

Investigation of hypersonic flow over double-cone geometries with varying nose bluntness

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(Received 5 August 2025; revised 3 February 2026; accepted 15 March 2026)

The hypersonic flow over 30°–50° double-cone configurations with three nose bluntness levels was experimentally investigated at Mach 6. High-speed schlieren photography, pressure sensors and pressure-sensitive paint were used to examine both global flow patterns and unsteady dynamics at a transitional Reynolds number. The experimental results indicate that the size of the separation region at the cone junction increases with increasing nose bluntness. Type V shock–shock interactions were observed in all three configurations, while the shock wave structures in the region below the triple point exhibited two patterns: Mach shock wave reflection in the sharp and small-blunt-nose cases, and regular shock wave reflection in the large-blunt-nose case. Spectral analysis of high-speed schlieren sequences revealed two types of unsteadiness across all cases: low-frequency shock oscillations and high-frequency unsteady structures along the boundary of supersonic jet on the second cone. For the low-frequency unsteadiness, shock oscillations displayed a broadband nature in the sharp and small-blunt-nose configurations, while a dominant frequency of approximately 2 kHz was observed in the large-blunt-nose case, characterised by shock motion and bubble breathing – an observation not experimentally reported before. Additionally, spectral analysis of wall pressure contours indicated that the low-frequency unsteadiness was primarily characterised by axisymmetric modes for all configurations. Global stability analysis and resolvent analysis further demonstrated noise-amplifier behaviour in all configurations, and the dominant low-frequency unsteadiness in the large-blunt-nose case is attributed to modal resonance induced by environmental noise.

Key words: hypersonic flow, shock waves, boundary layer separation

1. Introduction

Shock wave–shock wave interactions (SSIs) and shock wave–boundary layer interactions (SWBLIs) are prevalent phenomena in supersonic/hypersonic internal and external flows (Edney 1968; Green 1970). These interactions generate complex shock wave structures, significant airflow distortion, reduced total pressure and harmful unsteady aerodynamic forces and heat transfer, all of which significantly impact aircraft performance (Dolling 2001; Herrmann & Koschel 2002; Babinsky & Ogawa 2008; Détery & Dussauge 2009; Krishnan, Sandham & Steelant 2009; Babinsky & Harvey 2011). Therefore, deep comprehension of SSIs and SWBLIs is of significant importance to the design of hypersonic flight systems.

Investigations into the underlying mechanisms of SSIs and SWBLIs have frequently employed simplified canonical geometries, including the planar configurations like compression ramps, shock-impingement flat plates, double wedges (Schrijer *et al.* 2005; Détery & Dussauge 2009; Hashimoto 2009; Babinsky & Harvey 2011), alongside axisymmetric geometries such as spiked cylinders, double cones and cylinder–flare configurations (Maull 1960; Holden *et al.* 2003; Hornung, Gollan & Jacobs 2021; Lugin *et al.* 2021). A critical limitation of the planar geometries is the spanwise flow non-uniformities due to sidewall effects or flow leakage under the conditions without sidewalls (Hawboldt, Sullivan & Gottlieb 1995; Bleilebens & Olivier 2006; Grossman & Bruce 2018; Liu *et al.* 2024). In contrast, axisymmetric geometries inherently mitigate these issues, leading to their prominence as subjects for extensive experimental and numerical studies.

Experimental investigation of hypersonic flow separation on axisymmetric configurations can trace back to Maull's schlieren and shadowgraph studies of spiked bodies (Maull 1960). For hypersonic flows over the double-cone configuration, the characteristic features – oblique-bow shock interactions, triple points, shock-induced separation, laminar–turbulent transition and shear-layer instabilities – have been identified through extensive experimental and numerical efforts over decades. Early experiments by Wright *et al.* (2000) identified distinct Type VI and Type V SSIs on double cones in a Mach 8 blowdown tunnel. During the same period, Olejniczak, Candler & Hornung (1997) conducted a series of double-cone experiments in a shock tunnel using nitrogen at very high enthalpies of 27–31 MJ kg^{−1} to test the thermochemical models applied in numerical simulations. Subsequently, comprehensive experimental tests in LENS I and LENS XX tunnels yielded detailed datasets of the mean flow structures, pressure distributions and heat flux of the double-cone flows in air, nitrogen and oxygen across extensive Mach and Reynolds numbers, enabling rigorous computational validation (Holden 2003, 2015; Holden *et al.* 2003, 2008, 2018). However, computational validation succeeded only partially at low enthalpies. At high enthalpies, involving significant vibrational excitation and chemistry, computations substantially underestimated separation size (Knight *et al.* 2012). Additionally, in some high Reynolds number circumstances with Type V SSIs, one possible reason for the discrepancy between the experimental and numerical results is the misprediction for the turbulent behaviour (Sinha, Wright & Candler 1999; Holden 2015).

Regarding the factors influencing the flow characteristics of double-cone flows, the properties of the incoming flows play a significant role. Experiments conducted under both low- and high-enthalpy conditions with various Mach numbers showed that the size of the separation bubble increases with the Reynolds number for double-cone flows with a fully laminar interaction region (Olejniczak *et al.* 1997; Wright *et al.* 2000; Holden *et al.*

2008). Moreover, Holden *et al.* (2003) reported that, under 10 MJ kg^{-1} conditions, higher pressure and heat transfer peaks were detected in air rather than in pure nitrogen, due to partial dissociation of air while nitrogen remained largely undissociated. In contrast, the pressure and heat transfer distributions at a relatively lower-enthalpy condition of 5 MJ kg^{-1} were slightly influenced. Additionally, the properties of the double-cone model, particularly the bluntness of the nose, significantly influence flow characteristics. Holden (1997) demonstrated that the flow structures and aerothermal loads on blunted double-cone configurations are sensitive to thermochemical effects. Their study also observed that for small nose radii, the hypersonic flow remained attached both upstream and downstream of the separation region; conversely, increased nose bluntness caused the separation bubble to extend from the nose to the shoulder of the second cone. Recently, Hao & Wen (2020) conducted numerical simulations of hypersonic flows over blunted double-cone configurations. Their study identified a critical nose radius, below which the flow characteristics remained relatively constant. For nose radii exceeding the critical value, the separation region exhibited an initial increase followed by a decrease as the nose radius increased – a trend consistent with experimental and numerical studies on hypersonic compression-ramp flows featuring blunt leading edges (Holden 1971; Gray & Rhudy 1973; Cao *et al.* 2021a). Hao & Wen (2020) further proposed a prediction for this critical radius, based on the pressure recovery correlations of Blick & Francis (1966) for spherically blunt cones, which aligned well with both numerical and experimental observations (Holden *et al.* 2006). Although studies on the effects of bluntness specifically on boundary-layer transition over double cones are scarce, existing experimental and numerical research on hypersonic flow over circular cones has established that nose bluntness significantly influences transition (Stetson 1983; Marineau *et al.* 2014; Paredes *et al.* 2019). Compared with sharp cones, small bluntness delays transition; however, a transition-reversal phenomenon occurs when the bluntness exceeds a critical value. Jewell & Kimmel (2017) analysed transition on an 8° cone at Mach 6 using parabolised stability theory and reported that as nose bluntness increased, the computed N-factor at the experimental transition location dropped below the threshold for Mack's second mode, indicating a shift away from this conventional modal instability as the dominant mechanism. Paredes *et al.* (2019), in their overview of the NATO STO AVT-240 collaborative investigation on hypersonic boundary-layer transition, linked the transition reversal at high bluntness to non-modal disturbance amplification.

Unsteady flow structures represent another significant phenomenon in double-cone flows. Jagadeesh *et al.* (2003) observed large-amplitude shock oscillations in Mach 6.99 dissociated air flows over double-cone configurations with second cone angles ranging from 65° to 70° , using high-speed schlieren and double-exposure holographic interferometry. Through numerical investigations of supersonic double-cone flows across various inflow and geometric parameters, Hornung *et al.* (2021) revealed that configurations featuring small first-cone angles and large second-cone angles exhibit inviscid-dominated unsteady pulsating flows. Their extensive parametric study established an unsteadiness boundary defined by three independent dimensionless parameters. Employing a time-accurate direct simulation Monte Carlo method, Tumuklu *et al.* (2018a,b) performed numerical computations for 25° – 55° double-cone flows at Mach 16 under varying Reynolds numbers. At the highest Reynolds number examined in their study, the bow shock wave displayed oscillations at approximately 2 kHz. Meanwhile, Kelvin–Helmholtz instabilities emerged within the shear layers on the second cone, exhibiting dominant frequencies ranging from 45 kHz to 70 kHz. More recently, Posladdek *et al.* (2024) experimentally captured small-amplitude unsteadiness in Mach 7 flows over 25° – 35° and 25° – 55° double-cone models using high-speed schlieren. The unsteadiness

analysis conducted using the spectral proper orthogonal decomposition (SPOD) method revealed that for the 25° – 35° configuration, the Type VI SSI was primarily encompassed in a lower broadband frequency region. Conversely, in the 25° – 55° double-cone flow, while the flow structures associated with the Type V SSI dominated the low-frequency region, higher-frequency structures were observed, with frequencies reaching beyond 150 kHz in the underexpanded supersonic jet at the shoulder of the second cone.

