as seen in figure 6.20b. These fluctuations attain maximum amplitude at $x/\theta_i \approx 550$, which is where the amplitude of the forced fundamental instability is seen to attain a low level. Similarly, clear coherent structures pertaining to the second sub-harmonic of the fundamental instability ($St_{\theta_i} = 0.0055$, figure 6.20e) are observed further downstream at $x/\theta_i \approx 550$ as well. When compared to the unforced case, the amplitude of both these sub-harmonic instabilities are significantly lower. This clearly indicates the attenuation of sub-harmonic instability modes when the fundamental instability is forced.

6.4.5.2. First sub-harmonic instability forcing

The results of SPOD applied on the mixing layer forced at the first sub-harmonic of the fundamental instability ($St_f = St_0/2 = 0.011$) is presented in figure 6.21. As before, a sharp peak is seen in the TKE distribution at the forced instability as seen in figure 6.21a. These fluctuations contains the highest energy in the current mixing layer ($\approx 7\%$), with the first SPOD mode accounting for most of this energy. Concentration of energy at modes of fundamental instability is also observed. Finally, as observed in the power spectra of this case presented in figure 6.19b, high concentration of energy is observed in frequency bands around $St_{\theta_i} = 0.015$ and $St_{\theta_i} = 0.026$, here as well. In fact, the strength of fluctuations in the latter frequency band seem to be higher than those at the fundamental KH instability. These two non-harmonic fluctuation bands are discussed separately in section 6.4.5.5.

The vertical component of the eigenfunction $\Psi_v(0.011, 1)$ containing fluctuations of the forced sub-harmonic instability is shown in figure 6.21d and its amplitude along the $y_{0.5}$ trace is presented in figure 6.21b. Unlike in the unforced flow where the initial amplitude of these fluctuations were negligible (figure 6.11), here, the amplitude of this mode at $x/\theta_i = 0$ is $A_v(0.011, 1)/\overline{U} = 0.018$. This is a direct consequence of the momentum input of the DBD plasma actuator. These fluctuations then grow rapidly until $x/\theta_i \approx 50$, beyond which they am-



Figure 6.21.: SPOD decomposition of the mixing layer forced at $St_f = St_0/2 = 0.011$ ($c_{\mu} = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of selected frequencies along the $y_{0.5}(x)$ trace. (c)-(e) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) of these different frequencies (grey - negative; black - positive).

plify gradually and attain their maximum amplitude at $x/\theta_i \approx 320$, close to the location where the growth of this mixing layer came to a temporary halt as seen in figure 6.18b. This is also corroborated by the orientation of the coherent structures in the corresponding eigenfunction. Beyond this streamwise station, a gradual reduction in the amplitude of the mode is observed until about $x/\theta_i \approx 950$, where it reaches a low level. Similar to the previous case with the fundamental instability being forced, here too, the forced sub-harmonic instability attains its maximum amplitude at the same location as observed in the unforced mixing layer (figures 6.10b and 6.11). However, the maximum amplitude is higher compared to the unforced case and the amplitude reaches a low level further downstream as well.

The vertical component of the eigenfunction $\Psi_v(0.022, 1)$ containing fluctuations of the fundamental instability is shown in figure 6.21c. This eigenfunction bears some similarities with the corresponding one in the mixing layer where the fundamental instability was forced (see figure 6.20c). The amplitude of this instability is negligible initially as seen in figure 6.21b, suggesting that these fluctuations are not influenced by the applied forcing. It then grows rapidly until $x/\theta_i \approx 50$ at which station it attains its maximum amplitude in the investigated domain $(A_{v,max}(0.022, 1)/\overline{U} = 0.053)$ and then reduces gradually thereafter. Although, the maximum amplitude of this mode is lower here compared to not only the mixing layer forced at the fundamental instability (which is expectable), but also the unforced flow case.

A comparison of the evolution of the amplitudes of these two instability modes is now made. After the onset, these modes amplify at a similar rate ($\alpha_i \cdot \theta_i \approx 0.0014$) until $x/\theta_i \approx 25$ indicating that these fluctuations are amplifying independently. The amplification rate of the sub-harmonic is reduced from $x/\theta_i \approx 44$, just upstream of where the fundamental instability attains its maximum amplitude. This signifies the onset of sub-harmonic resonance (Monkewitz, 1988). As mentioned in section 2.3.1.2, the works of Monkewitz (1988) and Husain & Hussain (1995) established that the amplitude of the fundamental instability being greater than a critical value of $\alpha_{crit}/\overline{U} = 0.0153 \cdot R^2$ is a pre-requisite for the fundamental and sub-harmonic instabilities to phase-lock and result in the resonant growth of the latter. In the current mixing layer, the amplitude of the fundamental instability at $x/\theta_i = 44$ is $A_v(0.022, 1) = 0.051$, which is significantly higher compared to the critical value and thus is sufficient to establish a phase-lock between the fundamental instability and the forced sub-harmonic. Consequently, as the amplitude of fluctuations corresponding to the fundamental instability begins to reduce, that of the forced first sub-harmonic instability increases. This is representative of the sub-harmonic being reinforced by energy transfer from the mean flow field that is oscillating at the fundamental instability. This leads to resonant growth of the sub-harmonic instability between $x/\theta_i \approx 80$ and 300 at a rate of $\alpha_i \cdot \theta_i \approx 0.0003$. This forced instability mode attains a maximum amplitude of $A_{v,max}(0.011, 1)/\overline{U} = 0.099$.

The eigenfunction representative of the fluctuations pertaining to the second sub-harmonic of the fundamental instability $\Psi_v(0.0055, 1)$ is shown in figure 6.21e. Clear coherent structures are observed only after $x/\theta_i \approx 600$ which corresponds to the location where the current mixing layer begins to grow again after a temporary halt (see figure 6.18b). The maximum amplitude value of this mode is comparable to that in unforced mixing layer and this is attained at $x/\theta_i \approx 750$ which too is not very different from the unforced case. These observations suggest that the amplification of this mode may not be significantly influenced by the applied forcing.

6.4.5.3. First sub-harmonic instability forcing with detuning

The SPOD computations on the mixing layer forced at $St_f = 0.0117$, which is an instability that is marginally detuned (by $0.033 \cdot St_0$) from the first sub-harmonic of the fundamental instability are presented in figure 6.22. Similar to the previous case, highest concentration of TKE is contained in the modes pertaining to the fluctuations of the forced instability as can be seen in figure 6.22a. As observed in the power spectra in figure 6.19c, SPOD modes corresponding to the fundamental KH instability contain negligible energy (two orders-of-magnitude lower than the forced mode), while fluctuations at the first harmonic of the forced instability $St_{\theta_i} = 2 \cdot St_f =$ 0.0235, that is in the frequency band of the fundamental instability is amplified. Furthermore,



Figure 6.22.: SPOD decomposition of the mixing layer forced at $St_f = 0.0117$ ($c_\mu = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of selected frequencies along the $y_{0.5}(x)$ trace. (c) and (d) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) of these different frequencies (grey negative; black - positive).

significant energy is also observed in the non-harmonic frequency bands around $St_{\theta_i} = 0.0154$ and $St_{\theta_i} = 0.0271$. Finally, TKE of SPOD modes at the first sub-harmonic is one order-ofmagnitude lower compared to the forced instability.

The coherent structures corresponding to the forced instability and its first harmonic are presented through the vertical component of the SPOD eigenfunctions in figure 6.22 as is also the amplitude of these modes along the $y_{0.5}$ trace. Structures pertaining to the forced instability (figure 6.22d) attain a vertical alignment with respect to the streamwise direction at $x/\theta_i \approx 270$ which is upstream compared to the unforced mixing layer where the vertical alignment was achieved at $x/\theta_i \approx 320$ (see figure 6.10d). This is unlike the previous case when the first subharmonic was forced, where the structures bore close resemblance to those in the corresponding SPOD mode in the unforced mixing layer. The eigenfunction corresponding to the harmonic of the forced instability $\Psi_v(0.0235, 1)$ is presented figure 6.22c.

The initial amplitude of the forced instability is not negligible but $A_v(0.0117, 1)/\overline{U} = 0.019$ which is similar to the previous case. These fluctuations and those at its harmonic grow at the same rate of $-\alpha_i \cdot \theta_i \approx 0.0014$ until $x/\theta_i \approx 35$, at which point the amplification rate of the forced instability reduces. The amplitude of the modified fundamental (harmonic of the forced instability) at this streamwise stations is $A_v(0.0235, 1)/\overline{U} = 0.042$. While the rapid growth of the harmonic of the forced mode ends at $x/\theta_i \approx 45$, it continues to amplify until $x/\theta_i \approx 150$ although at a reduced rate. This is followed by a rapid decrease in amplitude of these fluctuations until $x/\theta_i \approx 300$, beyond which it reduces gradually and attains a low level at $x/\theta_i \approx 500$. The forced instability begins to amplify from $x/\theta_i \approx 70$ at a rate of $-\alpha_i \cdot \theta_i \approx 0.0004$, which is once again comparable to the previous case with sub-harmonic instability forcing. This instability reaches its maximum amplitude level of $A_{v,max}(0.0117, 1)/\overline{U} \approx 0.097$ at $x/\theta_i \approx 250$, which corresponds to the streamwise location where linear growth of current mixing layer at a higher growth rate ends in figure 6.18d. The amplitude of these fluctuations remains constant at the maximum level until $x/\theta_i \approx 300$ beyond which it gradually reduces and reaches a low level at $x/\theta_i \approx 600$.

The differences observed when comparing the current results to the case of first sub-harmonic instability forcing, support the proposition made in section 6.4.4. Due to the fluctuations being imparted in the splitter-plate boundary layer ensuing on the high-velocity side, it is probable that the forcing temporarily changes the fundamental instability in the mixing layer to the harmonic of the forced instability. This would explain the difference in the spatial organisation of the forced fluctuations here when compared to the unforced flow, which was not the case with forcing the first sub-harmonic. Finally, the growth curve of this mixing layer collapses on to that of the unforced case as observed in figure 6.18d at $x/\theta_i \approx 500$. This streamwise station is close to where the fluctuations of the harmonic ($St_{\theta_i} = 0.0235$) become negligible. This could result in the flow reverting back to its original fundamental instability of $St_0 = 0.022$. Although, the exact mechanism that leads to this sudden increase of thickness in this streamwise region cannot be discerned with the available data.

6.4.5.4. Second sub-harmonic instability forcing

The results of SPOD applied on the mixing layer forced at the second sub-harmonic of the fundamental instability ($St_f = St_0/4 = 0.0055$) is presented in figure 6.23. Fluctuations corresponding to the forced sub-harmonic instability is seen to contain the most energy (figure 6.23a). Corroborating well with the power spectra of this mixing layer shown in figure 6.19c and similar to the previous two cases with sub-harmonic forcing, significant amount of TKE is observed close to $St_{\theta_i} = 0.026$ here as well. Fluctuations corresponding to $St_{\theta_i} = 0.011$ which is the first harmonic of the forced instability are also seen to contain noticeable amount of TKE while fluctuations pertaining to the fundamental instability contain very little energy comparatively.

The eigenfunction $\Psi_v(0.0055, 1)$ of the first SPOD mode that corresponds to the forced subharmonic instability is shown in figure 6.23e. Based on the orientation of the structures in this eigenfunction and the amplitude of the mode shown in figure 6.23b, these fluctuations amplify linearly at a rate of $-\alpha_i \cdot \theta_i \approx 0.0002$ up to $x/\theta_i \approx 600$. This streamwise station is twice that where the first sub-harmonic instability attained its maximum amplitude previously (see figure 6.21b). Beyond this station, the rate of amplification is reduced and the mode attains its maximum amplitude of $A_{v,max}(0.0055, 1)/\overline{U} \approx 0.1$ at about $x/\theta_i \approx 700$. Post this, the strength of this mode continues to reduce until the end of the measured domain and can be expected reach a low level similar to previous cases further downstream. The streamwise station where this linear growth terminates in the case of the second sub-harmonic is twice that observed with the first sub-harmonic instability.



Figure 6.23.: SPOD decomposition of the mixing layer forced at $St_f = St_0/4 = 0.0055$ ($c_{\mu} = 0.22$). (a) Component of TKE contained in each mode normalized with the total TKE (in %). (b) shows the non-dimensional amplitude $(A_v(St_{\theta_i}, 1)/\overline{U})$ of selected frequencies along the $y_{0.5}(x)$ trace. (c), (d) and (e) show the vertical component of the eigenfunction of the first mode $(\Psi_v(St_{\theta_i}, 1))$ of these different frequencies (grey - negative; black - positive).

Fluctuations pertaining to the fundamental KH instabilities, as seen in the previous case, grow rapidly up to $x/\theta_i \approx 50$. It then gradually reduces in strength and reaches a low level at $x/\theta_i \approx 400$, similar to the unforced mixing layer. Finally, until $x/\theta_i \approx 150$ these fluctuations are stronger than the forced sub-harmonic instability. Fluctuations at the harmonic of the forced instability are organised in the current mixing layer as depicted by the structures in the eigenfunction $\Psi_v(0.011, 1)$ in figure 6.23d. Structures pertaining to this mode are observed only after $x/\theta_i = 80$. Based on the variation of the computed amplitude $A_v(0.011, 1)$, these fluctuations amplify and reach a first maxima at $x/\theta_i \approx 300$ and then decay gradually until



Figure 6.24.: Non-harmonic fluctuations at $St_{\theta_i} = 0.026$ and $St_{\theta_i} = 0.015$ in the mixing layer forced at $St_f = St_0/2 = 0.011$ ($c_\mu = 0.22$) (a) Non-dimensional amplitude ($A_v(St_{\theta_i}, 1)/\overline{U}$) along the $y_{0.5}(x)$ trace of the corresponding SPOD modes. Amplitude of other modes shown in figure 6.21b are repeated here for comparison. (b) and (c) show the vertical component of the corresponding eigenfunctions ($\Psi_v(St_{\theta_i}, 1)$).

 $x/\theta_i \approx 500$. It starts to amplify again beyond $x/\theta_i = 500$ and attains its maximum amplitude in the measured domain at $x/\theta_i \approx 900$. However, this amplitude modulation is due to the coherent structures not being centred along the $y_{0.5}$ trace in the entire domain. Thus, the amplitude $A_v(0.011, 1)$ computed along the $y_{0.5}$ trace might not accurately denote the evolution of this instability.

6.4.5.5. Amplification of non-harmonic fluctuations

In all the mixing layers forced with sub-harmonic instabilities, the amplification of certain nonharmonic fluctuations were observed in both the power spectra presented in section 6.4.4 and the SPOD computations being discussed here. For the mixing layer forced with the first-sub harmonic instability, these amplified non-harmonic fluctuations were in frequency bands centred around $St_{\theta_i} = 0.015$ and $St_{\theta_i} = 0.026$. With detuned first sub-harmonic forcing, these correspond to the frequency bands centred around $St_{\theta_i} = 0.0154$ and $St_{\theta_i} = 0.0271$ and finally with second sub-harmonic instability forcing, they were in the frequency band around $St_{\theta_i} = 0.0257$. In order to study the evolution and organisation of these fluctuations, the non-harmonic frequencies observed in the mixing layer forced with the first sub-harmonic instability is analysed here in more detail. The amplitudes computed along the $y_{0.5}$ trace of the first SPOD modes at $St_{\theta_i} = 0.026$ and 0.015 and the corresponding eigenfunctions obtained from applying SPOD over the mixing layer forced with the first sub-harmonic instability are presented in figure 6.24. The amplitudes of the first SPOD modes corresponding to the fundamental, the forced first sub-harmonic and the second sub-harmonic instabilities are reproduced from figure 6.21b and presented in 6.24a.