Linear stability analysis is a widely employed method for identifying the origins of unsteadiness in laminar and transitional flows. The pioneering work by Huerre & Monkewitz (1990) classified flow unsteadiness into two categories: flows can be characterised as oscillators exhibiting absolute instability or as noise amplifiers that filter and amplify ambient disturbances, which are found to be globally stable and unstable, respectively, through global stability analysis (GSA). In supersonic and hypersonic flows over two-dimensional or axisymmetric geometries, laminar separation bubbles generated by sufficiently strong shock–boundary-layer interactions (exceeding a critical threshold) have been shown to be globally unstable, sustaining intrinsic three-dimensional instabilities without external perturbations. For shock-impingement flows at supercritical shock angles, Robinet (2007) identified a Hopf bifurcation associated with low-frequency unsteadiness in Mach 2.15 flows. Hildebrand *et al.* (2018) observed stationary global modes generating streaks downstream of the reattachment of the laminar separation bubble in Mach 5.92 flows. Similarly, research on hypersonic compression corner, double wedge and double cone configurations has consistently demonstrated the emergence of stationary and oscillatory unstable modes when deflection angles exceed critical values (Sidharth *et al.* 2017, 2018; Cao *et al.* 2021b; Hao *et al.* 2021, 2022). In contrast, for globally stable flows that act as noise amplifiers, resolvent analysis, which examines the response of flow to external disturbances, has been extensively applied to shock-induced separated flows. Dwivedi *et al.* (2019) conducted a resolvent analysis on a Mach 8 compression corner flow and identified that the optimal response of the separated flow to upstream forcing corresponds to streaks downstream of the laminar separation bubble. This finding was further corroborated by subsequent studies conducted by Hao *et al.* (2023) on Mach 7.7 compression corner flows with various ramp angles. In a laminar shock-impingement flow at Mach 2.15, Bugeat *et al.* (2022) found that resonances to three-dimensional disturbances exhibited both modal and non-modal mechanisms at small and large wavenumbers, corresponding to the excitation of steady global modes and streamwise streaks, respectively. For transitional flows, Bonne *et al.* (2019) employed a Reynolds-averaged Navier–Stokes (RANS)-based resolvent approach on a Mach 1.6 shock-wave–transitional-boundary-layer interaction. This analysis revealed that low, medium and high-frequency modes are linked to the pseudoresonance process that implies amplification in the shear layer of the laminar separation bubble, the breathing motion of the bubble and turbulent fluctuations near the reattachment point, respectively.

In addition to laminar and transitional flows, stability analysis is also applicable to turbulent flows, particularly where the scale decoupling assumption between low-frequency large-scale motions and high-frequency small-scale turbulent structures holds (Rodi 1997; Iaccarino *et al.* 2003; Lawson & Barakos 2011). Toubert & Sandham (2009) conducted large-eddy simulations on a shock-impingement turbulent boundary layer flow at Mach 2.3. The linear stability analysis, based on time- and span-averaged LES data, identified a two-dimensional, stationary, globally unstable mode. This finding was later corroborated by Pirozzoli *et al.* (2010). Furthermore, under the scale decoupling assumption that URANS captures low- and medium-frequency SWBLI motions through time integration, while the turbulence model accounts for high-frequency fluctuations,

linear stability analysis based on linearised RANS equations has demonstrated success in transonic airfoil buffet (Crouch *et al.* 2009; Sartor, Mettrot & Sipp 2015*b*) and open cavity flows (Mettrot, Renac & Sipp 2014). Both studies exhibited clear spectral peaks that closely matched experimental results. Additionally, RANS-based resolvent analysis applied to globally stable transonic SWBLIs (Sartor *et al.* 2015*a*) and supersonic compression-corner flows (Hao 2023) indicated that low-frequency unsteadiness can be attributed to a preferred response of the separated flow to environmental noise.

Research on blunt-nose effects for double-cone configurations has predominantly focused on hypersonic laminar steady flow in existing literature (Holden 1997; Holden *et al.* 2006; Hao & Wen 2020). However, the impacts of nose bluntness on separation patterns, boundary-layer transition, and unsteady flow behaviour at transitional Reynolds numbers have been rarely reported. Therefore, the Mach 6 flow over three double-cone configurations with varying nose bluntness is investigated in the current study at a unit Reynolds number of $6.5 \times 10^6 \text{ m}^{-1}$. This Reynolds number is significantly higher than those considered in previous studies on blunt double-cone configurations (Holden 1997; Hao & Wen 2020), facilitating the occurrence of boundary-layer transition. A combination of schlieren photography, pressure measurements (using both sensors and pressure-sensitive paint (PSP)) and numerical simulations was employed to analyse the global flow characteristics and unsteady flow structures for the three double-cone configurations. Furthermore, based on the numerically reproduced base flow, GSA and resolvent analysis were performed to further explore the origins of the flow unsteadiness.

2. Methodology

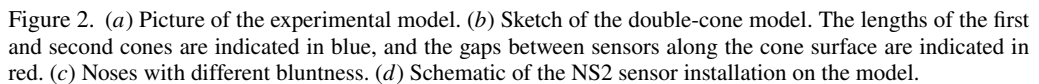
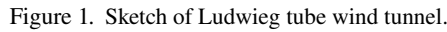
2.1. Experimental method

2.1.1. Wind tunnel

The hypersonic experiments on the double-cone configurations were conducted in a Mach 6 Ludwig tube wind tunnel at The Hong Kong Polytechnic University. Figure 1 provides a basic illustration of the wind tunnel, which primarily includes the storage tube, fast-action valve, Laval nozzle, test section and vacuum tank. The wind tunnel operated in a free-jet mode. The fast-action valve, located between the storage tube and the nozzle, controlled the wind tunnel's start and shutdown. Before testing, the storage tube was heated to prevent gas condensation downstream of the nozzle. When the fast-action valve was activated, under the pressure difference between the upstream storage tube and the downstream vacuum tank, the airflow accelerated to Mach 6 downstream of the circular Laval nozzle exit, which has a diameter of 300 mm. The hypersonic flow featured the total pressure of $700 \pm 5 \text{ kPa}$, the total temperature of $473 \pm 3 \text{ K}$, and the unit Reynolds number of $6.5 \times 10^6 \text{ m}^{-1}$, respectively. The wind tunnel operated in a noisy mode, with a noise level measured by a Pitot tube at approximately 2 %, which is typical for a conventional Ludwig tube (Munoz *et al.* 2019; Hoffman *et al.* 2023). The usable run time of this facility was around 30 ms, allowing for the measurement of global flow patterns and unsteady characteristics of hypersonic flows. Optical windows were installed on both the top and side of the test section to facilitate optical measurement techniques.

2.1.2. Double cone model

In this paper, three double cone configurations with varying nose bluntness were examined. Figure 2(*a*) shows one blunt double cone model in the wind tunnel. As illustrated in figure 2(*b*), the double-cone model consists of three components: a replaceable nose, the first cone and the second cone. The semiangles of the first and second



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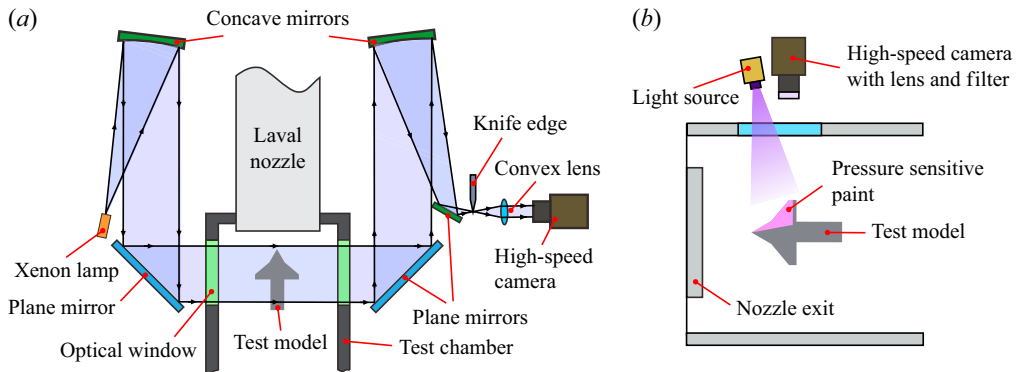


Figure 3. Schematics of (a) schlieren and (b) fast PSP techniques.

nose tip of the sharp case (shown in figure 2c), with the x - and r -axes representing the axial and radial directions, respectively, and φ indicating the azimuthal direction, following the right-hand rule (figures 2b and 2c).

2.1.3. Schlieren photography

The schlieren photography technique was employed to visualise the overall flow configuration and the unsteadiness of shock waves and shear layers. Figure 3(a) presents a schematic of the Z-type-like schlieren arrangement, which includes a Xenon lamp light source, a pair of concave mirrors with a diameter of 250 mm and a focal length of 2.5 m, a pair of plane mirrors with a diameter of 250 mm, a small plane mirror with a diameter of 50 mm, a knife edge and a Photron FASTCAM SA-Z high-speed camera. The light source was positioned at the focus of the left-hand concave mirror, directing emitted light to be reflected into a parallel beam by the mirror. This parallel light then passed through the test section after being reflected by the left-hand plane mirror and was subsequently cast onto the right-hand concave mirror after reflecting from another plane mirror. The light converged at the focus of the right-hand concave mirror, where a sharp knife edge was placed horizontally, allowing the schlieren technique to detect vertical density variations in the hypersonic flow fields. During the experiments, the high-speed camera operated at two sampling rates, 75 kfps and 300 kfps, with exposure times of 1.0 μs and 0.25 μs , respectively, to meet various experimental requirements. The lenses used for these sampling rates were the Nikon AF Micro-Nikkor 60 mm f/2.8D and the Nikon AF Nikkor 28 mm f/2.8, respectively, ensuring a sufficiently large observation area in different experiments. Under these two conditions, the spatial resolution of the schlieren images was 0.18 mm pixel⁻¹ and 0.38 mm pixel⁻¹, respectively. The maximum possible displacement of airflow during exposure, estimated based on the free stream velocity, corresponds to approximately 5 pixels and 0.6 pixels for the two imaging configurations, respectively.

2.1.4. Pressure measurements

In the pressure measurements, two techniques – pressure sensors and fast PSP – were employed to provide a quantitative understanding of the mean features and unsteady characteristics of hypersonic flows.

As shown in figure 2(b), 13 pressure sensors were arranged in two rows to measure wall surface pressure, with an azimuth angle interval of 45° between the two rows. This study utilised two arrangements for the pressure sensors. In the first arrangement, four XCQ-093

	Sensor 1	Sensor 2	Sensor 3	Sensor 4	Sensor 5	Sensor 6	Sensor 7
Row 1	0.452	0.589	0.726	1.094	1.246	1.398	—
Row 2	0.418	0.555	0.691	1.018*	1.170*	1.322*	1.474*

Table 1. Axial locations of the pressure sensors, normalised by the reference length L . The numbering of sensors follows the axial direction. Symbols marked with * denote the Kulite sensors installed using the first measurement arrangement. All other sensors are NS-2 sensors installed using the second measurement arrangement.

Kulite sensors (diameter 2.4 mm) were flush-mounted on the second cone. These sensors had a measurement range of 35 kPa and an accuracy of 0.1 % FS (± 35 Pa), with a natural frequency of 150 kHz. However, due to a protective B-screen over the sensors, the frequency response was limited to 30 kHz. The second arrangement was adopted because the Kulite sensors could not be directly installed on the first cone due to space constraints. To measure pressure on the first cone, three pressure tapings were machined just upstream of the Kulite sensors, and three steel tubes – each 25 mm long with outer and inner diameters of 1.0 and 0.5 mm, respectively, – were flush-mounted on the first cone. The other ends of these tubes were connected to NS-2 pressure sensors, which had a measurement range of 100 kPa and an accuracy of 0.1 % FS (± 100 Pa). A schematic of the NS2 sensor installation (the second arrangement method) on the model is depicted in [figure 2\(d\)](#). Additionally, another row of six sensors (Row 1) was arranged on the model using the second arrangement method, with their relative locations displayed in [figure 2\(b\)](#). [Table 1](#) lists the axial positions of the pressure sensors. All signals captured by the pressure sensors were collected using a National Instruments DAQ card (16 channels) with a sampling rate of 50 kS s^{-1} . The sampling rate of 50 kS s^{-1} was selected as it provides a Nyquist frequency (25 kHz) that sufficiently exceeds the frequency range of the low-frequency unsteadiness under investigation (typically below 5 kHz). This rate is also compatible with the effective bandwidth of the instrumented sensors. Before each test, the pressure sensors were recalibrated to eliminate any drift errors. It is important to note that the response time of the pressure measurements using the second arrangement was relatively long due to the presence of the steel tubes, which prevented the NS-2 sensors from detecting high-frequency pressure fluctuations. Therefore, the NS-2 sensors were used only to obtain time-averaged pressure, while dynamic pressure analysis relied on the signals obtained by the Kulite sensors.

Since the sensors can only measure pressure at specific points, the fast PSP technique – a global pressure measurement method widely used in unsteady three-dimensional supersonic and hypersonic experimental investigations (Burton & Babinsky [2012](#); Xiang & Babinsky [2019](#); Mason, Natarajan & Kumar [2021](#); Liu *et al.* [2022b](#); Jenquin, Johnson & Narayanaswamy [2023](#); Liu *et al.* [2024](#)) – was employed to obtain pressure distribution on the double cone surfaces. The preparation of the current fast PSP was based on formulations by Peng *et al.* ([2018](#)), utilising a platinum complex (Pt(II) meso-tetrakis (pentafluorophenyl) porphine) as the luminophore and mesoporous silica particles as the host. The use of these mesoporous particles effectively reduced the response time of the PSP, as their hollow structures extended the diffusion path of oxygen. The absorption and emission peaks of the fast PSP were around 390 nm and 650 nm, respectively. The frequency response of the fast PSP used in the current study was calibrated prior to the wind tunnel experiments. This calibration was performed using an acoustic resonance tube, following the established method proposed by Pandey & Gregory ([2016](#)). This set-up generates sinusoidal pressure oscillations of known amplitude and frequency, allowing for

direct measurement of the PSP luminescence response. The precalibration results show the amplitude attenuation was approximately 3.3 dB at 2 kHz and 5.6 dB at 4 kHz. In this study, the double cone surface was sprayed with PSP over an azimuthal range of approximately 120°. Before applying the PSP, a white base coating was sprayed to enhance the signal-to-noise ratio. Figure 3(b) illustrates the fast PSP measurement system, which includes a light source and a Photron FASTCAM SA-Z high-speed camera equipped with a Nikon AF Micro-Nikkor 60 mm f/2.8D lens and a 650 ± 25 nm bandpass filter. Visualisation was conducted at a sampling rate of 10 kfps with an exposure time of 98 μ s, resulting in a spatial resolution of 0.27 mm pixel⁻¹ for the captured images. To obtain data for *in situ* calibration as described by Liu *et al.* (2022b), the first row of sensors recorded pressure synchronously during the PSP measurements. Notably, during the schlieren measurement, pressure data collected from both the first and second rows of sensors were captured simultaneously. The experiments using schlieren and PSP techniques in different runs showed excellent repeatability, as confirmed by the pressure data obtained from the sensors of Row 1. Therefore, the pressure data collected from sensors of Row 2 were also utilised for *in situ* calibration. Given the short test duration, the high thermal conductivity of the aluminium model, and the use of an *in situ* calibration (Liu *et al.* 2022a), thermal effects on the PSP pressure measurements are considered negligible. Additionally, every instantaneous pressure image was spatially filtered using a 3 $\times$ 3 pixels square mean window to improve the signal-to-noise ratio.

2.2. Numerical method

To enhance the understanding of the complex phenomenon of hypersonic double-cone flow, this study systematically conducts numerical simulations, thereby supplementing experimental data with comprehensive flow field information and quantitative analysis. Notably, the current study identifies boundary-layer transition in all cases, which has a significant impact on the size of the separation bubble (Sinha *et al.* 1999; Wright *et al.* 2000; Holden 2015). Early numerical simulations of a Mach 8, 25°–50° double-cone flow by Wright *et al.* (2000) used the RANS method with a low Reynolds k – ϵ turbulence model. These simulations indicated that accounting for turbulence effects can reduce the size of the separation bubble, aligning results more closely with experimental data. However, accurately predicting the transition location remains a considerable challenge. In the recent numerical studies, Pezzella, de Rosa & Donelli (2015) and Knight *et al.* (2017) employed a novel method that fixed the transition at a specific streamwise location to rebuild the Mach 7.3 hypersonic double-wedge flow, ensuring it aligned with the experimental results. In their approach, the flow field upstream of the fixed transition point was simulated using the RANS method with turbulence production terms disabled, effectively modelling laminar flow, while the downstream region incorporated full turbulence modelling. In the present study, guided by the scale decoupling assumption under which RANS resolves the global flow structures of primary interest, the same strategy of prescribing a fixed transition location x_t is adopted to numerically reproduce the main shock and separation structures observed in the double-cone experiments. For the computational domain downstream of this transition, turbulence effects are accounted for using the RANS method with the Spalart–Allmaras model (Spalart & Allmaras 1992), incorporating modifications by Edwards & Chandra (1996). The fixed transition locations in different cases are estimated based on experimental results; further details can be found in Appendix A.

In the numerical simulation, the flow is governed by the RANS equation,

$$\frac{\partial \mathbf{U}}{\partial t} = \mathcal{N}(\mathbf{U}), \quad (2.1)$$

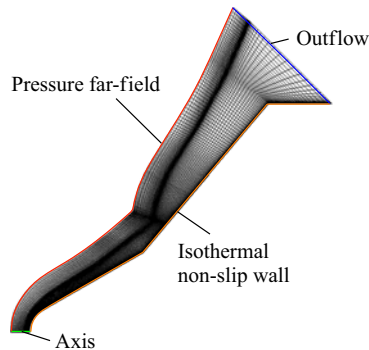


Figure 4. Computational grid for case RN8 (each 10th point is shown).

where $\mathbf{U}$ represents the vector of conservative variables, and $\mathcal{N}$ is the RANS operator. Note that the eddy viscosity is accounted for using the Spalart–Allmaras turbulence model in the region downstream of the transition point, while the source terms for turbulence are set to zero in the upstream region.

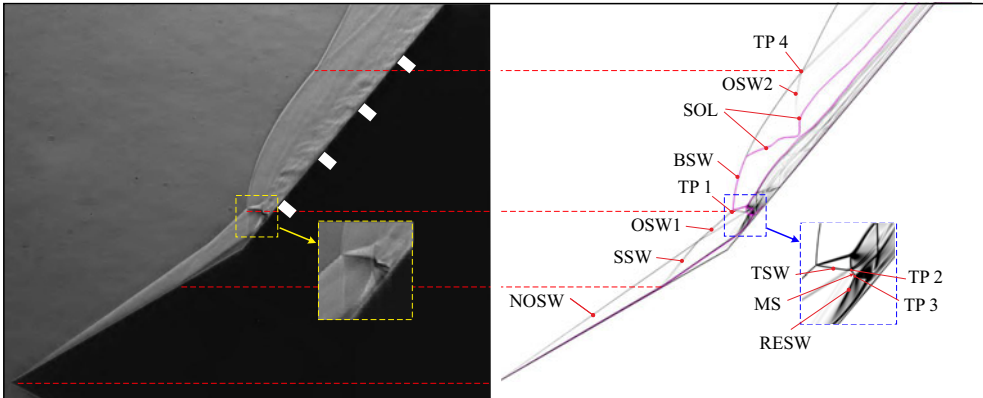
The two-dimensional axisymmetric numerical simulations are performed using an in-house multiblock parallel finite-volume computational fluid dynamics code – PHAROS, which has been effectively applied to hypersonic flows over compression ramp, double-wedge and double-cone configurations (Hao & Wen 2020; Hao *et al.* 2022; Hao 2023; Hao *et al.* 2023). Air is modelled as an ideal gas with a specific heat ratio γ of 1.4. The Prandtl number and turbulent Prandtl number are set at 0.72 and 0.90, respectively, while molecular viscosity is determined using Sutherland’s law. In the simulations, inviscid fluxes are calculated using the modified Steger–Warming scheme (MacCormack 2014), and viscous fluxes are computed with a second-order central difference scheme. An implicit line relaxation method (Wright, Candler & Bose 1998) is employed for pseudo time stepping. The computational grid for case RN8 is illustrated in figure 4, featuring 2200 nodes in the axial direction and 490 nodes in the radial direction. Isothermal non-slip wall boundary condition is applied to the cone surface. The wall temperature is set at 300 K. The grid is refined in regions near the wall, shock waves, and separation bubble. The height of the first grid layer in the wall-normal direction is set to 1.0×10^{-6} m. The grid-convergence validation is performed prior to the simulations; further details can be found in Appendix B.