The maximum amplitude of the fluctuations at $St_{\theta_i} = 0.026$ is the highest among all the modes considered, including the forced first sub-harmonic instability. Correspondingly, the initial amplification rate of these fluctuations are higher compared to the other modes as well. However, they attain their maximum amplitude at $x/\theta_i \approx 50$, similar to the fundamental instability. Coherent structures in the corresponding eigenfunction (6.24b) are observed only until $x/\theta_i = 200$ and the amplitude of these fluctuations reach a low level close to $x/\theta_i = 300$.

Fluctuations corresponding to $St_{\theta_i} = 0.015$ amplify until $x/\theta_i \approx 220$ at a rate similar to the forced instability after the onset of sub-harmonic resonance. Its amplitude then remains almost constant until $x/\theta_i = 300$ beyond which it gradually reduces and attains a low level after $x/\theta_i = 600$. The structures in the eigenfunction (figure 6.24c) however attain a vertical alignment with the streamwise direction at $x/\theta_i \approx 270$.

At this juncture, it is imperative to discuss whether these fluctuations at $St_{\theta_i} = 0.015$ can be a sub-harmonic sideband. As discussed in section 2.3.1.2, sub-harmonic sidebands are known to occur in the frequency band $St_0/2 \pm 0.08 \cdot St_0$ (Monkewitz, 1988; Thomas & Chu, 1989). In this forced mixing layer, apart from the forced first sub-harmonic itself, spectral content within the frequency band $St_f \pm 0.08 \cdot St_0 = 0.011 \pm 0.0018$ is very low (each mode containing about 0.15% of total TKE, see figure 6.21a). Fluctuations at $St_{\theta} = 0.015$ is $0.18 \cdot St_0$ greater than the forced sub-harmonic. This difference is comparable to the width of the frequency band $(0.16 \cdot St_0)$ where sideband excitation can occur. While these fluctuations can be considered to be a sideband at $St_f + 2 \cdot 0.09 \cdot St_0$, a bias is apparent in this scenario as fluctuations at $St_f - 2 \cdot 0.09 \cdot St_0$ are very weak. Furthermore, when the streamwise evolution of the first sub-harmonic is considered, no amplitude-modulation or suppression of the corresponding fluctuations is observed. Thus, these non-harmonic fluctuations seems to not correspond to sub-harmonic sidebands.

6.4.5.6. Brief remarks on important outcomes

Application of SPOD over the acquired PIV time-series of the forced mixing layer (and also the power spectra is section 6.4.4) demonstrate the effectiveness of the spanwise-uniform EFD forcing applied by the DBD plasma actuator in manipulation the spectral content of the mixing layer and consequently, its stability characteristics. Forcing the fundamental instability results in the fluctuations in the mixing layer being locked to this frequency and its harmonics, and attenuates sub-harmonic instability modes. When sub-harmonic instabilities are forced, the corresponding modes are observed to contain significant amounts of the the total TKE of the flow. However, it also leads to the amplification of other non-harmonic fluctuations.

While the reason for this amplification is not known at the moment, it is clear that they do not pertain to sub-harmonic sidebands (especially non-harmonic fluctuations with lesser frequency than the fundamental). In fact, the applied forcing here seems to not excite sub-harmonic

Forced instability		Selected SPOD modes	Number of phases
St_f	f_b (Hz)	$St_{\theta_i} \ (k=1 \text{ to } 20)$	(phase separation)
0.022	300	0.011, 0.022, 0.044, 0.066	$31~(0.065\pi)$
0.011	150	0.011, 0.022, 0.044	$62~(0.032\pi)$
0.0117	160	0.0117, 0.0235, 0.0351, 0.0468, 0.0585,	$59~(0.034\pi)$
		0.0702	
0.0055	75	0.0055, 0.011, 0.022	$126~(0.016\pi)$

Table 6.2.: SPOD modes selected for reconstruction and parameters of phase averaging for the different forced mixing layers.

sidebands. This indicates that the applied EFD forcing is not of a broadband nature but is rather locked to the selected frequency of the burst-modulation cycle and its harmonics. When a detuned first sub-harmonic forcing was applied it was expected to result in the amplification of sidebands as observed in the experiments of Husain & Hussain (1995). However, in this investigation, sideband amplification was not observed in this scenario either. The harmonic of the forced instability (which was in the frequency band of the fundamental instability) acted as the fundamental itself, making the forced mode the first sub-harmonic. This phenomenon is attributed to the method of forcing employed here, which is based on imparting fluctuations into the splitter-plate boundary layer ensuing on the high-velocity side.

Finally, by comparing the streamwise evolution of the sub-harmonic instabilities forced in the different test cases, a few interesting observations can be made. Firstly, the attained amplitude maxima of all three forced sub-harmonic instabilities are almost equal. As demonstrated by Benard & Moreau (2010), fluctuation amplitudes in the flow induced by the DBD actuator when operated with a burst modulation increases with reducing modulation frequency (i.e. f_b here). The constancy of the forced sub-harmonic instabilities' amplitude maxima suggests that the exercised control is not due to the level of imparted fluctuations but rather, the maximum value of the exerted unsteady EFD body force and of course the frequency of the fluctuations.

6.4.6. Investigation of vortex dynamics from reconstructed velocity fields

The effect of the applied EFD forcing on the dynamics and interactions of the coherent vortices within the mixing layer is now desired. As mentioned in section 6.4.1, the PIV acquisition was synchronized with the burst-modulation cycle. This permits phase-averaging of the time-series at different phases of the low-frequency burst modulation. The Q-criterion proposed by Hunt et al. (1988) (see section 3.3.4), is then applied on each of the phase-averaged fields to detect the coherent vortices in the field. This Q-criterion is non-dimensionalised with the maximum value in the respective field.

The phase-averaging can indeed be performed on the acquired PIV vector fields. In this thesis however, time-series reconstructed with select SPOD modes according to equation 3.15 is utilized for this purpose instead. This is motivated by the fact that such a reconstruction allows the selection of instabilities and to analyse their influence on the overall flow dynamics. Furthermore, this exercise also acts as a filter to remove any high-frequency noise from the acquired PIV time-series. The SPOD modes selected for the reconstruction of the velocity fields of different forced mixing layers are tabulated in table 6.2. The number of phases computed through the phase averaging (f_s/f_b) and the separation between adjacent phases are also presented in the table.



Figure 6.25.: Fluctuations of the v-velocity component at different phases $(v'_{\phi} = \overline{v}_{\phi} - \overline{v})$ nondimensionalised with ΔU , along the $y_{0.5}$ trace of the mixing layer forced at $St_f = St_0 = 0.022$ ($c_{\mu} = 0.22$).

6.4.6.1. Fundamental instability forcing

Results from computing the phase-average over the reconstructed time-series of the mixing layer forced with the fundamental KH instability are presented in figures 6.25 and 6.26. The v-velocity fluctuations at different phases of the burst modulation cycle v'_{ϕ} (= $\bar{v}_{\phi} - \bar{v}$) along the $y_{0.5}$ trace is presented in figure 6.25. This velocity field is non-dimensionalised with ΔU . Streaks oriented with a certain angle with the streamwise direction are observed in this figure. Each combination of streaks with positive vertical velocity (red), followed by a negative vertical velocity (blue), represent one spanwise coherent vortex in the mixing layer with the white region between them representing the vortex core (some of them highlighted with dash-dotted (black) line).

The orientation of the streaks correspond to the convection velocity of the spanwise coherent vortices which is computed as follows

$$u_c = \frac{\Delta x}{(\Delta \phi \cdot \pi)} \cdot \frac{St_f \overline{U}}{\theta_i} \cdot 2\pi \tag{6.5}$$

Here, the fraction $\Delta x/(\Delta \phi \cdot \pi)$ can be obtained by measuring the slope of the dash-dotted (black) lines in figure 6.25. For the current forced mixing layer, the convection velocity of the vortices was computed to be 7.54 ± 0.3 m s⁻¹. This is approximately equal to the average velocity of the two streams $\overline{U} = 7.5$ m s⁻¹ which corroborates well with previous experimental observations (Oster, Wygnanski, & Fiedler, 1977). Thus, the applied forcing has no effect on the convection velocity of the spanwise vortices in the mixing layer flow fields.

The Q-criterion computation on the different phase-averaged fields is shown in figure 6.26. The applied EFD forcing results in coherent vortices forming just downstream of the splitterplate trailing edge. All the vortices are separated from their neighbours by $\Delta x/\theta_i = 45.5$ which is equal to the wavelength of the forced fundamental instability ($\lambda_0/\theta_i = \overline{U}/f_0\theta_i = 1/St_0$, see section 6.3.4). Thus the vortex passage frequency is indeed equal to the forced instability as shown by the the power spectra computation in figure 6.19a. No instance of vortex pairing is observed here due to the attenuation of the sub-harmonic instability modes by the applied fundamental KH instability forcing. This corroborates well with previous research efforts, where low energy at sub-harmonic frequencies compared to the fundamental instability resulted in the inhibition of vortex pairing (Oster & Wygnanski, 1982; Ho & Huang, 1982; Husain & Hussain, 1995). This inhibition of vortex interaction results in the negligible



Figure 6.26.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0 = 0.022$ ($c_\mu = 0.22$).

growth of this mixing layer until $x/\theta_i = 200$ seen in figure 6.18a.

The coherent vortices are observed to grow in strength until the vortex at $x/\theta_i \approx 80$. This is more apparent in the profile of the non-dimensional v-velocity fluctuations $(v'_{\phi}/\Delta U)$ along the $y_{0.5}$ trace (black curve, axis on the right) in the figure. This streamwise station $(x/\theta_i \approx 80)$ is close to where the fluctuations pertaining to the forced fundamental instability was observed to attain its highest amplitude (see section 6.4.5.1). The strength of the coherent vortices seem to diminish further downstream and there are no clear vortices noticeable after $x/\theta_i \approx 450$. This suggests that the current mixing layer is not in phase with the burst-modulation frequency of the applied EFD forcing any more. The re-initiation of growth of the current mixing layer observed in figure 6.18a is a consequence of ambient disturbances and pairing feedback, which initiate vortex pairing.



Figure 6.27.: Fluctuations of the v-velocity component at different phases $(v'_{\phi} = \overline{v}_{\phi} - \overline{v})$ nondimensionalised with ΔU , along the $y_{0.5}$ trace of the mixing layer forced at $St_f = St_0/2 = 0.011$ $(c_{\mu} = 0.22)$.

6.4.6.2. First sub-harmonic instability forcing

The dynamics of large-scale spanwise vortices in the mixing layer forced with the first subharmonic instability is discussed here. The SPOD modes selected to reconstruct the time-series and the parameters of the phase-averaging are given in table 6.2. The fluctuations in the vertical component of the velocity field $(v'_{\phi}/\Delta U)$ at different phases along the $y_{0.5}$ trace is shown in figure 6.27. As observed in the previous case, alternate streaks with positive and negative velocity fluctuations oriented at a certain angle with the streamwise direction are observed after $x/\theta_i \approx 250$. The streaks are not arranged in this manner upstream of this location owing to the vortex pairing process. In this region (i.e. $x/\theta_i > 250$), the vortices convect downstream at a velocity of about 8 ± 0.22 m s⁻¹ (computed with equation 6.5) which is comparable to the expected value. This once again demonstrates that the applied sub-harmonic instability forcing does not have a significant influence on the convention speed of the vortices, especially after the completion of the pairing process.

The non-dimensional Q-criterion computed on the phase-averaged fields are presented in figure 6.28. Coherent vortices are observed to form just downstream of the splitter plate. The initial roll-up of vortices lasts until $x/\theta_i \approx 50$ which is close to the streamwise station where the fundamental KH instability attained its maximum amplitude (figure 6.21b). One instance of vortex pairing involving the two vortices (dashed red rectangle in the fields) is observed in the current mixing layer. The interaction of these two vortices along the burst modulation cycle is depicted in the snapshots presented in figure 6.28. Among these two vortices, the weaker vortex is generated by the naturally occurring fundamental KH instability in the flow field during the previous burst-modulation cycle. While the stronger vortex that is upstream, is influenced by the applied forcing. This is clearly evident when the first two vortices in the phase-averaged field at $\phi = 1\pi$, which marks end of the applied forcing, are observed.

The interaction of the forced sub-harmonic instability with the naturally present fundamental KH instability, especially the phase relation between them, results in the roll-up of these two vortices close to one another. The two vortices then undergo pairing due to mutual induction as discussed by Husain & Hussain (1995). The vortex pairing process begins at $x/\theta_i \approx 100$, which corresponds to the location where the amplitude of the forced sub-harmonic instability becomes more than that of the fundamental instability as observed in figure 6.21b. The pairing process occurs as the amplitude of the forced sub-harmonic instability increases linearly and ends close to where the amplitude reaches its maximum value, at about $x/\theta_i \approx 280$. This agrees with the



Figure 6.28.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0/2 = 0.011$ ($c_{\mu} = 0.22$). Snapshots of the vortex pairing process (vortices in dashed red rectangle in the fields) along the low-frequency burst modulation cycle is also shown.

observations of Ho & Huang (1982) who showed that the vortex pairing phenomenon occurs while the sub-harmonic instability grows linearly. The pairing process is completed within two cycles of the burst modulation, thus resulting in the higher growth rate in this streamwise region as seen in figure 6.18b for this mixing layer.

It is worth noting that SPOD modes corresponding to fluctuations in the two frequency bands around $St_{\theta_i} = 0.026$ and 0.015 that were seen to contain considerable TKE in section



Figure 6.29.: Fluctuations of the v-velocity component at different phases $(v'_{\phi} = \overline{v}_{\phi} - \overline{v})$ nondimensionalised with ΔU , along the $y_{0.5}$ trace of the mixing layer forced at $St_f = 0.0117$ $(c_{\mu} = 0.22)$.

6.4.5.2, are not included in the current reconstruction. However, the vortex interaction is still realized which indicates that these modes do not directly impact the vortex pairing process. The wavelength of fluctuations in the frequency band around $St_{\theta_i} = 0.015$ computed by $\lambda(St_{\theta_i})/\theta_i =$ $1/St_{\theta_i}$ ranges between 62 and 74. Interestingly, the distance between the weaker vortex and the stronger vortex downstream (not the pairing vortex-pair) is in that range. The fluctuations in the frequency band around $St_{\theta} = 0.026$ is due to the pairing process as well. This is a result of the passage of the pairing vortices which are displaced such that they are closer to one another from the onset of the mixing layer, resulting in a smaller streamwise wavelength compared to the fundamental instability. It is worth noting that this frequency is the sum of the frequency of the forced instability $St_f = 0.011$ and that of the fluctuations around $St_{\theta} = 0.015$, suggesting a favourable phase relation between these modes. Thus, the fluctuations in these nonharmonic frequency bands are in fact an artefact of the vortex passage frequency induced by the pairing process, rather than an instability. This explains the pairednature of their evolution with the forced sub-harmonic ($St_{\theta_i} = 0.015$ mode) and the fundamental ($St_{\theta_i} = 0.026$ mode) that was observed in section 6.4.5.5.

The coherent vortices show no interaction after the completion of the vortex pairing, which explains the temporary halt in the growth. This is expected to be due to the sub-harmonic resonance mechanism initiated by the imparted fluctuations (Oster & Wygnanski, 1982; Husain & Hussain, 1995). The sub-harmonic resonance renders the forced sub-harmonic instability to have a higher amplitude, especially in comparison with the fundamental instability as seen in figure 6.21b, and leads to a train of vortices that do not interact with each other. In fact, this region is bounded within $1 < RSt_f x/\theta_i < 2$, as proposed by Oster & Wygnanski (1982) based on their observations. Thus, it may be inferred here that, not just the presence of sub-harmonic instabilities, but their relative energy level with respect to the fundamental instability governs the vortex pairing process.