3. Experimental results

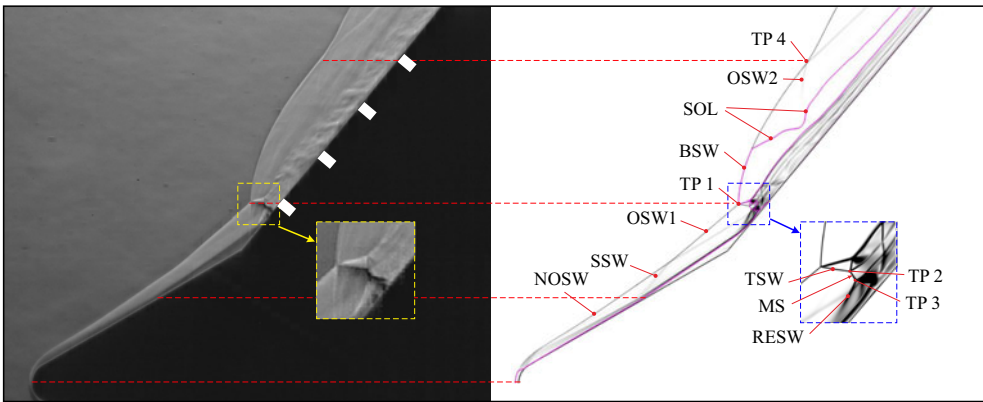
3.1. Global flow pattern of double-cone flows

The schlieren technique effectively visualises shock wave structures and flow separation according to the density changes in hypersonic flow fields. Figure 5 displays schlieren images captured experimentally at a sampling rate of 75 kfps, along with corresponding numerical schlieren results for the three scenarios. In the numerical simulations, the transition locations were fixed at $x_t/L = 0.86$, 0.85 and 0.59 for cases RN0, RN4 and RN8, respectively (details on the estimation of x_t can be found in Appendix A), yielding the closest alignment with experimental data. The schlieren results depicted in figure 5 distinctly reveal the presence of a separation bubble at the cone junction, alongside the complex shock wave system generated by the flow interaction. Both features are well reproduced by the numerical simulation, closely matching the experimental phenomena.

(a)



(b)



(c)

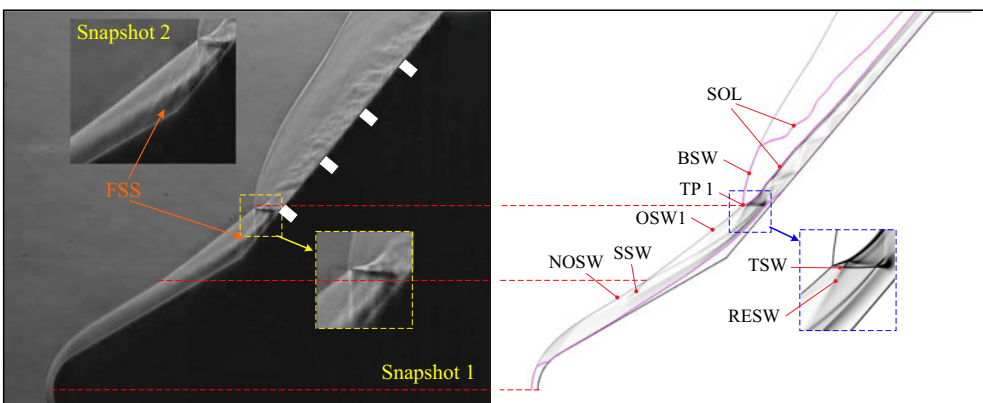


Figure 5. The experimental and numerical schlieren images of cases (a) RN0, (b) RN4 and (c) RN8, with flows structures annotated, including finescale structures (FSS), nose-induced oblique shock wave (NOSW), separation shock waves (SSW), transmitted oblique shock wave on the first cone (OSW1), bow shock wave (BSW), transmitted shock wave (TSW), reattachment shock wave (RESW), Mach stem (MS), sonic line (SOL), triple point (TP) and a relatively weak oblique shock wave on the second cone (OSW2). The white markers indicate the positions of the four Kulite sensors.

The schlieren results indicate that the separation point gradually moves upstream as the nose bluntness increases. The separation bubbles in cases RN0 and RN4 are very similar, exhibiting typical laminar boundary-layer separation morphology on the first cone. When the nose radius increases to 8 mm, the thick entropy layer downstream of the nose remains unswallowed by the boundary layer within a short streamwise distance, and its influence persists. The interaction between the shock wave and the boundary layer, along with the entropy layer, leads to a significant morphological alteration of the separation bubble compared with cases RN0 and RN4. Notably, case RN8 exhibits the emergence of finescale, unsteady vortical structures on the first cone, as shown by the two snapshots in [figure 5\(c\)](#). Nevertheless, a common feature across all three cases is a prominent Type V SSI. In each case, the separation shock wave first interacts with the nose-induced shock wave in a Type VI interaction, generating the transmitted oblique shock wave OSW1. The OSW1 subsequently intersects the detached bow shock ahead of the second cone, forming a distinct TSW and a TP (TP1) characteristic of a Type V pattern. In the flow field above TP1, the flow is subsonic immediately downstream of the bow shock, with the corresponding region enclosed by the sonic line (the purple solid line) indicated in the numerical schlieren images. Conversely, the airflow below TP1 remains supersonic as it passes through the TSW. A detailed inspection reveals two types of shock wave systems below TP1 in the three cases. For cases RN0 and RN4, the separation bubble is relatively small, and the compression waves rapidly coalesce into an oblique shock wave during the reattachment process. This shock wave intersects with TSW, resulting in a Mach shock wave reflection typified by a discernible Mach stem and downstream slip lines, observable in the close-up images of both experimental and numerical schlieren results in [figure 5](#). In contrast, the extended separation region in case RN8 inhibits full convergence of compression waves, yielding a weaker reattachment shock wave that engages in regular shock wave reflection with the TSW. Consequently, due to the two different wave structures below TP1, the downstream supersonic jet exhibits two patterns: RN0 and RN4 develop a three-layer jet stratified by slip lines emanating from TPs TP2 and TP3, whereas RN8 forms a single-layer structure. It is worth noting that in the experimental schlieren images, both the three-layer and single-layer structures of the supersonic jet maintain only a short streamwise distance before large-scale vortex structures emerge and disrupt the stratified flow. However, these large-scale vortical structures within the supersonic jet are not successfully reproduced in steady RANS simulations due to their inherent limitations in resolving unsteady vortical structures.

In fact, the detached bow shock wave undergoes progressive attenuation along the streamwise direction, transitioning into an oblique shock wave while maintaining supersonic flow conditions downstream. In cases RN0 and RN4, schlieren images indicate that the supersonic airflow traversing the middle region of the second cone decelerates again through a relatively weak oblique shock wave (designated OSW 2 in [figure 5](#)), aligning the airflow gradually parallel to the cone surface. This weak shock wave intersects with the outer oblique shock wave, resulting in a Type VI SSI, along with another TP (TP4) and a discernible slip line. In contrast, although the angle of the outer oblique shock wave in case RN8 shows an increasing trend near the tail region of the second cone, no typical Type VI SSI is observed. It can be inferred that this absence may be attributed to complex wave interactions wherein the outer oblique shock wave interacts with distributed weak compression waves. These compression waves, while physically present, possess insufficient density gradients for clear visualisation in schlieren images.

Regarding the separation bubble, a comprehensive analysis is carried out based on the numerical results. [Figure 6\(a,c,e\)](#) present the enlarged Mach number contours delineating

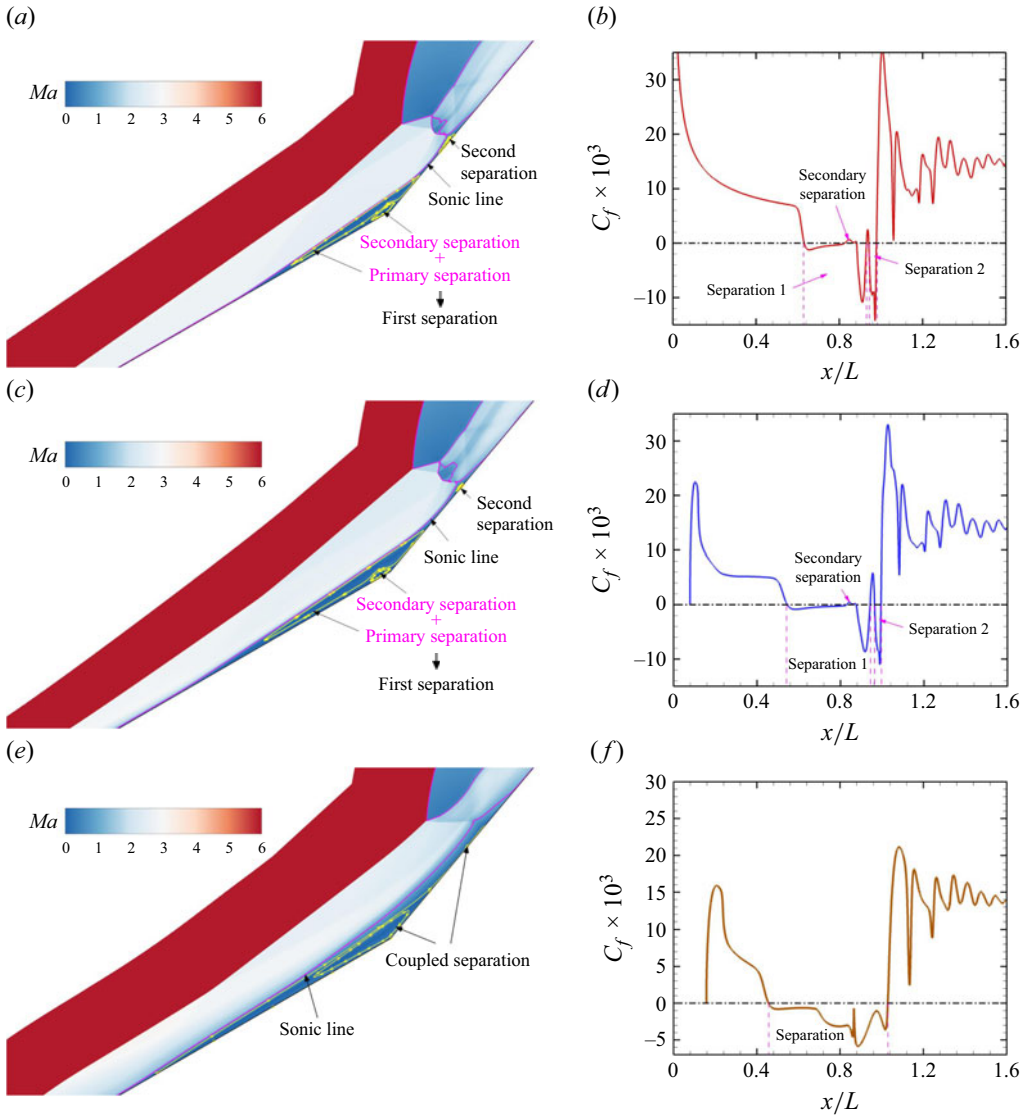


Figure 6. Enlarged Mach number contours near the cone junction and the friction coefficient distributions for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

flow structures near the separation zones. The streamline within the separation bubble is depicted through the yellow arrowed lines. Figure 6(b,d,f) display the friction coefficient (C_f) distributions; herein, $C_f = \tau_w / 0.5 \rho_0 u_0^2$, where τ_w , ρ_0 , u_0 denote the wall shear stress, free stream density and free stream velocity, respectively. The results indicate that for cases RN0 and RN4, the reflected shock wave originating from TP3 impinges on the boundary layer downstream of the reattachment of the separation bubble at the cone junction, leading to another separation bubble. The streamlines within the separation bubbles demonstrate the spatial segregation of these two recirculation regions, as further evidenced by the C_f distributions. It is worth noting that a secondary separation occurs within the first separation bubble, as shown in both the Mach number contours and C_f distributions. Conversely, for case RN8, the adverse pressure gradient caused by the geometric deflection

	Separation 1		Separation 2	
	x_{S1}/L	x_{R1}/L	x_{S2}/L	x_{R2}/L
Case RN0	0.631	0.930	0.941	0.981
Case RN4	0.544	0.941	0.963	0.993
Case RN8	0.455	1.030	—	—

Table 2. Axial locations of the separation and reattachment points. The symbols *S* and *R* in the subscripts denote the separation and reattachment points, respectively. The numbers 1 and 2 correspond to the first and second separation bubbles, respectively.

at the cone junction combines with that induced by the impingement of TSW, resulting in a relatively large-scale coupled separation bubble without any secondary separation occurring beneath it. The axial locations of the separation and reattachment points for all three cases are listed in [table 2](#). Results regarding the secondary separation beneath the first separation bubble are not included.

The PSP-measured pressure contours within a 120° azimuthal region are shown in [figure 7\(a,c,e\)](#). These results demonstrate a clearly axisymmetric flow, in contrast to the significant three-dimensional structures reported in the numerical work of Hao *et al.* (2022) when the flow is fully laminar. Notably, the observed minor curvature of high-pressure regions along the azimuthal direction is attributed to the optical image distortion, and the actual cone junction is marked by the red dashed lines. The streamwise pressure coefficient distributions ($C_p = p_w/0.5\rho_0u_0^2$, where p_w is the wall pressure), derived from numerical simulations, pressure sensors and PSP results along the central meridian line, are presented in [figure 7\(b,d,f\)](#), indicating good agreements between the experimental and numerical results. It is noteworthy that the PSP data along the model centreline were prioritised for optical distortion correction, yielding a streamwise displacement $\Delta x < 0.4$ mm. This ensures a negligible impact on the measured pressure peaks, as shown in [figure 7\(b\)](#). It is clear that the peak pressures in cases RN0 and RN4 are significantly higher than that in case RN8. Combining the flow structures shown in [figures 6](#) and [7](#), the pressure rises before reaching the peak value in cases RN0 and RN4, and can be divided into three stages. The first is the pressure rise induced by the airflow deflection at the separation point of the first bubble; the second pressure-rise stage occurs during the reattachment process of the first separation bubble; and the third stage is induced by the impingement of TSW on the boundary layer on the second cone. The latter two stages exhibit a steeper pressure rise compared with the first. In case RN8, the coupling of the two separation regions causes the latter two pressure rises to merge into a single gradual increase, primarily driven by a series of compression waves. Furthermore, the shock wave emitted from the TP TP2 in cases RN0 and RN4, along with the reattachment shock wave in case RN8, impinge on the shear layer downstream of TP TP1. These interactions result in strong reflected expansion waves, which subsequently impinge on the cone surface, causing a rapid pressure decrease. Additionally, these expansion waves and the shock waves within the supersonic jet undergo repeated reflections between the shear layer and the cone surface, leading to fluctuations in pressure distributions on the second cone.

[Figure 8](#) presents the comparison of C_p and C_f distributions for the three cases. While all cases exhibit a characteristic pressure plateau downstream of the separation point, distinct plateau magnitudes are observed. As the nose bluntness increases, the pressure plateau magnitude p_p gradually decreases. As shown in the C_p distributions, an over-expansion effect occurs downstream of the blunt noses in cases RN4 and RN8, leading to lower pressures compared with the sharp-nose case; correspondingly, the friction

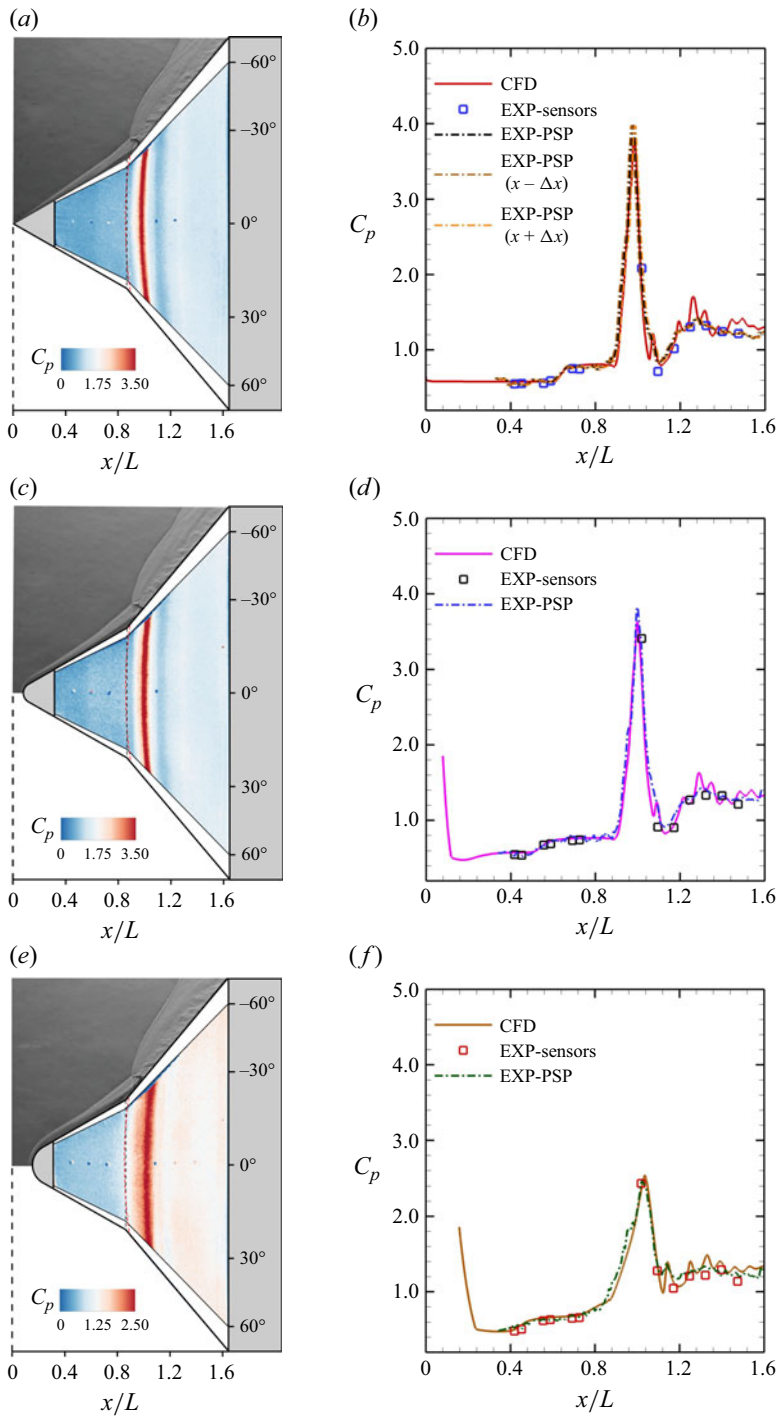


Figure 7. Pressure contours and pressure distributions along the meridian lines for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

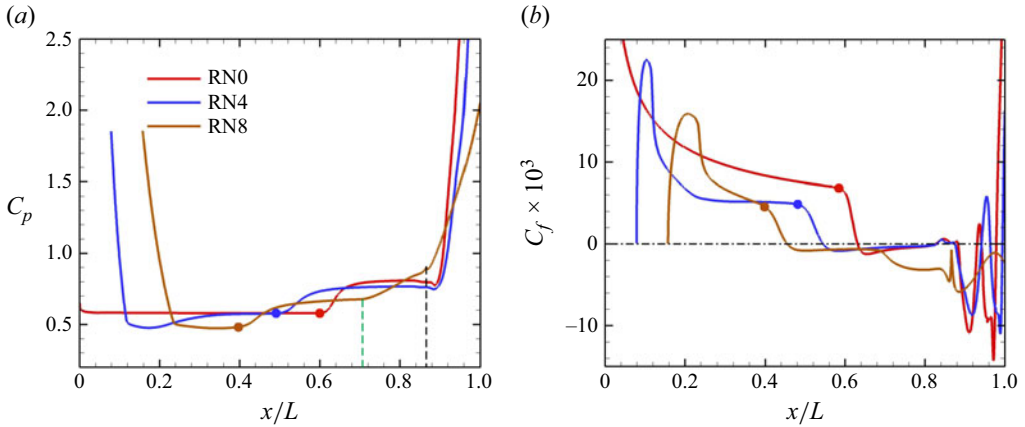


Figure 8. The comparison of (a) pressure coefficient C_p and (b) friction coefficient C_f distributions. The black dashed line indicates the location of the cone junction. The dots denote the upstream conditions prior to the interaction region.

coefficients in these blunt cases are also reduced relative to case RN0. Comparing cases RN0 and RN4 reveals that although the wall pressure upstream of separation recovers to the RN0 level in case RN4, the friction coefficient remains significantly lower. According to the free interaction theory (Chapman, Kuehn & Larson 1957), p_p is related to $p_0 C_{f,0}^{1/2}$ (where p_0 and $C_{f,0}$ represent the wall pressure and friction coefficient upstream of the separation region). The lower $C_{f,0}$ in case RN4 directly explains its reduced p_p relative to case RN0. In case RN8, insufficient streamwise distance prevents the wall pressure and friction coefficient from fully recovering to RN0 levels upstream of the separation point. These lower upstream values consequently contribute to a further reduction in the p_p compared with both cases RN0 and RN4. Moreover, the plateau terminations in cases RN0 and RN4 coincide with the cone junction (marked by the black dashed line), and the pressure rise downstream of this point is caused by the reattachment process. However, a notable deviation occurs in case RN8: both numerical simulation and PSP data reveal a gradual pressure rise in the region downstream of the plateau (at approximately $x/L = 0.7$, marked by the green dashed line) and before the cone junction. This subtle pressure increase precedes the sharper rise associated with reattachment. Considering that free interaction theory predicts higher plateau pressures for shock interactions involving transitional or turbulent boundary layers compared with purely laminar interactions (Babinsky & Harvey 2011), one possible explanation for this slow pressure rise in case RN8 is the onset of transition to turbulence within the shear layer overlying the separation bubble on the first cone. This developing turbulence would enhance momentum transport, gradually increasing the pressure upstream of the reattachment.