Here too, the vortices become weaker and eventually disappear after $x/\theta_i \approx 600$ (not shown here). The linear growth of the mixing layer after this streamwise station is mainly attributed to the natural onset of other instabilities with frequencies lower than the forced sub-harmonic instability (see for example amplitude of $St_{\theta_i} = 0.0055$ in figure 6.21b).

6.4.6.3. First sub-harmonic instability forcing with detuning

Vortex dynamics when the mixing layer was forced at $St_f = 0.0117 \ (= St_0/2 + 0.033 \cdot St_0)$ is now investigated, through the computation of phase-averaged vector fields from the reconstructed time-series. The fluctuations in the vertical component of the velocity field $(v'_{\phi}/\Delta U)$ at different phases along the $y_{0.5}$ trace is shown in figure 6.29. Here, the alternate streaks oriented at a certain angle with the streamwise direction are observed close $x/\theta_i \approx 200$, after the completion of the vortex paring process. These streaks are only visible until $x/\theta \approx 500$ in the current case. In this streamwise regions $(200 < x/\theta_i < 500)$, the vortices convect downstream at a velocity of about 7.8 ± 0.26 m s⁻¹.

The Q-criterion applied to these phase-averaged fields shown in figure 6.30 reveals the occurrence of pairing of two spanwise vortices. Although the applied forcing is slightly offset from the first sub-harmonic of the fundamental instability, the vortex interaction observed here bears great resemblance to the previous forcing case. In this mixing layer, the phase relation between the forced instability $St_f = 0.0117$ with its harmonic $St_{\theta_i} = 2 \cdot St_f = 0.0235$ that is close to the fundamental instability results in two vortices coming closer to each other compared to their respective neighbours. Here too, the weaker vortex is naturally occurring while the stronger vortex is influenced by the applied forcing as evident through the first two vortices in the phase-averaged field at $\phi = 1\pi$. As before, these two vortices undergo pairing due to mutual induction.

The vortex pairing process begins at $x/\theta_i \approx 100$ and is completed close to $x/\theta_i \approx 250$ (see snapshots at the bottom of figure 6.30). As in the case where the first sub-harmonic instability was forced, this streamwise region corresponds to that where the amplitude of the forced instability overtook that of its harmonic and attained its maximum amplitude after growing linearly, as shown in figure 6.22b. Additionally, this also corroborates well with observations in figure 6.18d, where the growth rate of this mixing layer was higher compared to the unforced flow. The strength of the vortices diminish further downstream which is also indicated by the *v*-velocity fluctuations along the $y_{0.5}$ trace (black line, axis on the right). Finally, no vortices are observed beyond $x/\theta_i \approx 500$ which corresponds to the location where the current mixing layer's growth collapsed on to the unforced flow configuration.

As in the previous case with first sub-harmonic instability forcing, SPOD modes pertaining to fluctuations around $St_{\theta_i} = 0.0271$ and 0.0154 are not included in the reconstruction of these fields as well. Yet, the vortex interaction is well accounted for. Additionally, the distance between the weak vortex and the stronger vortex downstream is once again in the wavelength range of the frequency band around $St_{\theta_i} = 0.0154$ ($\lambda_{\theta_i}/\theta_i = 59$ to 68). This once again shows that the amplification of these non-harmonic fluctuations is an artefact of the vortex pairing process and not any primary instability mechanism. The variation of frequencies of these non-harmonic fluctuation in the different forced mixing layers is in fact expectable, as forcing different instabilities will result in different vortex dynamics and hence different vortex passage frequencies.

6.4.6.4. Second sub-harmonic instability forcing

Velocity fields reconstruction for the mixing layer being forced with the second sub-harmonic of the fundamental instability was carried out with SPOD modes presented in table 6.2. The results of computing the phase average over these fields are presented in figures 6.31 and 6.32. As observed in the previous case, alternate streaks with positive and negative v-velocity fluctuations


Figure 6.30.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = 0.0117$ ($c_\mu = 0.22$). Snapshots of the vortex pairing process (vortices in dashed red rectangle in the fields) along the low-frequency burst modulation cycle is also shown.

oriented at a certain angle with the streamwise direction are observed in figure 6.31 after $x/\theta_i \approx 400$. The streaks are not arranged in this manner upstream of this location owing to the vortex pairing process. The convection velocity of the spanwise vortices observed after $x/\theta_i = 400$ was computed with equation 6.5 and found to be 7.8 m s⁻¹, which is comparable to the average velocity of the two streams \overline{U} .

The non-dimensional Q-criterion computed on the phase-averaged fields are presented in figure



Figure 6.31.: Fluctuations of the v-velocity component at different phases $(v'_{\phi} = \overline{v}_{\phi} - \overline{v})$ nondimensionalised with ΔU , along the $y_{0.5}$ trace of the mixing layer forced at $St_f = St_0/4 = 0.0055$ $(c_{\mu} = 0.22)$.

6.32. A total of four coherent vortices are observed to merge together in the current mixing layer. The first instance of vortex pairing involves the three vortices in the dashed (red) rectangle. The streamwise profile of v-velocity fluctuations in the field pertaining to $\phi = 0$ shows that the amplitude of these fluctuations close to the second vortex, and hence its strength, is lower compared to the vortices adjacent to it. This seems to result in vortex shredding where the stronger vortices on either sides shear the weaker vortex in between them and capture part of the latter. This vortex shredding process along the burst modulation cycle is depicted in the snapshots presented in figure 6.32 (snapshots with red axis, first row). Given that the first vortex (left) is the strongest among the three, it seems to capture a larger part of the central vortex, when compared to the last vortex (right). This vortex interaction begins at $x/\theta_i \approx 80$ and is completed at $x/\theta_i \approx 250$ within one cycle of the burst modulation.

These vortices then interact with a vortex that is upstream (dash-dotted (blue) rectangle). The interaction of these three vortices along the burst modulation cycle is also shown in the snapshots (with blue axis, second row). This vortex interaction begins at $x/\theta_i \approx 250$ and is completed at $x/\theta_i \approx 600$ within two burst modulation cycles. This streamwise station $(x/\theta_i \approx 600)$ approximately corresponds to the location where the linear growth of the forced instability mode was observed to end in figure 6.23b. Furthermore, this also corresponds to the streamwise location where the growth of this mixing layer was reduced as observed in figure 6.18c.

6.4.7. Effect of DBD actuator operating parameters

The discussion until now in this chapter shows that the EFD forcing applied by the DBD plasma actuator, especially the low-frequency burst modulation, manipulates the primary KH instabilities within the plane mixing layer. Hence, it has a profound influence on both, the dynamics and (consequently) growth of the mixing layer. It is now envisaged to analyse the effect of other operating parameters of the DBD plasma actuator on the flow field. Two other parameters that govern the working of the DBD actuator apart for the frequency of the burst modulation are the duty-cycle of the burst modulation and the voltage amplitude (V_{ac}) of the sinusoidal AC waveform. The effect of these parameters on the mixing layer are discussed here. As mentioned previously, the AC frequency (f_{ac}) is responsible for the ionisation process and does not influence the spectral content of the mixing layer in any way. It contributes to only the production of the time-averaged EFD body force and as such will have a similar effect to that of varying the voltage amplitude. Thus, the effect of varying f_{ac} is not investigated here.



Figure 6.32.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0/4 = 0.0055$ ($c_{\mu} = 0.22$). Snapshots of the vortex pairing process (vortices in dashed red and dash-dotted blue rectangles in the fields) along the low-frequency burst modulation cycle is also shown.

The momentum coefficients corresponding to the different cases investigated here are estimated as discussed in section 6.4.1 and are presented in table 6.1 as well.

Low-frequency burst modulation was applied at $St_f = 0.0117$ to excite the corresponding primary KH instabilities, in all the test cases presented here. The choice of forcing this primary instability for this parametric investigation is motivated by the following. These fluctuations correspond to a detuned first sub-harmonic instability. Forcing this instability represents a more "realistic" scenario, as forcing the sub-harmonic cannot always be guaranteed in practical applications. Furthermore, forcing this instability should have lead to the amplification of subharmonic sidebands as demonstrated by the experimental investigation of Husain & Hussain (1995). However, acquired results (see section 6.4.5.3) did not indicate the occurrence of sideband frequencies. Instead, the method of forcing (i.e., into the high-velocity splitter-plate boundary layer) seems to modify the fundamental KH instability temporarily, making the forced instability the first sub-harmonic. Thus, forcing this instability is chosen here to investigate if amplification of sidebands and sub-harmonic amplitude modulation can occur at other forcing conditions.

The results presented henceforth are compared firstly with the unforced mixing layer, but also with the results pertaining to forcing this instability discussed until now. As observed in the SPOD computations (figure 6.22) and phase-averaging (figure 6.30, the influence of forcing this instability is confined to the streamwise region between the splitter-plate trailing edge and $x/\theta_i \approx 500$. Thus, only the near-wake PIV FOV (see figure 3.9) is considered to carry out this parametric study.

6.4.7.1. Burst-modulation duty cycle

The duty cycle of the burst-modulation dictates for how long, within one burst modulation cycle, will the EFD forcing be applied. As commented upon in section 2.1.3, the flow is accelerated during the *on* phase of the burst-modulation cycle, while it relaxes/decelerates naturally during the *off* phase. Varying the duty cycle impacts the time-period of both these phases and thus affects the amplitude of induced velocity fluctuations, when it is measured in quiescent conditions. The amplitude of the induced fluctuations increases with reducing burst-modulation duty cycle.

The results presented and discussed until now pertained to forcing the mixing layer with the low-frequency burst-modulation with its duty cycle set at 50%. Here, the effect of the duty cycle on the mixing layer is analysed by applying the forcing at two burst-modulation duty cycles, 20% and 80%. Thus, the time period (i.e. time for which the DBD actuator is operational and applies a force) per burst-modulation cycle is 1.25 ms and 5 ms for the 20% and 80% duty-cycle cases respectively. The AC voltage amplitude and carrier frequency are held constant at $V_{ac} = 14$ kV and $f_{ac} = 2$ kHz for both cases. The corresponding momentum coefficients are $c_{\mu} = 0.06$ and 0.37 respectively (see table 6.1).

The growth of the mixing layers forced at different duty cycles is presented in figure 6.33. As discussed in section 6.4.3, the onset of the mixing layer is upstream compared to the unforced case for all tested cases. This occurs at $x/\theta_i \approx 15$ (as observed in figure 6.17a) in all the three tested cases, suggesting that burst-modulation duty cycle has no influence on the onset of the forced mixing layer. Growth of the layer post its onset varies with the burst-modulation duty cycle. The mixing layer forced at $c_{\mu} = 0.37$ (80% duty cycle) shows linear growth as soon as the initial roll-up of spanwise vortices with a spreading rate of $\approx 0.37 \cdot R$ which is lower compared



Figure 6.33.: Growth of the mixing layer forced at $St_f = 0.0117$, with varying low-frequency burstmodulation duty cycle.

to the mixing layer forced at $c_{\mu} = 0.22$ (50% duty cycle, $\approx 0.41 \cdot R$). It grows at this rate until $x/\theta_i \approx 280$, post which the curve collapses on to that of the unforced mixing layer. While forcing at $c_{\mu} = 0.06$ (20% duty cycle) results in a temporary halt in the growth from $x/\theta_i \approx 50$ and to ≈ 80 . The mixing layer eventually attains a spreading rate of $\approx 0.43 \cdot R$ until $x/\theta_i \approx 350$, beyond which it grows at the same rate as the unforced mixing layer. This forcing leads to a higher thickness than the unforced case and all other forcing cases discussed until now in the measured spatial domain.

The results obtained from applying SPOD on the fluctuating fields of these two cases are presented in figure 6.34. As seen previously in figure 6.22a, in both these cases, significant TKE is contained within the modes pertaining to the forced instability of $St_{\theta_i} = St_f = 0.0117$. Additionally, amplification of fluctuations at the harmonic of the forced instability ($St_{\theta_i} = 0.0235$) that is close to the fundamental KH instability in the current mixing layer flow configuration, instead of the fundamental instability itself is observed in both these cases as well. In the mixing layer forced at $c_{\mu} = 0.37$ (figure 6.34a), strong fluctuations are also observed around $St_{\theta_i} = 0.0286$ and 0.0169. This is similar to results of $c_{\mu} = 0.22$ forcing (figure 6.22a), where fluctuations about $St_{\theta_i} = 0.0271$ and 0.154 were strong. Interestingly however, with the lower duty-cycle case of $c_{\mu} = 0.06$ (see figure 6.34c), the fluctuations within the mixing layer are locked at the frequency of the forced instability and its harmonics, with no other modes showing any concentration of TKE. Furthermore, unlike all other cases with sub-harmonic instability forcing discussed until now, fluctuations at $St_{\theta_i} = 2 \cdot St_f = 0.0235$ contains more TKE (about 1.5 percentage points) compared to the forced instability.

The coherent structures in the SPOD eigenfunctions of the mixing layer forced at $c_{\mu} = 0.37$ (figures 6.34e and g) bear some resemblance to those in the mixing layer forced at $c_{\mu} = 0.22$ shown in figure 6.22. Structures in the eigenfunction corresponding to the forced instability attain vertical alignment with respect to the streamwise direction at $x/\theta_i \approx 300$ which is slightly downstream compared to the $c_{\mu} = 0.22$ forcing case (at $x/\theta_i \approx 270$) but very close to that in the unforced cases (at $x/\theta_i \approx 320$, see figure 6.10d). The eigenfunction $\Psi_v(0.0235, 1)$ shows no significant different compared to that in the mixing layer forced at $c_{\mu} = 0.22$. The maximum amplitude of these two SPOD modes (figure 6.34b) are lower compared to that observed in the previous case $(A_{v,max}(0.0235, 1)/\overline{U} \approx 0.038, A_{v,max}(0.0117, 1)/\overline{U} \approx 0.065)$. The mode corresponding to the forced instability attains its maximum amplitude at $x/\theta_i \approx 275$ and starts to decay thereafter. This streamwise station is where the growth of the mixing layer collapsed on to that of the unforced flow in figure 6.33.



Figure 6.34.: SPOD computations of the mixing layers forced with a burst-modulation cycle with 80% duty cycle, $c_{\mu} = 0.37$ (figures a, b, e and g) and 20% duty cycle, $c_{\mu} = 0.06$ (figures c, d, f and h).

SPOD computation on the mixing layer forced at $c_{\mu} = 0.06$ show significant differences compared to not only the different cases pertaining to this instability reported until now, but also with forcing other sub-harmonic instabilities discussed until now. Turbulent coherent structures corresponding to selected fluctuations in this mixing layer are presented in figure 6.34f and h and the variation of the amplitude of these SPOD modes are in figure 6.34d. Structures in the eigenfunction corresponding to the forced instability (figure 6.34h) attain a vertical alignment between $x/\theta_i = 300$ and 350. Correspondingly, these fluctuations reach their maximum amplitude of $A_{v,max}(0.0117,1)/\overline{U} \approx 0.117$ at $x/\theta_i \approx 300$. However, significant differences are observed with the fluctuations at $St_{\theta_i} = 0.0235$. Structures corresponding to this mode (figure 6.34f) are not present beyond $x/\theta_i \approx 200$. Although not being forced, these fluctuations amplify rapidly until $x/\theta_i \approx 60$ almost similar to the case where the fundamental KH instability was forced (see figure 6.20b). The attained amplitude maxima is $A_{v,max}(0.0235,1)/\overline{U} \approx 0.324$ which too is comparable to the forced fundamental instability in the section 6.4.5.1. At this streamwise station, amplitude of these fluctuations are much greater than the forced instability



Figure 6.35.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = 0.0117$ while varying the burst-modulation duty cycle.