3.2. Dynamic analysis of double-cone flows

In the previous section, the global flow patterns for the three configurations were introduced. Notably, examination of the schlieren sequences revealed the presence of shock-wave oscillations across all cases (refer to the Supplementary materials for the movies are available at <https://doi.org/10.1017/jfm.2026.11443>). However, the intensity of these oscillations is significantly weaker than the large-scale pulsations reported in earlier experimental and numerical studies (Jagadeesh *et al.* 2003; Hornung *et al.* 2021). To delve deeper into the unsteady characteristics of these cases, we conduct a comprehensive

analysis that integrates data obtained from both schlieren visualisation and pressure measurements in this section.

Unlike the two-dimensional geometries, such as compression corners and double-wedge configurations, the axisymmetric double-cone models employed in the current study offer a distinct advantage when utilising schlieren techniques. The axisymmetric features effectively minimise spatial integration effects along the spanwise direction, thereby enhancing the fidelity of the captured flow dynamics. This set-up enables spectral analysis via the SPOD method. This technique has been widely applied to investigate unsteadiness in supersonic and hypersonic flows, including applications to second-mode instabilities (Butler & Laurence 2021), shear-layer structures (Benitez *et al.* 2023, 2024) and shock wave-separation bubble dynamics (Hoffman *et al.* 2022; Li *et al.* 2024; Posladak *et al.* 2024). The SPOD method decomposes the flow field into a hierarchy of orthogonal modes, each oscillating at distinct frequencies, which collectively capture the coherent spatiotemporal evolution of dominant flow structures. Motivated by the identification of relatively high-frequency structures (approaching 100 kHz) in the supersonic jet within a Mach 7 double-cone flow, as reported in the previous experimental study by Posladak *et al.* (2024), this study carried out the spectral analysis based on the schlieren sequences at sampling rates of both 75 kfps and 300 kfps. The schlieren visualisation at the higher sampling rate of 300 kfps primarily focuses on the flow field around the second cone, targeting the potential capture of these high-frequency phenomena.

Regarding the SPOD analysis applied to schlieren sequences acquired at 75 kfps, the Welch method was employed with a Hamming window of length 512 points to mitigate spectral leakage. The schlieren sequences were segmented into 14 overlapping blocks using a 75 % overlap ratio. This set-up yielded a frequency resolution of 146 Hz. Normalising this resolution using the characteristic length L and free stream velocity u_0 results in a Strouhal number resolution of $St_L = 0.008$. Consistent with typical practice for identifying dominant flow features, the analysis focuses exclusively on the leading SPOD mode due to its highest energy content. Figure 9(a) present the modal energy distributions of the leading SPOD modes for all three cases. The results indicate that for cases RN0 and RN4, the modal energies decrease with increasing frequency, reflecting broadband low-frequency characteristics. Interestingly, when the nose radius increases to 8 mm, a significant peak emerges at around $St_L = 0.12$ ($f = 2.197$ kHz) in the leading mode energy, signifying the presence of dominant low-frequency flow unsteadiness in this case. Figure 9(b,c,d) show the selected SPOD mode shapes at $St_L = 0.12$ for all three cases. In cases RN0 and RN4, the broadband low-frequency unsteadiness is primarily associated with the oscillation of shock wave system, encompassing the separation shock wave, bow shock wave, transmission shock wave and the Mach shock wave reflection occurring downstream of the first separation bubble, as well as the shear layer of the separation bubble. In case RN8, although a dominant frequency in the low-frequency region is observed, the SPOD mode shape indicates that the unsteady structures corresponding to this frequency are remarkably similar to those observed in cases RN0 and RN4, predominantly characterised by wave structures and the shear layer of the separation bubble.

As for the unsteadiness observed on the second cone, figure 10 presents the SPOD results derived from the schlieren sequences captured at the sampling rate of 300 kfps. The energy distributions of SPOD leading mode depicted in figure 10(a) reveal the presence of broadband low-frequency characteristics in cases RN0 and RN4, alongside the notable peak frequency at $St_L = 0.113$ ($f = 2.05$ kHz) in case RN8. This is attributed to the inclusion of the bow shock wave, transmission shock wave and portions of the separated shock wave within the field of view. In the higher-frequency region, the analysis

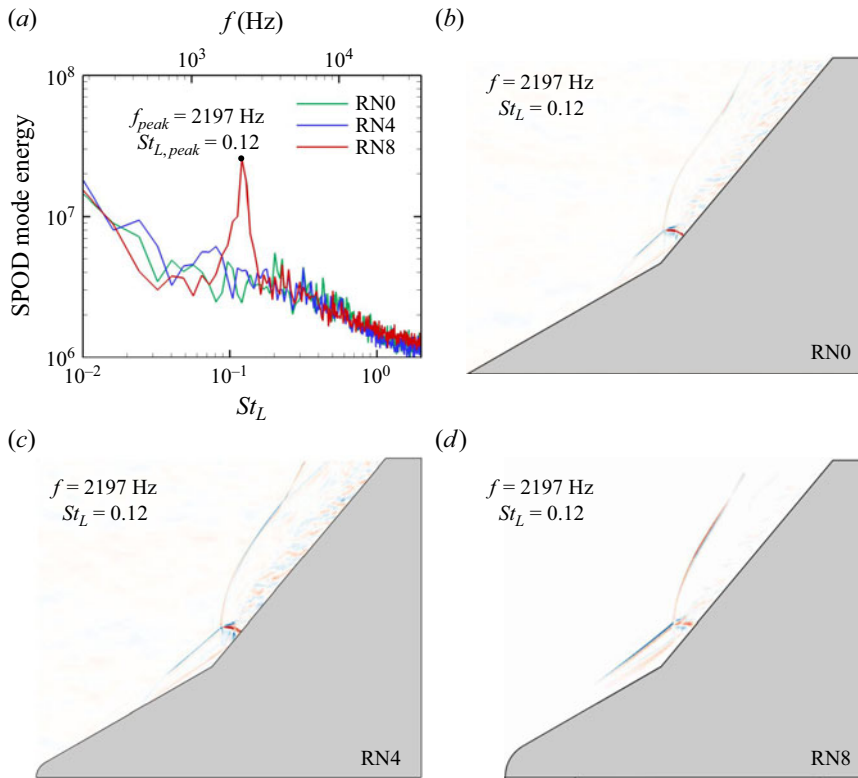


Figure 9. (a) The SPOD leading mode energy distributions for all cases. Selected SPOD mode shapes at $f = 2197$ Hz for cases (b) RN0, (c) RN4 and (d) RN8, derived from the schlieren sequences at the sampling rate of 75 kfps.

indicates that the leading modal energy begins to rise at approximately $St_L = 2$ for cases RN0 and RN4, and at $St_L = 1$ for case RN8. In cases RN0 and RN4, the peak frequencies of the leading modal energy are very close, approximately falling within the range around $St_L = 3.9\text{--}4.4$ ($f = 70\text{--}80$ kHz). In comparison, in case RN8, the peak of the leading mode energy in the high-frequency range occurs at slightly lower frequency around $St_L = 3.3\text{--}3.9$ ($f = 60\text{--}70$ kHz). Figure 10(b,c,d) illustrate the selected modal shapes corresponding to $St_L = 4.1$ ($f = 75$ kHz) for cases RN0 and RN4, and $St_L = 3.6$ ($f = 65$ kHz) for case RN8. The SPOD mode shapes indicate that, in these cases, high-frequency unsteadiness is primarily associated with the flow structures along the outer boundary of the supersonic jet. Based on their location within the high-shear region and their coherent vortical appearance in schlieren snapshots (figures 5 and 10), these high-frequency structures are suggestive of an instability intrinsic to the shear layer. Their characteristics are plausibly linked to a Kelvin–Helmholtz type instability (Knisely 2016; Tumuklu *et al.* 2018a,b).

Consequently, the SPOD analysis of flow structures captured through schlieren visualisation reveals that the unsteady phenomena in the three cases manifest in two primary forms. The first form consists of low-frequency unsteady oscillations associated with the shock wave system and the separation bubble. The second form involves large-scale vortices along the outer boundary of supersonic jet corresponding to the high-frequency unsteadiness, which plausibly arises from the Kelvin–Helmholtz instability of the strong shear layer on the second cone. Notably, unlike the broadband low-frequency

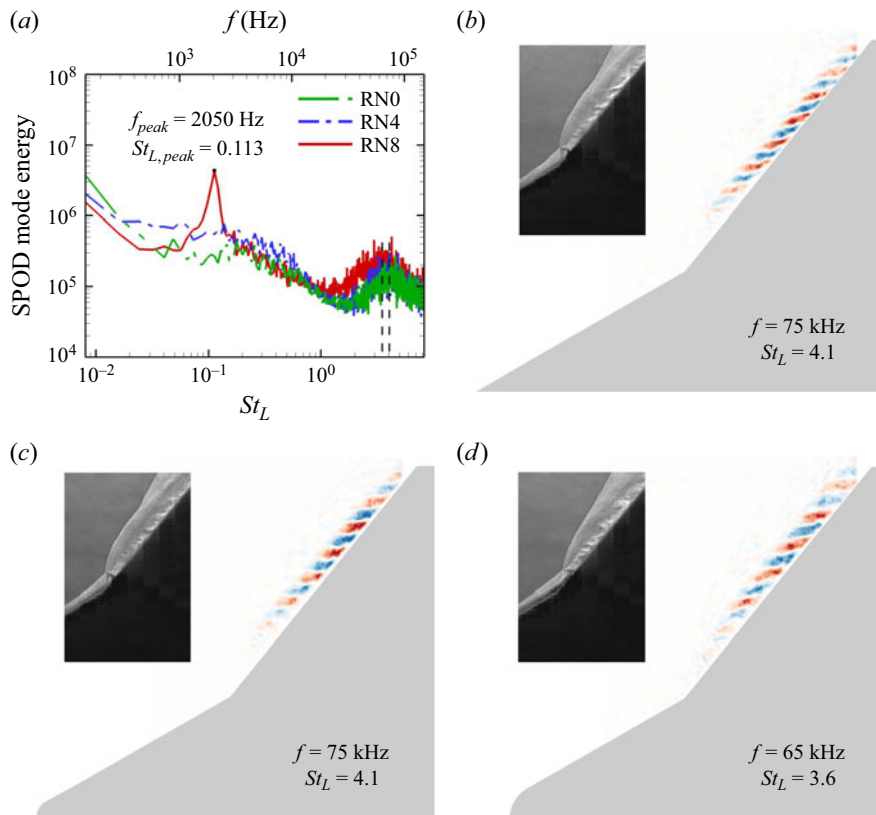


Figure 10. (a) The SPOD leading mode energy distributions for all cases. Selected SPOD mode shapes for cases (b) RN0, (c) RN4 and (d) RN8, derived from the schlieren sequences at the sampling rate of 300 kfps.

oscillations observed in cases RN0 and RN4, case RN8 exhibits a distinct dominant frequency within the low-frequency range – a phenomenon not previously documented. The underlying mechanism for this dominant frequency will be further explored in § 4.

Notably, the spectral analysis of the schlieren images primarily highlights the spatial unsteady structures in the meridian plane. In this study, dynamic wall pressure measurements were also conducted using pressure sensors and fast PSP technique. Actually, the wall surface parameters are intrinsically linked to the spatial flow field structures, prompting us to perform a dynamic analysis of the pressure fluctuations on the wall surface. As mentioned in § 2.1.4, the frequency response of pressure measurements from the NS2 sensors is influenced by the steel tube, leading us to focus solely on the pressure signals obtained from the four Kulite sensors and fast PSP technique. In the current study, the Kulite sensors and PSP system were sampled at 50 kHz and 10 kHz, respectively. Consequently, the effective upper frequency for spectral analysis was governed by the corresponding Nyquist frequencies (25 kHz and 5 kHz), restricting the dynamic analysis of wall pressure to low-frequency unsteady phenomena in this section. Figure 11(a,c,e) present the normalised pressure fluctuation signals (relative to local mean pressure) recorded by the four Kulite sensors on the second cone. In all three cases, the first sensor is positioned near the peak pressure location, while the remaining three sensors are situated beneath the supersonic jet, as illustrated in figure 5 (the sensor positions are indicated by white markers). It is evident that the normalised fluctuation amplitude at

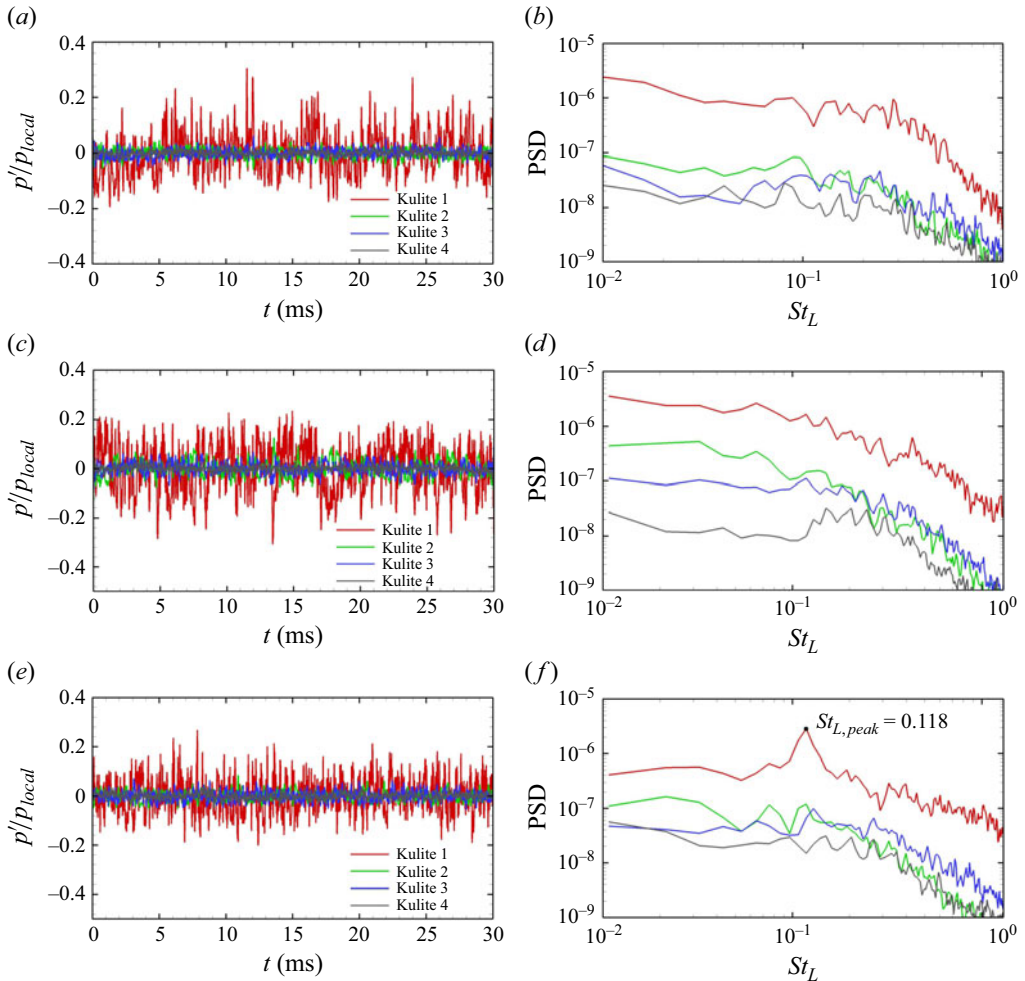


Figure 11. Pressure signals captured by the Kulite sensors and PSD distributions of the pressure signals for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

the first sensor is significantly higher than those at the other three locations. To quantify these observations, power spectral density (PSD) analysis of p'/p_{local} was conducted using the Welch method, as shown in figure 11(b,d,f). The results reveal that, in the low-frequency range, the energy content at the first sensor is approximately one to two orders of magnitude greater than those of the other three sensors. This finding indicates that the supersonic jet on the second cone exhibits very weak low-frequency unsteadiness, which aligns well with the SPOD results presented in figure 9. Additionally, the PSD distributions for the first sensor in case RN8 reveal a distinct peak at around $St_L = 0.118$ ($f = 2148$ Hz), whereas cases RN0 and RN4 display broadband low-frequency characteristics.

It is well established that the spectral analysis of schlieren sequences captures unsteady oscillations in the meridian plane. However, previous studies on shock-induced separations have demonstrated that low-frequency unsteadiness can also manifest in spanwise movements within nominally two-dimensional flows (Sidharth *et al.* 2018; Cao *et al.* 2021b; Hao *et al.* 2021, 2022). Although the time-averaged pressure contours shown in figure 7 do not reveal obvious large-scale three-dimensional flow structures, this

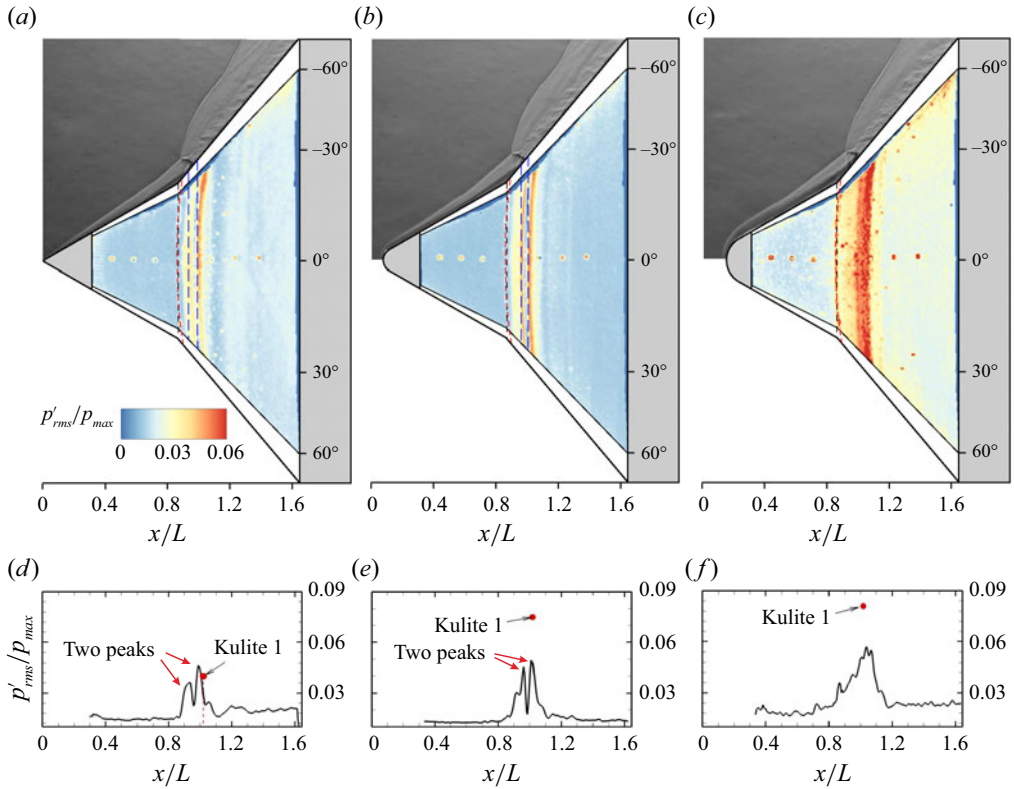


Figure 12. Root mean square (r.m.s.) contours of wall-pressure fluctuations obtained from the PSP and the corresponding distribution along the meridian lines for cases (a,d) RN0, (b,e) RN4 and (c,f) RN8. The red dots denote the r.m.s. values measured by the first Kulite sensor.