(by approximately 9 times). It then decays and reaches a low level at $x/\theta_i \approx 320$.

In order to analyse the effect of the burst-modulation duty cycle on the vortex dynamics within the mixing layer, velocity fields are reconstructed according to equation 3.15 with all the SPOD modes given in table 6.2 for this instability (as done with forcing this instability in section 6.4.6). The *Q*-criterion computed on the phase-averaged fields of this reconstructed dataset are presented in figure 6.35.

The vortex dynamics in the mixing layer forced at $c_{\mu} = 0.37$ (figure 6.35a) is not significantly different compared to that forced at $c_{\mu} = 0.22$ (figure 6.30). The applied forcing results in two vortices that are close to one another. The weaker vortex among the two is generated naturally by the flow field during the previous burst-modulation cycle. While the other vortex is stronger due to being influenced by the applied forcing. The interaction of these vortices begins just upstream of $x/\theta_i \approx 100$ and is completed close to $x/\theta_i \approx 250$, within two cycles of the burst



Figure 6.35.: Caption in previous page.

modulation. The stronger vortex among the two mentioned above is in fact a combination of two vortices in very close proximity. With 80% duty cycle, the time for which the plasma forcing is applied by the DBD actuator per burst-modulation cycle is 5 ms, which is greater than the time period of the instability at $St_{\theta_i} = 0.0235$ (3.1 ms). Thus, the flow field produces one vortex that is influenced by the applied forcing as before. The second vortex (the more upstream one) is purely a consequence of the additional time-period of the applied forcing. These vortices amalgamate into one vortex at $x/\theta_i \approx 100$ as seen in the field at phase $\phi = 1\pi$ in figure 6.35a. This streamwise station is close to that where the instability at $St_{\theta_i} = 0.0235$ attains maximum amplitude as observed in figure 6.34b. The combination of these vortex dynamics results in a higher spreading rate of the current mixing layer compared to the unforced flow between $x/\theta_i = 50$ and 250 (see figure 6.33).

The dynamics and interaction of vortices when forcing is applied at $c_{\mu} = 0.06$ with 20% duty burst modulation, shown in figure 6.35b, is significantly different compared to the other two cases of forcing the $St_f = 0.0117$ instability. Based on the distance between the initial vortices, these are found to be organised according to the $St_{\theta_i} = 0.0235$ instability mode and results in



Figure 6.36.: Growth of the mixing layer forced at $St_f = 0.0117$, with varying AC voltage amplitude.

the corresponding fluctuations being more amplified than the forced instability (figure 6.34d). In this case, the time-period of the applied forcing (1.25 ms) is lower than that of fluctuations at $St_{\theta_i} = 0.0235$ (3.13 ms). Thus, one of the two vortices is marginally stronger due to the momentum input. Additionally, these two vortices produced during one burst-modulation cycle of the forced instability are marginally closer to one another, than to their neighbours, which is expected to be due to the phase difference between the forced instability and its harmonic, the $St_{\theta_i} = 0.0235$ mode. Just downstream of $x/\theta_i \approx 200$, where the amplitude of the forced instability became higher than fluctuations at $St_{\theta_i} = 0.0235$, they revolve around each other and amalgamate into one larger vortex. This amalgamation process is completed around $x/\theta_i \approx 350$, resulting in the higher growth rate compared to the unforced flow reported in figure 6.33.

6.4.7.2. AC voltage amplitude

The maximum amplitude value of the fluctuations imparted into the mixing layer flow field by the DBD plasma actuator is dictated by the voltage amplitude (V_{ac}) of the input sinusoidal AC waveform. To carry out this exercise, the frequency of the AC signal was held constant at 2 kHz and so was the duty-cycle of the $St_f = 0.0117$ burst-modulation which was fixed at 50%. Impact of applying EFD forcing with AC voltage amplitudes $V_{ac} = 12$ kV and 16 kV is discussed here. The corresponding momentum coefficients are $c_{\mu} = 0.11$ and 0.35 respectively (see table 6.1). As done earlier, with the burst-modulation duty cycle, these forcing cases are compared to that where the AC voltage amplitude was $V_{ac} = 14$ kV corresponding to a momentum coefficient $c_{\mu} = 0.22$.

The evolution of the vorticity thickness of the mixing layers forced with different AC voltage amplitudes is shown in figure 6.36. As observed previously, the onset of forced mixing layers is shifted upstream to $x/\theta_i \approx 15$ compared to the unforced flow case. Growth of the mixing layer post its onset differs with voltage amplitude as well, although less pronounced compared to varying the duty cycle. Forcing the mixing layer at $c_{\mu} = 0.35$ ($V_{ac} = 16$ kV) results in a similar growth compared to the previous case of $c_{\mu} = 0.22$ forcing case until $x/\theta_i \approx 100$. The current mixing layer then grows linearly at a rate of $0.4 \cdot R$ similar to the $c_{\mu} = 0.22$ case ($0.41 \cdot R$) until $x/\theta_i \approx 300$. Beyond this, the growth curve collapses on to that of the unforced mixing layer. The mixing layer forced at $c_{\mu} = 0.11$ ($V_{ac} = 12$ kV) too grows at a rate of $0.4 \cdot R$ until $x/\theta_i \approx 300$. From here to $x/\theta_i \approx 360$ the spreading rate of the layer is reduced to $0.22 \cdot R$, beyond which the growth is halted until almost the end of the measured spatial domain.



Figure 6.37.: SPOD computations of the mixing layers forced with an AC voltage amplitude of $V_{ac} = 16$ kV, $c_{\mu} = 0.35$ (figures a, b, e and g) and $V_{ac} = 12$ kV, $c_{\mu} = 0.11$ (figures c, d, f and h).

The results obtained by applying SPOD on the two forced mixing layers being analysed here is presented in figure 6.37. In both these cases, fluctuations corresponding to the forced instability $St_{\theta_i} = St_f = 0.0117$ contains the highest concentration of TKE compared to all other modes (see figures 6.37a and c). Furthermore, as seen with all the previously presented cases concerning the forcing of this instability, fluctuations at the harmonic $St_{\theta_i} = 2 \cdot St_f = 0.0235$ that is in the frequency band of the fundamental KH instability are also amplified. Application of forcing at $V_{ac} = 16$ kV ($c_{\mu} = 0.62$) leads to the amplification of other fluctuations as observed with the $c_{\mu} = 0.22$ forcing case (figure 6.22a) and the 80% burst-modulation duty cycle case ($c_{\mu} = 0.37$, figure 6.34a). In the current mixing layer, these fluctuations are centred around $St_{\theta_i} = 0.0257$ and 0.0139. With the application of EFD forcing with $V_{ac} = 12$ kV ($c_{\mu} = 0.11$), similar to the case with 20% burst-modulation duty cycle ($c_{\mu} = 0.06$, figure 6.34b), fluctuations are locked at the frequency of the forced instability and its harmonics, with no other modes showing any concentration of TKE. The coherent structures in the SPOD eigenfunctions pertaining to fluctuations of the forced instability and its harmonic in the forced mixing layers being discussed here are also shown in figure 6.37. In the case of the mixing layer forced with $V_{ac} = 16$ kV ($c_{\mu} = 0.35$), the forced instability grows linearly from $x/\theta_i \approx 80$ to 300 where it attains its maximum amplitude of $A_{v,max}(0.0117, 1)/\overline{U} = 0.082$ and decays gradually further downstream. While fluctuations at its harmonic ($St_{\theta_i} = 0.0235$) amplifies rapidly until $x/\theta_i \approx 50$ ($A_{v,max}(0.0235, 1)/\overline{U} \approx 0.07$), where it stagnates. After $x/\theta_i \approx 100$, these fluctuations decay and reach a low level just downstream of $x/\theta_i = 300$. The structures corresponding to both these instability modes in the mixing layer forced at a voltage amplitude of $V_{ac} = 12$ kV ($c_{\mu} = 0.11$) bear great resemblance to those observed in the mixing layer forced with a 20% burst-modulation duty cycle ($c_{\mu} = 0.06$, figures 6.34f and h). The evolution of the amplitudes of these instability modes presented in figure 6.37d too are similar to those in the 20% duty-cycle case (figure 6.34d). Although, while the maximum amplitude attained by the $St_{\theta_i} = 0.0235$ mode is comparatively lower in the current mixing layer ($A_{v,max}(0.0235, 1)/\overline{U} = 0.269$), that of the forced instability is higher ($A_{v,max}(0.0117, 1)/\overline{U} = 0.135$).

Velocity fields are now reconstructed with all the SPOD modes given in table 6.2 for this instability. The Q-criterion computed on the phase-averaged fields of this reconstructed dataset are presented in figure 6.38. Application of EFD forcing at $c_{\mu} = 0.35$ ($V_{ac} = 16$ kV, figure 6.38a) results in vortex dynamics and interactions that is very similar to the $c_{\mu} = 0.22$ forcing case (see figure 6.30). The applied forcing generates a pair of vortices close to one another, with one being stronger than the other. These vortices begin to undergo pairing through mutual induction at $x/\theta_i \approx 100$, which is completed by $x/\theta_i \approx 300$. Consequently, this results in a higher growth rate of the current mixing layer compared to the unforced flow as observed in figure 6.36.

When the mixing layer is forced at a lower voltage of $V_{ac} = 12$ kV ($c_{\mu} = 0.11$, figure 6.38b), the vortex dynamics is similar to the mixing layer forced with 20% burst-modulation duty cycle ($c_{\mu} = 0.06$) presented in figure 6.35b. The vortices here show the action of two instabilities, $St_{\theta_i} = 2 \cdot St_f = 0.0235$ close to the splitter-plate trailing edge in the rolling-up of vortices and the forced $St_{\theta_i} = St_f = 0.0117$ in the latter half of the measured spatial domain. As before, two vortices with different strengths are produced per burst-modulation cycle of the forced instability. The distance between these vortices gradually reduce along the streamwise direction. Eventually, they come so close to each other that they start revolving around each other and amalgamate into one large vortex. This vortex interaction is completed just downstream of $x/\theta_i = 300$. The newly formed vortices then convect downstream without interacting with one another until the end of the measured domain while reducing in strength.

6.4.7.3. Brief comparison of low-momentum input actuation tests

A brief comment about the application of forcing at $c_{\mu} = 0.06$ (20% burst-modulation duty cycle) and $c_{\mu} = 0.11$ ($V_{ac} = 12$ kV) is necessary. The SPOD computations on both these forced mixing layers showed that the fluctuations within are locked only to the forced instability and its harmonics. No amplification of non-harmonic fluctuations are observed. This is essentially due to the ensuing vortex formation and dynamics, which is similar in both these cases. The phase-averaged fields of both these cases (figures 6.35b and 6.38b) indicate an important feature. When the time-period of the applied forcing is less than that of the fundamental instability (at least when forcing first sub-harmonic instabilities), its effect is similar to the influence of the forcing amplitude. The comparison of these two cases once again point towards the fact that the control authority exercised by DBD actuators operated with burst modulation is primarily



Figure 6.38.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = 0.0117$ while varying the AC voltage amplitude V_{ac} .

due to the maximum-bound of the unsteady force exerted over the flow field and of course the burst-modulation frequency.

6.5. MANIPULATION OF SECONDARY INSTABILITIES THROUGH SPANWISE-MODULATED EFD FORCING

A preliminary investigation is now carried out to study the effect of exerting a spanwisemodulated EFD body force over the same mixing layer configuration. The pursuit here is to take advantage of the spatial nature of the applied forcing to excite secondary instabilities. As introduced in section 2.3.1.3, the streamwise vortical flow structures arising from this instability mechanism, combined with the vortex pairing phenomenon that is moderated by the primary



Figure 6.38.: Caption in previous page.

KH instability mechanism can result in enhanced mixing. The results reported in section 6.4 have demonstrated the capability of the DBD plasma actuator exerting a spanwise-uniform EFD body force to exercise control over the primary KH instability mechanism. The modulated actuator here will be operated with a similar input signal as in the previous chapter (i.e. a burst-modulated sinusoidal AC signal). If successful, the flow control will be realized at a lower input power cost compared to the spanwise-uniform forcing as demonstrated previously in chapter 5 (specifically, see section 5.7), which can be beneficial for practical applications.

The design of this actuators is as reported in section 4.5 (see figure 6.39). It is such that the encapsulated electrode is composed of equi-spaced stems (see figure 4.1b). Each of these stems are $w_S = 3$ mm wide and are 37 mm apart from each other. This results in a spanwisemodulated body force with a wavelength of $\lambda_z = 37$ mm and a spatial duty of $w_S/\lambda_z = 8\%$. The flow conditions and the coordinate/reference system for these investigations are the same as the rest of this chapter (introduced in section 6.1).

A thorough investigation of such a forced mixing layer necessitates the usage of a measure-



Figure 6.39.: PMMA splitter-plate attachment housing the DBD plasma actuator applying spanwisemodulated EFD body force. The two spanwise stations where the PIV fields were captured is represented by the green vertical lines.

ment system that can provide time-resolved three-dimensional velocity vector data, such as tomographic PIV. However, realizing a FOV that has dimensions of even just the near-wake PIV field reported in section 6.4.7 in the x - y directions ($\approx 280 \text{ mm}^2$) and at least 40 mm along the spanwise direction (in order to capture the $\lambda_z = 37 \text{ mm}$ flow organisation) comes with some complications. Firstly, these secondary vortical structures are much smaller than the spanwise-vortices dealt with in the previous section. Thus, it necessitates a tomographic PIV system that has a high enough spatial resolution as well (while capturing the above mentioned FOV), to measure these vortical flow structures. Indeed the amount of data acquired and computational costs are also key issues.

Thus, for this preliminary investigation, such a measurement system was not deployed. Instead, the time-resolved planar PIV set-up introduced in section 3.2.2.4 and utilized until now in this chapter is employed for here as well. Only the near-wake PIV field of view (see figure 3.9) was captured here. However, instead of translating the entire PIV set-up along the streamwise direction to capture the far-wake FOV, it was translated along the spanwise (z) direction to acquire two parallel x - y planes. The first near-wake field was at the centre of one of the stems of the DBD plasma actuator (z = 0 mm). While the second field was between two stems, with the PIV system being translated by (z =) 14 mm. Results pertaining to the statistical convergence and measurement error are the same as that reported earlier in section 6.2, and hence is not repeated here.

6.5.1. Unforced mixing layer

The flow conditions at which the plane mixing layer configuration is realized is the same as given in section 6.1. The high-velocity (top) stream was maintained at $U_2 = 10 \text{ m s}^{-1}$ while the low-velocity (bottom) stream was at $U_1 = 5 \text{ m s}^{-1}$. Thus, the velocity difference between the two streams is $\Delta U = 5 \text{ m s}^{-1}$ and the average velocity of the two streams is $\overline{U} = 7.5$ m s⁻¹. The velocity ratios being $r = U_1/U_2 = 0.5$ and $R = \Delta U/\overline{U} = 0.33$. However, the PMMA attachment fixed to the aluminium (splitter) plate is not the same as the one used for the spanwise-uniform forcing discussed in section 6.4. As mentioned in section 4.5, given than that spanwise-modulated forcing is realized by altering the shape of the encapsulated electrode,



Figure 6.40.: Vorticity thickness of the unforced mixing layer along the streamwise direction at the two spanwise stations. Also shown is the vorticity thickness of the unforced mixing layer measured with the PMMA attachment generating a spanwise-uniform EFD forcing (referred to as UML 2D, reproduced from figure 6.4).

a second PMMA attachment was fabricated for this study. Thus, the condition of the incoming boundary layer and the unforced mixing layer are briefly characterised here once again, before investigating the effect of the applied forcing.

6.5.1.1. Incoming boundary layer conditions

The splitter-plate boundary layers on either sides were not tripped here as well. As done in section 6.3.1, the condition of the splitter-plate boundary layers was checked at x = -1 mm. The shape factors were computed with the velocity data from the splitter-plate surface until $y = 2 \cdot \delta_{99}$. At both spanwise stations (i.e. z = 0 and z = 14 mm), the shape factors of the two boundary layers are $H_1 \approx 2.6$ and $H_2 \approx 2.9$, respectively. This confirms that both boundary layers are in laminar conditions. Additionally, the momentum thickness of the splitter-plate boundary layer ensuing on the high-velocity side, at this location (x = -1 mm) was measured to be about $\theta_i = 0.55$ mm, here as well. Thus, this value is utilized henceforth to non-dimensionalise the spatial coordinates and other length parameters wherever necessary. It is worth noting that the two PIV frames will henceforth be referred to by their non-dimensional spanwise location, $z/\theta_i = 0$ and $z/\theta_i = 25.5$.

6.5.1.2. Unforced mixing layer growth

The vorticity thickness of the ensuing mixing layer is presented in figure 6.40. At both spanwise stations, no increase in vorticity thickness is observed until $x/\theta_i \approx 40$ due to the occurrence of the splitter-plate wake. Beyond this streamwise station, a sudden increase in the vorticity thickness is observed which marks the formation of spanwise-oriented vortices due the roll-up of the shear layer at the fundamental frequency (Michalke, 1965a; Winant & Browand, 1974) and thus the onset of the mixing layer. This is very similar to the previous unforced mixing layer (see figure 6.4, also reproduced in figure 6.40). In the current unforced mixing layer, a linear growth is observed from $x/\theta_i \approx 100$ and until the end of the measured spatial domain. The growth of the mixing layer measured at both spanwise stations are very similar. This shows that the growth of the current mixing layer is controlled by the dynamics of the two-dimensional spanwise vortices arising from the KH instability mechanism.



Figure 6.41.: Non-dimensional, normalised power spectra $(\Phi_N^{v'} \cdot \Delta f/\overline{U}^2 \text{ (dB)}, \Delta S t_{\theta_i} = 3.67 \cdot 10^{-4}, \Delta f = 5 \text{ Hz})$ of the v'-velocity component sampled along the $y_{0.5}$ trace of the current unforced mixing layer at $z/\theta_i = 0$.

The growth of the unforced mixing layer previously considered for the spanwise-uniform forcing is also shown in figure 6.40 (reproduced from figure 6.4). When comparing this with the current unforced mixing layer, a significant difference in the growth curves, especially between $x/\theta_i \approx 100$ and 350, is noticeable. This is expected to be a consequence of the surface finish of the two different PMMA attachments employed. This seems to suggest that development of the mixing layer after the roll-up is very receptive to the surface conditions of splitter-plate. However, this is beyond the scope of the current thesis and is not pressed upon. Given the similarity of the growth curves at both spanwise stations of the current mixing layer, the growth curve at $z/\theta_i = 0$ will henceforth be used as the baseline to compare the impact of the applied spanwise-modulated EFD forcing.

6.5.1.3. Spectral distribution and organisation of fluctuations through SPOD

The power spectra of the time series sampled along the $y_{0.5}$ of the current mixing layer at $z/\theta_i = 0$, normalized according to equation 3.11 and non-dimensionalised with \overline{U}^2 is presented in figure 6.41. Similar to the previous unforced mixing layer (see figure 6.8), fluctuations around $St_{\theta_i} = 0.022$ (300 Hz) is highly amplified here close to the splitter-plate trailing edge. This confirms that the fundamental KH instability in the current flow configuration is the same as earlier and has a Strouhal number of $St_0 = 0.022$ ($f_0 = 300$ Hz). Fluctuations at the harmonic of this frequency $St_{\theta_i} = 0.044$ also contains high energy close to the splitter-plate trailing edge. While sub-harmonic frequencies around $St_{\theta_i} = 0.011$ are amplified further downstream from about $x/\theta_i = 250$ until the end of the measured spatial domain.

This is also confirmed by the SPOD computations carried out on the fluctuating fields of this unforced mixing layer, presented in figure 6.42. These SPOD computations (and those that follow) were performed with the same parameters as reported in section 6.3.5.1 ($N_{fr} = 1900$ with 50% overlap leading to $N_k = 20$ (number of blocks) realizations of the plane mixing layer). At both spanwise stations, maximum turbulent kinetic energy (TKE) is contained within the fluctuations around $St_{\theta_i} = St_0 = 0.022$, as can be observed through figures 6.42a and b. Structures within the eigenfunction $\Psi_v(0.022, 1)$ corresponding to these fluctuations (figures 6.42c and d) are similar to one another. Through the orientation of these structures, it can be concluded that these fluctuations amplify until about $x/\theta_i \approx 80$. They then start to decay and no clear structures are observed downstream of $x/\theta_i = 300$. These structures bear good resemblance



Figure 6.42.: SPOD of the unforced mixing layer. Figures on the left correspond to $z/\theta_i = 0$ (a, c, e and g) and figures on the right correspond to $z/\theta_i = 25.5$ (b, d, f and h). (a) and (b) show component of TKE contained in each mode normalized with the total TKE (in %). (c)-(h) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) at different frequencies (grey - negative; black - positive).

to those observed in the corresponding eigenfunction of the previous unforced mixing layer case (see figure 6.10a). The eigenfunctions of the first SPOD mode corresponding to fluctuations at $St_{\theta_i} = 0.0117$ and 0.011 are also shown in figure 6.42.

6.5.2. Forcing configuration, operating conditions and momentum coefficient

The spanwise-modulated actuator that applies forcing at a wavelength of $\lambda_z = 37 \text{ mm} (\lambda_z/\theta_i = 67.3, \text{ spatial duty } w_S/\lambda_z = 8\%)$ whose design was introduced in section 4.5 (figure 6.39) is employed here to provide the necessary forcing amplitudes. As mentioned at the start of this section, the input electrical signal to operate the actuator is similar to that used in section 6.4 to force the primary KH instability modes. It comprises of a sinusoidal AC waveform and a low-frequency (f_b) square-wave burst modulation. This provides a unique advantage where the employed DBD actuator, while exerting the spanwise-modulated EFD body force to excite the secondary instabilities, also imparts fluctuations that can potentially excite the primary KH

instabilities within the mixing layer.

Through the low-frequency burst modulation, fluctuations were imparted at two different frequencies. The first corresponds to the fundamental KH instability of the mixing layer $St_f = St_0 = 0.022$ whose streamwise wavelength is $\lambda_0/\theta_i = 45.5$ (see section 6.3.4). Thus, in this case the spanwise wavelength being imparted is $(\lambda_z/\lambda_x =)$ 1.48 times the streamwise wavelength of the primary instability. The second low-frequency burst modulation corresponds to imparting fluctuations at $St_f = 0.0117$. This detuned first sub-harmonic instability excitation is chosen here for the same reasons given in section 6.4.7 (also explained in section 6.1). The wavelength of this instability wave along the streamwise direction is $\lambda_x/\theta_i = 85.5$. Thus the ratio $\lambda_z/\lambda_x = 0.79$.

The frequency and voltage amplitude of the sinusoidal AC waveform that causes the ionisation process was $f_{ac} = 2$ kHz and $V_{ac} = 12$ kV respectively for all tested cases. The duty-cycle of the low-frequency burst modulation was also held constant at 50%. At these forcing conditions, the momentum coefficient with spanwise-uniform DBD plasma actuator deployed in section 6.4 is $c_{\mu} = 0.11$ (see table 6.1). The momentum coefficient of the current spanwisemodulated actuation is approximated by multiplying the spatial duty of the actuator with that of the spanwise-uniform actuator, as shown in equation 6.4. This is similar to the procedure followed with the spanwise-modulated actuators employed for swept-wing boundary layer control in chapter 5 (see section 5.4.2).

The momentum coefficient of the spanwise-modulated actuator is estimated by multiplying the spanwise duty of the actuator's construction w_S/λ_z to the corresponding one of the spanwiseuniform actuator as shown in equation 6.4. In the current investigation, the spanwise-modulated DBD plasma actuator is composed of ($w_S =$) 3 mm stems that is spaced ($\lambda_z =$) 37 mm from one another. This results in a spanwise duty of $w_S/\lambda_z = 8\%$. The momentum coefficient of the current forcing configuration is then estimated to be $c_{\mu} = 0.11 \cdot w_S/\lambda_z = 0.009$.

6.5.3. Forced mixing layer growth

Growth of the forced mixing layers, at both the measured spanwise stations, are presented in figure 6.43. In the mixing layer where the applied forcing imparted fluctuations corresponding to the fundamental instability $St_f = St_0 = 0.022$ shown in figure 6.43a, the thickness of the mixing layer at $z/\theta_i = 0$ is significantly greater than the unforced flow. The initial roll-up of vortices itself (until $x/\theta_i \approx 50$) seems to result in a thicker mixing layer. It then spreads at a higher rate ($\approx 0.47 \cdot R$, from $x/\theta_i \approx 150$) compared to the unforced flow until $x/\theta_i \approx 325$, after which the growth is temporarily halted until $x/\theta_i \approx 430$. Beyond this, the growth of the layer, although marginally higher in magnitude, is at the same rate as the unforced flow until the end of the measured spatial domain. The growth of this layer at $z/\theta_i = 25.5$ is significantly different compared to that at the previous spanwise station. It is qualitatively similar to that observed with imparting these fluctuations with spanwise-uniform forcing shown in figure 6.18a. After the initial onset, the growth of the mixing layer is temporarily halted until $x/\theta_i \approx 200$, resulting in a thinner layer compared to the unforced flow. It then grows rapidly ($\approx 0.47 \cdot R$) until $x/\theta_i \approx 450$. Beyond this, the layer grows at a rate similar to the unforced flow.

Imparting fluctuations at $St_f = 0.0117$ results in a mixing layer growth as shown in figure 6.43b. At $z/\theta_i = 0$, the mixing layer grows at a linear rate ($\approx 0.35 \cdot R$) higher than the unforced case from $x/\theta_i \approx 80$ and until ≈ 370 . Beyond this streamwise station it grows at a rate marginally lower than the unforced flow until the end of the domain. However, thickness of the



Figure 6.43.: Growth of the mixing layer when fluctuations pertaining to different primary KH instabilities are imparted using the spanwise-modulated DBD plasma actuator ($c_{\mu} = 0.009$). The unforced mixing layer growth presented here corresponds to that at $z/\theta_i = 0$ from figure 6.40.

mixing layer at this spanwise station is greater than the unforced flow in the entire measured spatial domain. Here too, the layer's growth $z/\theta_i = 25.5$, the growth of the mixing layer is not similar and as continuous as in the previous spanwise station. Between $x/\theta_i \approx 150$ to 310, it grows at a rate of $0.33 \cdot R$. At this point, it starts to grow at almost double the rate $(0.65 \cdot R)$ for a short distance until $x/\theta_i \approx 350$. The growth is then halted until $x/\theta_i \approx 470$ post which the layer seems to grow again. Here, the growth of the forced mixing layer seems to be parallel to the unforced mixing layer until the end of the measured domain.

These results once again demonstrate that the influence of the applied EFD forcing, even when spanwise-modulated, is noteworthy. What holds utmost interest is the compelling difference observed in the growth of the corresponding mixing layers at the two spanwise stations. This is a first indication of a spanwise modulation of the thickness of the forced mixing layer due to the exerted forcing.

6.5.4. Spatial organisation of fluctuations through SPOD

Spatial organisation of the fluctuating fields corresponding to the forced mixing layers is now sort, which is achieved through SPOD computations. These computations were performed on the acquired time-series of the forced mixing layer, at both spanwise stations (as done with the unforced mixing layer in section 6.5.1.3). The parameters chosen for these computations are the same as all the previous SPOD computations and is as given in section 6.3.5.1.



Figure 6.44.: SPOD of the mixing layer forced at $St_f = St_0 = 0.022$ ($c_\mu = 0.009$). Figures on the left correspond to $z/\theta_i = 0$ (a, c, and e) and figures on the right correspond to $z/\theta_i = 25.5$ (b, d, and f). (a) and (b) show component of TKE contained in each mode normalized with the total TKE (in %). (c)-(f) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) at different frequencies (grey - negative; black - positive).

6.5.4.1. Fundamental KH instability forcing

Results obtained from applying SPOD on the time-series of the mixing layer imparted with fluctuations at the fundamental instability $St_f = St_0 = 0.022$ is presented in figure 6.44. The distribution of energy in different SPOD modes compared to the TKE of the flow is shown in figures 6.44a and b. Similar to the corresponding case with spanwise-uniform forcing (section 6.4.5.1), the fluctuations at both spanwise stations are locked at the frequency of the forced fundamental instability and its harmonics. Additionally, no amplification of sub-harmonic instabilities is observed in either of the spanwise stations.

The strongest fluctuations at $z/\theta_i = 25.5$, the spanwise station between two stems of the employed DBD plasma actuator, is at the frequency of the forced fundamental KH instability. The eigenfunction $\Psi_v(0.022, 1)$ at this spanwise station is shown in figure 6.44d. The structures attain a vertical alignment with the streamwise direction at $x/\theta_i \approx 80$. They remain oriented so until $x/\theta_i \approx 300$, beyond which clear structures are not noticeable. The spatial organisation of these fluctuations is very similar to that when this instability is forced with the spanwise-uniform actuator (see figure 6.20c). This indicates that, although the fundamental KH instability is forced with a spanwise-modulated EFD body force, it still influences the two-dimensional primary instability mechanism. The eigenfunction containing fluctuations at the

first sub-harmonic of the fundamental KH instability shown in figure 6.44f too, is similar to that in the spanwise-uniform forcing case. Clear structures are present downstream of $x/\theta_i = 200$ and these structures attain a vertical alignment at $x/\theta_i \approx 450$. This streamwise region is where the mixing layer growth rate was significantly higher at this spanwise station as seen in figure 6.43a.

The coherent structures within the eigenfunction of the imparted fluctuations ($\Psi_v(0.022, 1)$) at $z/\theta_i = 0$ are as shown in figure 6.44c. Here, clear structures are observed only until $x/\theta_i \approx 170$. Additionally, these structures have a different shape compared to those in this eigenfunction in previous cases and also in the field at $z/\theta_i = 25.5$. They seem to be elongated along the positive *y*-direction, with the maxima still remaining close to $y/\theta_i = 0$. Downstream of $x/\theta_i = 50$, these elongations gradually get aligned with the streamwise direction, eventually reaching above the head of the coherent structure downstream. Structures pertaining to fluctuations at the first sub-harmonic too shows some stark differences compared to previous cases. While clear structures are observed only downstream of $x/\theta_i = 200$, these structures initially have two maxima. The maxima at the top seems to eventually disappear beyond $x/\theta_i = 400$. However, it is noteworthy that the TKE of these sub-harmonic fluctuations are significantly low (two orders of magnitude) compared to the imparted fluctuations at the fundamental instability. Thus, the thickness of the mixing layer at this spanwise station being larger is not expected to be due to the vortex pairing mechanism.

6.5.4.2. Detuned first sub-harmonic instability forcing

The results obtained from applying SPOD at both spanwise stations of the mixing layer forced at $St_f = 0.0117$ are presented in figure 6.45. Similar to imparting these fluctuations with spanwise-uniform forcing at these operating conditions ($V_{ac} = 12$ kV, $c_{\mu} = 0.11$, see figure 6.37), the fluctuations in the current mixing layer are locked at the forced instability and its harmonics only, as seen in figures 6.45a and b. Furthermore, fluctuations at the fundamental instability are not amplified in this mixing layer as well. Rather, fluctuations at the harmonic of the forced instability in the frequency band of the fundamental contains significant TKE.

The spatial organisation of fluctuations pertaining to the forced sub-harmonic instability and its harmonic at $z/\theta_i = 25.5$, represented by the eigenfunctions in figures 6.45d and f, is similar to that in the corresponding case with spanwise-uniform forcing (see figures 6.37f and h). Fluctuations at the harmonic of the forced instability amplify until $x/\theta_i \approx 80$ and clear structures are not visible in this eigenfunction downstream of $x/\theta_i = 200$. The alignment of structures in the eigenfunction $\Psi_v(0.0117, 1)$ signify the amplification of the imparted fluctuations until $x/\theta_i \approx 350$ and its decay thereafter. These results once again point to the re-organisation of the two-dimensional flow field according to the imparted fluctuations, although through spanwisemodulated EFD forcing.

The eigenfunction $\Psi_v(0.235, 1)$ depicting structures of the harmonic of the imparted fluctuations at $z/\theta_i = 0$ is very similar to that at $z/\theta_i = 25.5$. However, structures in the eigenfunction $\Psi_v(0.0117, 1)$ at this spanwise station, presented in figure 6.45e, are very different compared to that at $z/\theta_i = 25.5$ and also in other mixing layers forced with this instability. Two distinct types of structures are detected here. The first set of structures are detected just downstream of the splitter-plate trailing edge and until $x/\theta_i \approx 250$. Similar to that in the previous case (but clearer here), the first three structures (until $x/\theta_i \approx 150$) comprise of an extrusion from their head, that extends to the top of the structure downstream. The last two of these structures



Figure 6.45.: SPOD of the mixing layer forced at $St_f = 0.0117$ ($c_\mu = 0.009$). Figures on the left correspond to $z/\theta_i = 0$ (a, c, and e) and figures on the right correspond to $z/\theta_i = 25.5$ (b, d, and f). (a) and (b) show component of TKE contained in each mode normalized with the total TKE (in %). (c)-(f) show the vertical component of the eigenfunction of the first mode ($\Psi_v(St_{\theta_i}, 1)$) at different frequencies (grey - negative; black - positive).

between $x/\theta_i = 150$ and 250 do not show such a feature. This could be due to the fluctuations that these extrusions represent becoming very weak. While it is difficult to ascertain this, the core of these structures seem to attain a vertical alignment at $x/\theta_i \approx 80$, beyond which they remain so aligned.

The second set of structures are more elongated and resemble those at $z/\theta_i = 25.5$, at least in size. Furthermore, the structure just downstream of $x/\theta_i = 250$ is oriented with an obtuse angle with respect to the streamwise direction. These structures attain a vertical alignment at $x/\theta_i \approx 370$ marking the corresponding fluctuations attaining maximum amplitude. These structures consist of two maxima, with the lower maxima roughly aligned with the cores of the previous set of structures. The maxima on top develops gradually, becoming very evident at $x/\theta_i \approx 380$ and disappearing further downstream.

6.5.5. Vortical flow structures and their dynamics

Investigating the vortex dynamics within the mixing layer at both these spanwise stations is now desired. As done in section 6.4.6 (and also in section 6.4.7), velocity fields at both spanwise stations were reconstructed from select SPOD modes according to equation 3.15. The SPOD modes selected for reconstructing the time-series of the different forcing cases are as given in



Figure 6.46.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = St_0 = 0.022$ with the spanwise-modulated DBD plasma actuator ($c_{\mu} = 0.009$). Phase-averaged fields at both spanwise stations (a) $z/\theta_i = 0$ and (b) $z/\theta_i = 25.5$ are presented. Contours in (b) represent $Q/Q_{max} = 0.01$ (green) and 0.1 (black) from the corresponding field in (a). The dash-dotted (black) lines represent (from bottom to top) the $y_{0.05}$, $y_{0.5}$ and $y_{0.95}$ traces at the corresponding spanwise stations.

table 6.2. Phase averaging is then performed along different phases of the burst-modulation cycle and the Q-criterion (section 3.3.4) is employed to detect vortical flow structures. Results pertaining to this exercise is presented here.

6.5.5.1. Fundamental instability forcing

The non-dimensional Q-criterion computation on the phase-averaged fields at both spanwise stations of the mixing layer forced with the fundamental KH instability $St_f = St_0 = 0.022$ using the spanwise-modulated actuator are presented in figure 6.46.



Figure 6.46.: Caption in previous page

6.5.5.1.1. Effect on primary vortices

At the spanwise station $z/\theta_i = 25.5$ which is between two stems of the DBD plasma actuator, the applied EFD forcing results in strong vortices just downstream of the splitter-plate trailing edge as shown in figure 6.46b (coloured contours). These are the spanwise vortices arising from the primary KH instability mechanism. As observed with the spanwise-uniform forcing of this instability mode (see 6.26), these vortices advect downstream without interacting with one another resulting in the negligible growth of the layer at this spanwise station observed in figure 6.43a. The separation between neighbouring vortices is $\Delta x/\theta_i \approx 45.5$ which is the wavelength of the forced fundamental instability. This confirms the conclusions drawn from the **SPOD computation that although fluctuations are imparted here using a threedimensional spanwise-modulated plasma discharge, it is still able to re-organise the two-dimensional flow field that is controlled by the applied burst-modulation cycle.**

At $z/\theta_i = 0$ as well (figure 6.46a), which is at the centre of one of the DBD plasma actuator's stems, the applied EFD forcing generates strong spanwise (primary) vortices just downstream of the splitter-plate trailing edge. As observed in the previous spanwise station, these vortices

advect downstream without interacting with one another, with the separation between them once again being $\Delta x/\theta_i \approx 45.5$. In figure 6.46b corresponding to the spanwise station $z/\theta_i = 25.5$, the contours of $Q/Q_{max} = 0.01$ and 0.1 (green and black respectively) from the corresponding phase-averaged fields at $z/\theta_i = 0$ are also shown. This indicates two important features. The first is that the streamwise location of the primary vortices at these spanwise stations are not identical. The spanwise vortex core at $z/\theta_i = 0$ lags that in the other station and the the amount of this lag increases progressively along the streamwise direction. This demonstrates the spanwise undulation of the primary vortices due to the presence of the secondary vortices which was observed first by Breidenthal (1981, referred to as *wiggle*). This undulation is a result of the translative instability introduced by Pierrehumbert & Widnall (1982). Secondly, while the spanwise vortices are centred along the respective $y_{0.5}$ traces of each spanwise station, they are displaced along the *y*-direction when compared to one another. This could be an indication of the concentration of vorticity at the bottom or top, at the bends of the undulated primary vortices (Nygaard & Glezer, 1991; Rogers & Moser, 1992).

6.5.5.1.2. Effect on secondary vortices

A secondary weak vortex core is observed above every primary vortex at the spanwise station $z/\theta_i = 0$, approximately along the $y_{0.95}$ trace, which is not seen at $z/\theta_i = 25.5$. These are the signatures of the secondary vortices that are superimposed on the primary vortex cores observed by previous research efforts such as Bernal & Roshko (1986), Lasheras & H. Choi (1988), and Nygaard & Glezer (1991) to name a few. These vortices are known to have a hairpin or λ shape (see figure 2.21, reproduced from Lasheras & H. Choi, 1988). As the current PIV measurement is two-dimensional, a section of these vortices, as they loop around the primary vortex cores and where their vorticity is perpendicular to the measured plane, is captured here. Thus, they seem like a second set of point vortices at the top edge ($y_{0.95}$ trace) of the mixing layer at this spanwise stations. It is worth noting that the extrusions of the coherent structures observed in the eigenfunction of the forced instability at this spanwise location (figure 6.44c) seems to be an artefact of the fluctuations induced by these secondary vortices.

At the spanwise station $z/\theta_i = 0$, these secondary structures become apparent at $x/\theta_i \approx 70$ in the field at phase $\phi = 0.5\pi$. However, these vortical structures develop with the primary vortices (Huang & Ho, 1990), or even before the roll-up when forcing is applied (Nygaard & Glezer, 1991). Thus, they are present further upstream in the current mixing layer as well. Additionally, they are clearly identifiable at this spanwise station until $x/\theta_i \approx 250$. Their existence correlates well with the streamwise region where the mixing layer growth was significantly higher in figure 6.43a at this spanwise station. This higher growth is expected to be due to greater entrainment caused by the secondary vortices in this region. No vortical structures are observed here post $x/\theta_i = 300$ which corresponds to where the growth of the mixing layer came to a temporary halt. The absence of the secondary flow structures at $z/\theta_i = 25.5$ clearly results in a growth pattern similar to that of forcing the fundamental instability using a spanwise-uniform actuator (figure 6.26).

6.5.5.2. Detuned sub-harmonic instability forcing

In the case of the mixing layer forced with primary instability at $St_f = 0.0117$, results of the phase-averaging performed over the reconstructed time-series along the low-frequency burst modulation, carried out at both the spanwise stations, are presented in figure 6.47.

6.5.5.2.1. Effect on primary vortices

The phase-averaged fields at $z/\theta_i = 25.5$ (figure 6.47b) once again show the authority of the spanwise-modulated forcing on the dynamics of the spanwise (primary) vortices arising from the two-dimensional KH instability mechanism. Spanwise vortices are generated just downstream of the splitter-plate trailing edge and are organised according to the $St_{\theta_i} = 0.0235$ primary instability. One of the two vortices generated per cycle of the imparted fluctuations is marginally stronger than the other. These vortices come close to one another along the streamwise direction and amalgamate into one vortex. This vortex interaction is completed by $x/\theta_i \approx 400$ which corresponds to the streamwise location where the growth of the current mixing layer came to a halt at this spanwise stations as shown in figure 6.43b. Interestingly, the dynamics of these vortices are very similar to that observed in the mixing layer forced with this instability through the spanwise-uniform actuator being operated at the same conditions ($V_{ac} = 12$ kV, figure 6.38b). This once again confirms the manipulation of the two-dimensional primary KH instability mechanism through these spanwise-modulated plasma actuators.

The non-dimensionalised Q-criterion computed on the phase-averaged fields at $z/\theta_i = 0$ is as shown in figure 6.47a. As in the previous spanwise station, initial vortices are generated and organised according to the fluctuations corresponding to the harmonic of the imparted fluctuations that is in the frequency band of the fundamental KH instability. One instance of vortex pairing involving the two primary vortices generated per burst-modulation cycle occurs which is completed by $x/\theta_i \approx 400$, although marginally upstream compared to the previous spanwise stations.

The contour lines in figure 6.47b represent the levels $Q/Q_{max} = 0.005$ and 0.05 (green and black respectively) from the corresponding phase-averaged field at $z/\theta_i = 0$. Here too, the undulation of the primary vortices along the span is evident. It is important to note that the streamwise wavelength of the forced primary instability is greater than the wavelength of the spanwise-modulation of the EFD body force $(\lambda_z/\lambda_{x,f} = 0.79)$. Initially however, the primary vortices are organised according to the forced mode's harmonic $St_{\theta_i} = 0.0235$ whose wavelength is smaller than that of the spanwise modulation $(\lambda_z/\lambda_{x,2f} = 1.58)$. Thus the spanwise undulation of the primary vortices does occur. This corroborates with the observation of Nygaard & Glezer (1991). However, the pairing process seems to negate this spanwise undulation or at least reduce it to a minimal level. Once the pairing process is complete, the vortices are organised according to the forced primary instability $St_f = 0.0117$. Now, the streamwise wavelength surpasses the spanwise wavelength, thus reducing the spanwise undulation. This forcing condition demonstrates the dependence of the translative instability and the corresponding primary vortex core undulation on the ratio λ_z/λ_x . Whatever be the forced instantiates (both primary and secondary), this ratio must be greater than 1 locally (in a certain streamwise region) for the primary vortex core undulation to occur.

6.5.5.2.2. Effect on secondary vortices

As in the previous case, secondary vortical flow structures are observed along the $y_{0.95}$ trace at the spanwise station $z/\theta_i = 0$ but not at $z/\theta_i = 25.5$. This once again indicates the occurrence of secondary vortices that wrap around the primary vortex cores. Given their hairpin of λ shape, they are not observed at both spanwise stations. The difference in the growth patterns at these spanwise stations observed in section 6.43b, can be attributed to the differences in entrainment rate at these two stations due to the secondary vortex cores.



Figure 6.47.: Non-dimensional Q-criterion (Q/Q_{max}) and non-dimensional v' profile along the $y_{0.5}$ trace (black line, vertical axis on the right) computed on the averaged field at different phases (ϕ) of the burst-modulation cycle for the mixing layer forced at $St_f = 0.0117$ with the spanwise-modulated DBD plasma actuator ($c_{\mu} = 0.009$). Phase-averaged fields at both spanwise stations (a) $z/\theta_i = 0$ and (b) $z/\theta_i = 25.5$ are presented. Contours in (b) represent $Q/Q_{max} = 0.01$ (green) and 0.1 (black) from the corresponding field in (a). The dash-dotted (black) lines represent (from bottom to top) the $y_{0.05}$, $y_{0.5}$ and $y_{0.95}$ traces at the corresponding spanwise stations.

As discussed in section 6.4.6, given that the imparted fluctuations are in the frequency band of the first sub-harmonic, the applied EFD forcing influences one of the two vortices generated per burst-modulation cycle while the other is naturally generated by the flow. Consequently, the secondary vortical structures observed at $z/\theta_i = 0$ too, are evident on top of every second primary vortex core here. The extrusions observed on the head of the first set of coherent structures in the SPOD mode eigenfunction of the forced instability until $x/\theta_i \approx 200$ (figure 6.45e) is an artefact of the fluctuations induced by these secondary flow structures. Corresponding to the second set of coherent structures in the eigenfunction, these secondary vortices are not clearly noticeable after $x/\theta_i = 250$. The vortex pairing process engulfs the secondary vortical structures into the newly formed primary vortex core which will result in small-scale transition and enhanced mixing.



Figure 6.47.: Caption in previous page

6.6. CONCLUDING REMARKS

In this chapter, results pertaining to the manipulation of instability modes within the plane mixing layer configuration, through the spanwise-uniform and modulated EFD forcing applied by a DBD plasma actuator, is presented and discussed. A thorough characterisation of the unforced mixing layer was first carried out. The baseline flow configuration was found to be a canonical mixing layer that agreed well with previous studies. Then, spanwise-uniform EFD forcing was applied to study its influence on the growth, spectral distribution, spatial organisation of the fluctuating fields and finally the vortex dynamics within the mixing layer configuration. Finally, a brief study on the effect of different operating parameters of the DBD plasma actuator on the mixing layer was performed.

It is observed that exciting these instabilities by introducing perturbations into the boundary layer on the high-velocity side of the splitter plate is effective in conditioning the mixing layer flow configuration that ensues downstream. The results demonstrate that the applied EFD forcing, specifically the low-frequency burst modulation, has a significant effect on the generation of spanwise vortices, their subsequent pairing and as a consequence, the growth of the mixing layer. Amplification of the fundamental KH instability results in the inhibition of vortex pairing and consequently a local negligible growth of the ensuing mixing layer. This is due to the attenuation of sub-harmonic instabilities, which reiterates their necessity for the occurrence of vortex pairing (Ho & Huang, 1982; Mansour et al., 1988; Husain & Hussain, 1995). Imparting fluctuations at the first sub-harmonic of the fundamental instability, as expected, results in one instance of vortex pairing involving two vortices that amalgamate into one through mutual induction. The growth rate of the mixing layer is observed to be higher than the unforced case in the streamwise region where the vortex pairing process takes place. The growth stagnates thereafter owing to the forced instability mode being highly amplified due to sub-harmonic resonance. Application of forcing at the second sub-harmonic of the fundamental KH instability, as expected, results in the amalgamation of four spanwise vortices. Most interestingly, one instance of vortex pairing seems to be through the shredding mechanism.

Forcing was also exerted at a frequency that is slightly detuned from the first sub-harmonic instability. The effect of the DBD actuator's operating parameters was also investigated through the forcing of this instability mode as this detuned forcing is more realistic, given the accurate phase-locking of the sub-harmonic instability might not always be possible in practical applications. Contrary to previous experimental observations (Husain & Hussain, 1995), this instability mode did not behave as a sub-harmonic sideband, but rather as a sub-harmonic itself. Forcing this instability resulted in its first harmonic that was in the frequency band of the fundamental instability acting as the fundamental, whatever be the magnitude of the applied forcing. This is attributed to the location of imparting fluctuations in the current study, which is in the high-velocity-side splitter-plate boundary layer. Imparting fluctuations here leads to the applied forcing effectively manipulating the stability characteristics of the mixing layer flow field. This reiterates the findings of Ho & Huang (1982), who demonstrated that the high-velocity boundary layer containing higher shear defines the stability characteristics of the mixing layer.

The parametric investigation that was carried out reveals an interesting trend. While the observed vortex evolution and interaction is different, the authority of these actuators in manipulating the primary KH instabilities is more pronounced at lower forcing amplitudes. In fact, the highest increase in mixing layer thickness with spanwise-uniform forcing is achieved with weakest forcing case investigated. Furthermore, reduction of forcing amplitudes by reducing the AC voltage amplitude (which reduces the maximum force input) and burst-modulation duty cycle (which reduces the time-period of applied forcing) result in similar vortex dynamics. This, and other results point towards the realized instability control being an artefact of the maximal-steady component of the momentum input rather than the fluctuating component of the applied EFD body force.

A preliminary study was also carried out to investigate the effect of spanwise-modulated EFD forcing on the mixing layer flow configuration. The use of DBD plasma actuators provided a unique advantage here. The geometry of the actuator provided the necessary spanwise-modulated force to manipulate secondary instability modes while through the input electric signal, low-frequency fluctuations were imparted to excite primary KH instability modes. The velocity in the flow field was measured at two spanwise stations, the first being at the centre of one of the plasma discharge while the second being between two spanwise discharges.

The acquired results indicate a spanwise modulation of the mixing layer thickness. Weak vortical flow structures are observed at the top edge of the mixing layer in the plane aligned with the EFD forcing location which is not present in the frame between two spanwise discharges.

Furthermore, a spanwise undulation of the primary vortex cores is also observed. This indicates that the applied EFD forcing was successful in manipulating the secondary instabilities in the mixing layer and generate the corresponding vortices that loop around the primary vortex cores.

One of the most interesting outcomes of this investigation however is that, while spanwisemodulated forcing was being exerted here, the low-frequency burst modulation manipulated the two-dimensional KH instability mechanism. By doing so, it controlled the dynamics and interactions of the spanwise primary vortices. This holds significant importance for engineering applications as controlling the primary instabilities and the corresponding vortex interaction can potentially be achieved with significantly lower input power cost. Although, it is important to note the effectiveness of this control strategy depends on the desired outcome. If the motive of the flow control application is to reduce turbulence or avoid transition, this method will not be feasible. The secondary vortices produced will cause small-scale transition after the completion of a number of primary vortex-pairing iterations and as such will render the applied control as completely ineffective. On the other hand, if the end goal is to anticipate turbulence and enhance mixing, deploying spanwise-modulated forcing holds significant advantages.

Perspectives for the future

A recurring theme in the current study, especially pertaining to controlling the primary instability modes, is the effect of the phase relation between the forced sub-harmonic and the fundamental instabilities. Thus, using composite input signals that manipulate both the fundamental and sub-harmonic instabilities simultaneously should be attempted in the future. This exercise will not only help govern the phase relation between the two instability waves, but also open the possibility to analyse the effect of the relative energy levels of these instabilities. As such, it has the potential to achieve more efficient flow/instability control.

When operating the actuator in quiescent conditions with a burst modulated AC signal, the amplitude of velocity fluctuations in the induced flow increases with reducing burst-modulation frequency and duty cycle (Benard & Moreau, 2010). However, mixing layers forced with the spanwise-uniform DBD plasma actuator indicate in many instances that the this flow configuration is more receptive to the steady component of the applied forcing rather than the fluctuating component. This is a key observation that can have an impact on not only this flow configuration but many other flow control applications of DBD plasma actuators. One possible explanation for this is that the flow relaxing during the *off* phase of the burst modulation (see figure 2.5b, reproduced from Benard & Moreau, 2010) could be a feature only in quiescent conditions. When the actuator is immersed into a boundary layer, the flow field only experiences the EFD body force, which is non-existent during the *off* phase of the burst modulation cycle. Obtaining time-resolved measurements of the induced flow field in quiescent conditions and comparing these fields with those acquired when the actuator is immersed into a boundary layer will shed more light into this phenomenon.

With regards to the preliminary investigation with spanwise modulated forcing application, the ability of the actuator to control the primary instability modes and the dynamics of the twodimensional flow field while generating secondary vortical structures must be capitalized upon in the future. This can shed more light on the interaction of the flow structures arising from the primary and secondary instability mechanisms in the mixing layer. Furthermore, due the advent of PIV technology, acquisition of a three-dimensional velocity measurement of the mixing layer flow field must be attempted. Replacing the planar PIV fields in the current methodology with such a measurement technique will provide a complete picture of the mixing layer flow field and can be a useful dataset for future flow control studies.

Finally, the importance of the current experimental investigation also lies in its methodology. The procedure followed here involves using DBD plasma actuators to manipulate the instabilities, employing PIV to acquire spatially and temporally correlated velocity measurements, spectral proper orthogonal decomposition (SPOD) to analyse fluctuations dominating the flow field and finally vortex identification schemes to investigate the vortex dynamics within the mixing layer flow field. It is worthy to point out the versatility of SPOD (as also demonstrated by Towne et al., 2018) in providing the spatial organisation of different fluctuations and simplifying the analysis of their evolution and interaction, which is key to studying such flow configurations that are governed by different instability modes. As such, the current methodology can be deployed in a number of flow configurations that are dominated by frequency-dependent instability mechanisms.

7

CONCLUDING REMARKS & RECOMMENDATIONS FOR THE FUTURE

In this thesis, the proficiency of dielectric barrier discharge (DBD) plasma actuators to exercise instability control over flow configurations whose evolution is governed by hydrodynamic instability mechanisms are investigated upon. Flow instability control pivots around imparting low-amplitude perturbations in order to manipulate the stability characteristics of the flow configuration, so as to control it and realize the necessary beneficial effect. Flow instabilities can be broadly classified into two categories. The first subset includes instabilities that scale according to temporal wavelengths and thus are frequency dependent. Excitations of these instabilities therefore require an unsteady/fluctuating forcing. The second category are instability waves that occur spatially and are defined based on their spatial wavelength/wavenumber. Manipulation of such instabilities necessitates a spatially-modulated, three-dimensional forcing input. Given the inherent unsteadiness of the forcing applied by DBD plasma actuators, these actuators can manipulate the first category of instabilities. Through the variation of the geometry of their construction (especially that of the electrodes), a three-dimensional forcing scenario (although still unsteady) can also be realized with DBD plasma actuator. Hence, it can be capable of manipulating the second category of instabilities.

In order to investigate the capabilities of DBD plasma actuators in this regard, two actuators with different geometric configurations were designed and utilized. The first is the commonly used actuator that exerts an unsteady/fluctuating, two-dimensional, spanwise-uniform electrofluid-dynamic (EFD) body force. The second is a specially designed actuator that exerts a three-dimensional spanwise-modulated EFD body force. These actuators were deployed in two flow configurations that are governed by hydrodynamic stability mechanisms. The first being the swept-wing boundary layer dominated by the inviscid cross-flow (CF) instability mechanism which causes laminar-to-turbulent transition on the swept wing's surface. Control of the stationary CF instability modes that manifest as steady co-rotating vortices along the span of the swept-wing is attempted, which necessitates a spanwise-modulated, three-dimensional EFD forcing. The second flow configuration chosen for this thesis is the plane mixing layer, the dynamics of which is defined by the underlying vortical flow structures and their interaction with one another. The large-scale spanwise vortices are governed by the *Kelvin-Helmholtz* (KH) instability mechanism which manifests itself in this flow configuration as frequency-dependent waves. The secondary instabilities manifest themselves as streamwise vortical flow structures spread along the span of the layer with specific spanwise-wavelength, which once again necessitates the use of spanwise-modulated forcing.

Although plasma actuators have been praised for their non-obtrusive character, conventional construction techniques typically followed in laboratory settings result in electrode thicknesses of tens of microns, which can prohibitive and thus have a destabilising effect on the flow field. This is especially the case here with CF-dominated swept-wing boundary layer, which is extremely sensitive to surface features such as roughness and element-thickness. In this thesis a prototype automated fabrication technique was developed to construct metallic electrodes whose thickness is only a few micrometers. The proposed technique pivots around computer-controlled spraying of conductive silver particles on to the dielectric surface. The technique is capable of producing actuators with extremely thin electrodes. It is also capable of fabricating electrodes with intricate geometries in order to generate three-dimensional EFD forcing whenever necessitated by the flow control application. Additionally, this method is cost-effective and easy to realize in a laboratory setting. Installation of DBD actuators thus produced, on the swept-wing model, showed no influence on the ensuing boundary layer characteristics. This demonstrates the success of the proposed technique in meeting the objectives set during its conception.

Concerning the exerted forcing to manipulate the instabilities in the corresponding flow configurations, the DBD plasma actuators were only partially successfully. The employed actuators demonstrated great authority over the plane mixing layer. The corresponding actuators successfully excited the intended instability modes and manipulated the stability characteristics of the layer. In the swept-wing boundary layer, the deployed actuators were successful in delaying laminar-to-turbulent transition. In fact, the results reported here mark the first realization of a transition delay over a swept wing in an experimental setting with plasma-based active flow control. Although, conclusive proof of plasma-based instability control being the cause of the realized transition delay was not observed when spanwise-modulated actuation was employed to control stationary CF instabilities in the boundary layer. The conditioning of the boundary layer by an array of discrete roughness elements (DREs) seems to render the employed actuators as weak and thus ineffective in forcing the intended instability modes. Contrarily however, in the mixing layer flow field, small-amplitude forcing was observed to be preferred. Thus, the receptivity of the flow field seems to dictate the effectiveness of the deployed instability control.

An important observation comes from the deployment of spanwise-modulated EFD forcing over the two flow configurations. Exerting this forcing over the plane mixing layer revealed its ability to manipulate the primary KH instabilities that govern the two-dimensional dynamics of the layer. This, while exciting the intended secondary instabilities along the span of the layer. This finding has some notable implications for plasma-based instability control in other fluid systems. It demonstrates that while the three-dimensional dynamics of the flow field is influenced by these modulated actuators as intended, the two-dimensional dynamics of the flow can be impacted as well. On the one hand this can be a potential issue as the twodimensional dynamics can be influenced unintentionally. On the other hand, if this is not an issue or preferred, deploying such actuators can be very advantageous as they require significantly lower power expenditure to operate. The positive impact of this lower power consumption is brought forward well in the swept-wing boundary layer investigations. While transition delay was realized with both spanwise-uniform and modulated forcing, the control efficiency of the spanwise-modulated forcing was significantly higher than the corresponding uniform actuation. Similarly in the mixing layer flow configuration, if the end-goal is mixing enhancement, deploying this spanwise-modulated forcing can help realize this objective with very low energy expenditure.

The scope of DBD plasma actuators for engineering applications can also be drawn from the current investigations. The actuators demonstrated good authority over the formation and dynamics of vortical flow structures in the shear layer. As these vortices and the ensuing turbulence are a major source of noise, DBD actuators show potential for noise reduction applications. Secondly, enhanced mixing can also be achieved, especially with the spanwise-modulated actuators, by manipulating the flow field accordingly so as to advance small-scale transition. Finally, in the context of the increasing layer's thickness, the smallest forcing amplitude was the most effective in the mixing layer flow field. In the swept-wing boundary layer, while the achieved transition delay was very small, the realized control efficiency of small-amplitude forcing was significantly high. These observations signal that the low-momentum forcing of DBD plasma actuators can indeed have a beneficial effect in high Reynolds number flows which are encountered more commonly in engineering applications.

Recommendations for the future

As mentioned earlier, any potential flow control system, for many flow configurations that are highly receptive to surface features, needs to adhere to strict roughness and thickness thresholds. While the computer-controlled spraying fabrication technique developed in this thesis is a positive step towards this goal, it should be further improvised, characterised and scaled.

Secondly, the spanwise-modulated DBD plasma actuators that was designed during the course of this thesis shows potential to realize very efficient flow-control systems. This geometric configuration needs some dedicated attention so as to better understand the underlying phenomenon in terms of plasma physics. In the current investigations, the spanwise-wavelength of the applied forcing was within the receptivity band of the corresponding flow configurations. The impact of such a forcing when its wavelength is outside receptivity band must be examined. In such a scenario, it is expected that the forcing will have no effect on the three-dimensional dynamics of the flow configuration. However, if the applied forcing is still capable of influencing the two-dimensional dynamics of the flow field, a significantly more efficient flow control can be realized. Furthermore, in this thesis, all the stems of the encapsulated electrode protruded from a common two-dimensional metallic strip which was grounded. This fixes the spanwise wavelength of the produced EFD body force. Grounding each stem separately can help in varying the EFD force's spanwise wavelength or in imparting sweeping spanwise waves. This scenario must be investigated as it will increase the applicability of these actuators. This will provide an additional control dimension, especially if these actuators are included in closed-loop control algorithms.

Finally, while the current investigations show successful instability control with DBD plasma actuators, these were carried out in low-speed, sub-sonic conditions. Investigation of the efficacy of these actuators in high-speed flow conditions and practical Reynolds numbers is absolutely necessary. Indeed, if a "brute-force" flow-control approach is necessitated by the application, DBD actuators are not a feasible candidate. However, if the flow-control path of manipulating the hydrodynamic stability mechanisms can be taken, DBD actuators are very well suited. Their electrical operation, large frequency-bandwidth of actuation and fast response make them viable device for closed-loop control algorithms. Deploying DBD plasma actuator forcing operated by such intelligent control algorithms has the potential to realize effective and efficient control in high Reynolds number flows.

A

VALIDATION OF SPOD MODE AMPLITUDE COMPUTATION

Amplitude of fluctuations with different frequencies were computed through the corresponding SPOD eigen modes and used for investigating the dynamics of the flow field extensively in this thesis. A brief validation of this computation is carried out here.

To this effect, a set of 2000 composite fields were generated, the assumed sampling frequency being 1 kHz (see figure A.1b for field at t = 0). Here, x represents the horizontal axis and y represents the vertical axis, with the fields being invariant along y. These fields are composed of two sinusoidal waves (equation A.1a) and are defined as shown in equations A.1b and A.1c.

$$f(x,t) = A_w \cdot \sin\left(\frac{2\pi}{\lambda_x}x - 2\pi ft\right)$$
(A.1a)

$$f_1(x,t) = 5 \cdot \sin\left(\frac{2\pi}{0.5}x - 2\pi \cdot 200 \cdot t\right)$$
 (A.1b)

$$f_2(x,t) = 3 \cdot \sin\left(\frac{2\pi}{1}x - 2\pi \cdot 100 \cdot t\right) \tag{A.1c}$$

Finally, an amplitude modulation is superimposed on these waves which results in the waveforms presented in figure A.1a (at t = 0). The first wave (equation A.1b, blue curve in figure A.1a) has a streamwise wavelength of $\lambda_x = 0.5$ units in x. Its amplitude increases from a non-zero value initially and attains a maximum amplitude of $A_w = 5$ units at x = 2.6. Beyond this, its amplitude begins to reduce and becomes zero at x = 8. The second sinusoidal wave (equation A.1c, red curve in figure A.1a) with a unit streamwise wavelength amplifies from x = 2, attaining its maximum amplitude of $A_w = 3$ units at x = 6. It decays further on, becoming zero at the end of the field. These waves are superimposed on one another to generate the final output wave, which is repeated along the y direction to generate the composite fields (see figure A.1b) at t = 0).

The power spectra computed at three-different x-location (at y = 0) are shown in figures A.1c-e. Initially, when only the first sinusoidal wave is present, the power spectra (figure A.1c) shows a peak at f = 200 Hz, which is the carrier frequency of this wave as defined in equation A.1b. Similarly, at x = 9 where only the second sinusoidal wave (equation A.1c) is active, the


Figure A.1.: (a) The two amplitude-modulated sinusoidal waveforms used to generate the composite signal at t = 0. (b) The generated composite field at t = 0. (c)-(e) Power spectra computed at locations in the composite field marked by the black 'x' in (b).

power spectra (figure A.1e) shows a peak at f = 100 Hz. As expected, in the centre of the field at x = 5, the power spectra presented in figure A.1d shows two peaks at f = 200 Hz and 100 Hz due to superimposing the two waveforms.

Spectral POD (see section 3.3.3 for formulation) was then applied upon this composite field dataset, the results pertaining to which are presented in figure A.2. The eigenvalues of the computed modes, normalized with the sum of eigenvalues of all modes, is as shown in figure A.2a. This is representative of the energy contained in each eigen mode compared to the total energy of fluctuations in the field. Unlike in the thesis (see figure 6.9b for example), the first SPOD mode (k = 1) contains all the energy in the field. Furthermore, concentration of energy is observed only at the frequencies of the two sinusoidal waves used to generate the fields. The modes representing fluctuations at f = 200 Hz and f = 100 Hz contain 77% and 23% of the total energy respectively. This is expectable as there is no noise whatsoever in these fields.



Figure A.2.: SPOD computation on the generated composite field dataset. (a) shows the normalized eigenvalue of each SPOD mode representing the energy contained in each eigen mode (in %). (b) and (d) show the the eigenfunctions of the first mode ($\Psi(f, 1)$) at different frequencies (blue - negative; red - positive).



Figure A.3.: Amplitude (A(f, 1)) of the SPOD modes corresponding to the eigenfunctions presented in figures A.2 (b) and (d) along the y = 0. Dashed white curves represent the amplitude of the corresponding sinusoidal waves used to generate the composite field (dash-dotted curves in figure A.1a).

The eigenfunctions of the first SPOD modes at these frequencies are shown in figures A.2b and c. It is important to note that these eigenfunctions illustrate the spatial organization of the corresponding fluctuations only and under no circumstances indicate the strength of these fluctuations at a given location. Fluctuations at f = 200 Hz ($\Psi(200 \text{ Hz}, 1)$), figure A.2b) are observed between x = 0 and 7.8 which corroborates well with the spatial extent of the first sinusoidal wave that is part of the composite fields (blue curves in figure A.1a). Similarly, the eigenfunction $\Psi(100 \text{ Hz}, 1)$ (figure A.2c) indicates that these fluctuations are present only after x = 2 and until the end of the field. This too corroborates well with the location where the second sinusoidal wave introduced fluctuations into the generated composite fields (red curves

in figure A.1a).

The amplitude of these fluctuations is now computed from the corresponding SPOD modes. This is done through the system of equation 6.3 introduced in section 6.3.5. The computed amplitudes are shown in figure A.3. The amplitudes of the modulated sinusoidal waveforms used to generate the composite fields are also presented in the figure for comparison (dashed white curves). As can be observed, the amplitudes computed from the SPOD eigenmodes matches with that of the input waveforms precisely. Thus, this method can be employed to estimate the amplitude of different fluctuations in the corresponding flow configurations to a high degree of accuracy.

B

COMPARISON BETWEEN SPECTRAL AND SPATIAL POD

A form of proper orthogonal decomposition (POD) that has dominated literature on various flow configurations since its inception is the *snapshot* POD technique proposed by Sirovich (1987). The advantage of this method is that it saves computational time for calculating the eigen modes. It is based on decomposing the flow field through the spatial correlation of instantaneous snapshots of the flow. As such, it provides spatially orthogonal eigen modes that are modulated in time by expansion coefficients with random time dependence. It is important to note that the *snapshot* POD technique must not be confused with *classical* POD introduced by Lumley (1967), but it is rather an approximation of the latter (George, 2017; Towne et al., 2018).

B.1. FORMULATION

As done for the SPOD computation, the formulation of which was introduced in section 3.3.3, let us consider a stochastic variable $\tilde{Q}(\Omega, t)$, which in the current work corresponds to the velocity fluctuations in the acquired PIV time-series. Here, Ω represents the measurement spatial domain and t is time. In the case of spatial *snapshot* POD, each instantaneous vector field is considered to be a realization of a stochastic process and as such, temporal correlation is not accounted for directly. Thus, instead of the cross-spectral tensor computed at each frequency for SPOD (equation 3.13), a spatial correlation tensor is defined here as

$$C(\Omega, \Omega') = \widetilde{Q}^*(\Omega)\widetilde{Q}(\Omega) \tag{B.1}$$

where * represents the complex conjugate. The integral eigenvalue problem to be solved to decompose the flow field then reduces to

$$\int_{\Omega} C(\Omega, \Omega') W(\Omega') \zeta(\Omega') d\Omega = \gamma \zeta(\Omega)$$
(B.2)

with $\zeta(\Omega)$ being the spatial eigenfunctions and γ being the corresponding eigenvalues which is representative of the contribution of a particular mode to the total TKE of the flow field. The weight $W(\Omega)$ is a positive-definite Hermitian matrix.



Figure B.1.: Component of TKE contained in each spatial POD mode normalized with the total TKE of the mixing layer forced at $St_f = 0.011$.

The orthogonality of the eigenvectors provides a complete basis for \tilde{Q} . This orthogonality can be exploited to compute the temporal coefficients a(t), as it implies that these coefficients are only dependent on the spatial modes $\zeta(\Omega)$. Thus, the flow field can be expanded according to equation B.3a. As such, the expansion or temporal coefficients $a_k(t)$ can be computed according to equation B.3b with k being the mode index.

$$\widetilde{Q}(\Omega,t) = \sum_{k=1}^{\infty} a_k(t) \zeta^k(\Omega)$$
(B.3a)

$$a_k(t) = \int \tilde{Q}(\Omega, t) \zeta^{k*}(\Omega) d\Omega$$
 (B.3b)

B.2. RESULTS

A comparison between spectral POD (SPOD) extensively used in the current thesis, and spatial POD computed using the *snapshot* technique (Sirovich, 1987) is presented here. To this end, spatial POD was applied on the full set of PIV vector fields of the mixing layer forced with the first sub-harmonic of the fundamental KH instability $St_f = St_0/2 = 0.011$. The results presented here will be compared to that obtained through spectral POD, discussed in section 6.4.5.2. As the influence of forcing this instability is confined to the streamwise region between the splitter-plate trailing edge and $x/\theta_i \approx 500$, only the near-wake PIV field of view (see figure 3.9) is considered here.

The contribution of the first 20 eigen modes to the total TKE of the flow field $(\gamma_k / \sum_k \gamma)$ is shown in figure B.1. The first two modes, which contain the most energy of all modes, contribute only about 6% each to the total TKE. Results pertaining to these modes are presented in figure B.2. The similarity of their relative energies is indicative of these modes being paired. This is also evident from the phase portrait of the corresponding temporal coefficients in figure B.2a which shows the temporal organisation of these modes. The circular shape of the phase portrait suggests that the temporal coefficients of these modes are of the same amplitude and have a phase difference of quarter of a wave $(\pi/2)$. This phase quadrature is also demonstrated by the vertical component of the spatial eigenfunctions of these modes ($\zeta_v^{k=1}$ and $\zeta_v^{k=2}$) presented in figures B.2c and d respectively. While the structures within both modes are very similar,



Figure B.2.: Results pertaining to POD eigen mode pair 1-2 of the mixing layer forced at $St_f = 0.011$. (a) Phase portrait of the temporal (expansion) coefficients $a_{k=1}$ and $a_{k=2}$. (b) Normalized power spectra of these temporal coefficients. (c) and (d) shows the vertical component of the eigenfunctions $\zeta_v^{k=1}$ and $\zeta_v^{k=2}$ respectively (grey - negative; black - positive).

a translation of approximately a quarter of a wavelength is noticeable. This phase quadrature is a necessary condition for these POD modes to represent a periodic phenomenon such as the convection of spanwise vortices in the current flow configuration.

The normalized power spectra (normalized by equation 3.11) of the temporal coefficients of the two modes is presented in figure B.2b. While dominant fluctuations in these modes correspond to the frequency band around $St_{\theta_i} = 0.026$, a peak in the power spectra at the forced first sub-harmonic and a smaller peak at the frequency of the fundamental instability are also evident. However, the distance between structures of the same sign in the eigenfunctions is about $\Delta x/\theta_i \approx 38$ which is equal to the wavelength of fluctuations at $St_{\theta_i} = 0.026$. These eigenfunctions bear great similarity to that computed through SPOD $\Psi_v(0.026, 1)$ presented in figure 6.24b. Structures have an obtuse-angle orientation with the streamwise direction initially indicating a net transfer of energy from the mean flow to the fluctuations pertaining to this mode (Rajaee et al., 1994). They attain a vertical alignment at $x/\theta_i \approx 60$ indicating the corresponding fluctuations attaining their maximum amplitude. Furthermore, no clear structures are noticed after $x/\theta_i = 200$.



Figure B.3.: Results pertaining to POD eigen mode pair 3-4 of the mixing layer forced at $St_f = 0.011$. (a) shows the phase portrait of the temporal (expansion) coefficients $a_{k=3}$ and $a_{k=4}$. (b) shows the normalized power spectra of these temporal coefficients. (c) and (d) shows the vertical component of the eigenfunctions $\zeta_v^{k=3}$ and $\zeta_v^{k=4}$ respectively (grey - negative; black - positive).

POD modes k = 3 and 4 are the second most dominant modes in this forced mixing layer, containing 4.7% of the total TKE of the flow each, as shown in figure B.1. The pairing of these two modes and the phase quadrature between them is apparent from the phase portrait of the corresponding temporal coefficients shown in figure B.3a. The normalized power spectra of these temporal coefficients (figure B.3b) indicate that the dominant fluctuations in this mode-pair correspond to the forced first sub-harmonic. Spectral content in a frequency band centred around the forced sub-harmonic is also seen to be high. Additionally, spectral content is also observed in the frequency band around $St_{\theta_i} = 0.026$, although it is significantly lower in magnitude compared to the dominant fluctuations.

The vertical component of the eigenfunctions of this mode-pair are presented in figures B.3c and d. The distance between structures with the same sign in these eigenfunctions, especially after $x/\theta_i = 200$, is approximately equal to $\Delta x/\theta_i \approx 90$ which corroborates well with the wavelength of the forced sub-harmonic instability ($\lambda_f/\theta_i = 1/St_f = 91$). Here too, the structures bear a great resemblance to those in the corresponding eigenfunction $\Psi_v(0.011, 1)$, computed through the application of SPOD on these vector fields, presented in figure 6.21d. The structures

within the eigenfunctions here are oriented with an obtuse angle to the streamwise direction initially indicating the amplification of the corresponding fluctuations through energy transfer from the mean flow. They attain a vertical alignment at $x/\theta_i \approx 320$ (in $\zeta_v^{k=4}$) signifying the fluctuations reaching their maximum amplitude. Beyond this streamwise station, the structures develop an acute-angle orientation with the streamwise direction until the end of the shown spatial domain.

B.3. DISCUSSION

Spectral POD provides eigen modes that are coherent in both space and time. On the other hand, with spatial POD computed through the *snapshot* technique, the time dependence of the velocity fluctuations (which is the stochastic variable here) is not accounted for, as snapshots of the flow at different instances of time are each considered to be a member of a collection of spatially-dependent fields. Thus, the sequential ordering of the snapshots is lost and as such, the temporal evolution of the flow fluctuations has no impact on the POD eigen modes. Due to this, the computed POD modes do not intrinsically represent structures that evolve coherently in time but only in space. Although, one must not assume that spatial POD modes can never exhibit space-time coherence. An example to this effect is the experimental work of Rajaee et al. (1994), who applied spatial POD on velocity data of a forced mixing layer acquired through hot-wire measurements. The mixing layer was forced acoustically at two frequencies simultaneously, that of the fundamental KH instability and its first sub-harmonic. Their results demonstrate that the mode-pairs 1-2 and 3-4 capture a single frequency each, that of the forced sub-harmonic and fundamental instabilities respectively.

Contrarily in the current work, the power spectra of the temporal coefficients of the first four modes (figure B.2b and B.3b) indicate that these modes are not pertaining to just one particular frequency, but represent a combination of different fluctuations. Equations B.4, reproduced from Towne et al. (2018), show the relationship between spatial (with mode index j) and spectral POD eigen modes (with mode index k).

$$\zeta^{j}(\Omega) = \int_{-\infty}^{\infty} \sum_{k} \frac{\Lambda(f,k)}{\gamma_{j}} c_{jk}(f) \Psi(\Omega, f, k) \cdot df$$
(B.4a)

$$\gamma_j = \int_{-\infty}^{\infty} \sum_k \Lambda(f,k) \left| c_{jk}(f) \right|^2 \cdot df$$
(B.4b)

Here, the variable $c_{jk}(f)$ is the inner product of some spatial POD mode $\zeta^{j}(\Omega)$ and SPOD spatial mode $\Psi(\Omega, f, k)$. These equations show, that each spatial POD mode can be interpreted as a combination of different spatio-temporal SPOD eigen modes. As seen in the current spatial POD results, the organisation of the structures itself might represent fluctuations of a particular frequency that is dominant in those modes. However, the contribution of other spatio-temporal coherences at different frequencies is manifested within the temporal coefficients as observed in the presented power spectra.

This brief discussion demonstrates that each spatial POD eigen mode can represent flow fluctuations at different time scales. In flow configurations such as the mixing layer, where flow fluctuations arising from different instabilities play a key role, this feature of spatial POD computations complicates the investigation and hence, is a drawback. On the other hand, spectral POD computations while decomposing the flow field into different spatial modes, also decouples the different time scales successfully. As such, it is more suitable for investigating such flow configurations where flow dynamics arising from different fluctuations have to be studied for understanding the underlying phenomenon better.

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