analysis aims to further specifically investigate the potential presence of low-frequency unsteady oscillations in the azimuthal direction. Figure 12 presents the r.m.s. of the wall pressure fluctuations (p'_{rms}). Figure 12(a – c) show contours of p'_{rms} normalised by the maximum of mean pressure (p_{max}), obtained from the PSP data. Figure 12(d – f) plot the corresponding p'_{rms} distribution along the central meridian line. For comparison, the r.m.s. values measured by the first Kulite sensor are marked by red symbols. The comparison shows that the PSP derived p'_{rms} levels are approximately 40 % lower than the sensor measurements. Actually, this difference aligns with the frequency-dependent attenuation inherent to the fast PSP technique. As noted in § 2.1.4, the PSP attenuation is approximately 3.3 dB at 2 kHz (amplitude reduction to approximately 70 %) and 5.6 dB at 4 kHz (amplitude reduction to approximately 50 %). It is important to note that the frequency-dependent attenuation is expected to be spatially uniform, a condition intrinsically linked to the homogeneity of the PSP coating itself (Funderburk & Narayanaswamy 2019). This expectation is supported by the absence of any non-physical artefacts in the time-averaged PSP pressure fields, indicating a high degree of spatial uniformity in the applied paint layer. Given this consistent PSP application and optical set-up, the data are employed in this study primarily for the qualitative observation of spatial mode shapes, effectively complementing the pointwise measurements provided by conventional sensors. A quantitative comparison of unsteady characteristics across different cases, which would require accounting for the precise frequency-dependent attenuation in amplitude, is beyond

the scope of the present work. In [figure 12](#), it is of interest that distinct dual peaks are observed within the p'_{rms} contours on the second cone for cases RN0 and RN4. In contrast, for case RN8, the first peak diminishes significantly, becoming barely discernible. For case RN0, the two high p'_{rms} regions are located at approximately $x/L = 0.90\text{--}0.95$ and $x/L = 0.98\text{--}1.02$, while for case RN4 they occur at approximately $x/L = 0.93\text{--}0.97$ and $x/L = 1.00\text{--}1.04$. As illustrated in the Mach number contours and C_f distributions in [figure 6](#), both cases RN0 and RN4 display two individual separation bubbles. According to the axial locations of the separation and reattachment points listed in [table 2](#), it can be observed that the first high p'_{rms} region corresponds to the area between points R1 and S2, while the second high p'_{rms} region corresponds to the region downstream of point R2 (here, S and R denote the separation and reattachment points, respectively; the numbers 1 and 2 correspond to the first and second separation bubbles, respectively). The spatial correlation indicates that the dual-peak structure in the pressure fluctuations corresponds to the two spatially distinct separation regions. Critically, in the two cases, the second, much smaller separation bubble is directly induced by the impinging shock wave emanating from the TP TP3 (as visible in [figures 5a](#) and [5b](#)). Consequently, the oscillation of this bubble serves as a direct manifestation for the unsteadiness of the impinging shock, which is the primary cause of the distinct second p'_{rms} peak on the cone surface. In contrast, case RN8 features a large, coupled separation bubble. The thickened boundary layer at the shock impingement point substantially attenuates the shock's direct impact on the cone surface. Here, the shock's influence is indirectly transmitted by driving the low-frequency, large-scale oscillation of the entire bubble. This integrated dynamic results in a single, broad p'_{rms} peak, thereby explaining the absence of the dual-peak configuration in the p'_{rms} distribution observed in the other cases.

Spectral analysis for the surface pressure field was performed using the SPOD method. The datasets were segmented into 14 blocks with 75 % overlap, yielding a frequency resolution of 156 Hz. As illustrated in [figure 13\(a\)](#), the leading SPOD modes for cases RN0 and RN4 exhibit broadband low-frequency characteristics while case RN8 displays a pronounced spectral peak at $St_L = 0.112$ ($f = 2031$ Hz). The selected SPOD mode shapes at $St_L = 0.112$ are also shown in [figure 13\(b,c,d\)](#). Analysis of these modes reveals two key observations: no significant large-scale low-frequency three-dimensional unsteady structures are evident in any case, and the observed low-frequency unsteadiness primarily manifests as axial-direction oscillations, i.e. the axisymmetric oscillation mode. Moreover, the spatial distribution of modal energy demonstrates distinct patterns: in cases RN0 and RN4, the most energetic components are predominantly localised near the high p'_{rms} regions identified previously, with other regions exhibiting less distinct modal energy. The corresponding SPOD mode on the first cone exhibits a diffuse and less organised spatial structure, manifesting as a weak and indistinct alternating pattern. This stems from the relatively small separation bubble, which is characterised by inherently broadband, low-amplitude unsteadiness. In comparison, case RN8 is dominated by a strong, narrowband unsteadiness associated with a large separation bubble. This bubble effectively couples and amplifies the unsteady signal, transmitting vigorous dynamics from the downstream impinging shock upstream. The result is a highly coherent and spatially organised mode at a dominant frequency, with primary energy concentrated on the second cone and a visible pattern linked to bubble breathing on the first cone. This mode is clearly visualised as a strong red/blue alternation. Differences are also evident in the downstream region beneath the supersonic jet. In cases RN0 and RN4, the jet has a complex three-layer structure (as shown in [figures 5a](#) and [5b](#)). Reflections of upstream unsteady shock waves within this irregular environment are highly disturbed. This complexity, combined with the inherent broadband oscillations in this region, yields a diffuse SPOD mode with indistinct spatial

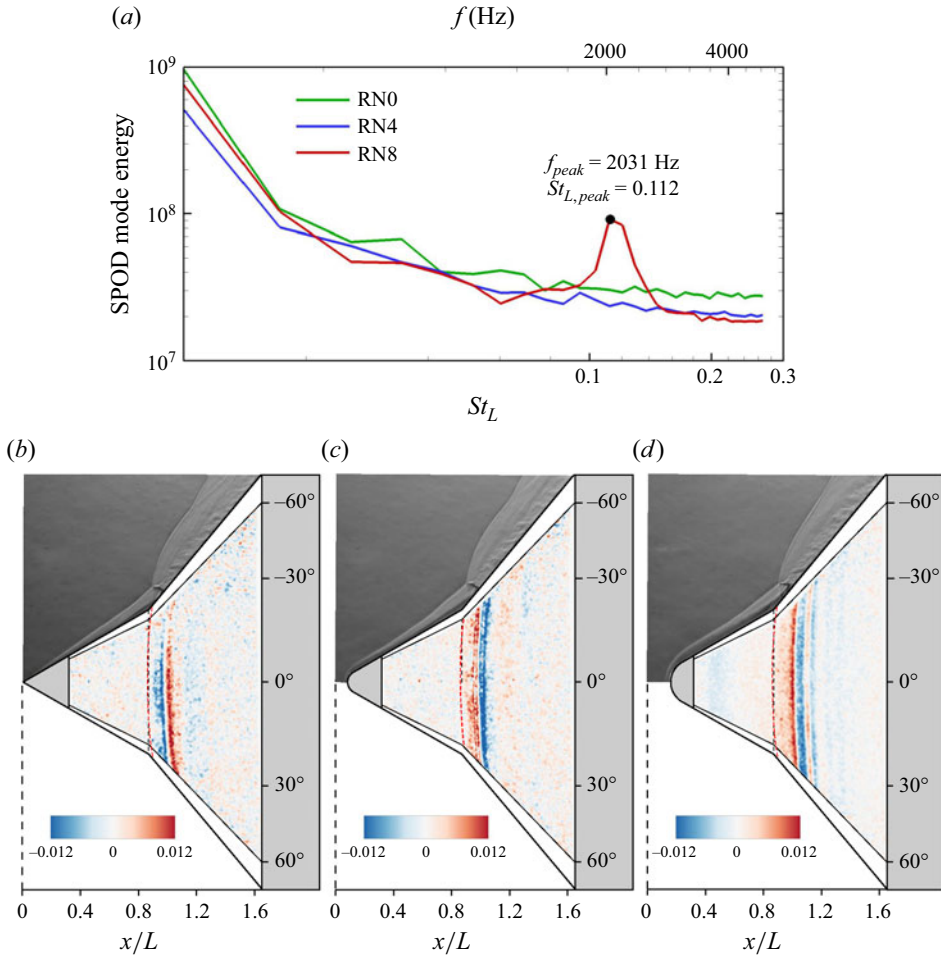


Figure 13. (a) The SPOD leading-mode energy distributions. Selected SPOD mode shapes at $f = 2031$ Hz for cases (b) RN0, (c) RN4 and (d) RN8, derived from the pressure field data obtained using PSP technique.

organisation. Conversely, in case RN8, the jet possesses a simpler single-layer structure (figure 5c), which supports clean and regular shock reflections that form a distinct wave pattern. Consequently, the wall-pressure mode shapes in figure 13(d), particularly the clear alternating red/blue bands in case RN8, serve as a direct spatial imprint of the regular reflection of upstream unsteady shock waves on the second cone surface.

4. Global stability analysis and resolvent analysis

The preceding section demonstrated that the axisymmetric base flows for all three cases were successfully reproduced by assuming a fixed transition location in the numerical simulation. Downstream of this transition point, the flow was simulated using the RANS method, incorporating turbulence effects through the Spalart–Allmaras model. Building upon the scale-decoupling assumption, this section employs both GSA and resolvent analysis on the simulated base flows to elucidate the physical mechanism driving the experimentally observed flow unsteadiness. Critically, this investigation focuses

exclusively on the axisymmetric oscillation mode, as demonstrated by the experimental observations in the spectral analysis of the wall-pressure field.

4.1. Global stability analysis

In the GSA method, the flow field is assumed can be decomposed into two components: the steady base flow and small-amplitude unsteady perturbations. This decomposition is formally expressed as

$$\mathbf{U} = \bar{\mathbf{U}} + \mathbf{U}', \quad (4.1)$$

where $\bar{\mathbf{U}}$ and $\mathbf{U}'$ denote the steady base flow and unsteady perturbations, respectively. Since the steady base flow satisfies the governing (2.1), substituting the decomposition (4.1) into (2.1) and neglecting higher-order terms yields the linearised perturbation equation,

$$\frac{\partial \mathbf{U}'}{\partial t} = \mathcal{J} \mathbf{U}', \quad (4.2)$$

where $\mathcal{J}$ is the Jacobian matrix, which corresponds to the linearisation of the RANS operator and can be evaluated based on the steady base flow. The perturbation $\mathbf{U}'$ is assumed to take the form

$$\mathbf{U}' = \hat{\mathbf{U}} e^{\lambda t}, \quad (4.3)$$

here, $\hat{\mathbf{U}}$ and λ represent the eigenfunction and eigenvalue, respectively. By substituting (4.3) into (4.2), we derive an eigenvalue problem that can be expressed in the following general form:

$$\lambda \hat{\mathbf{U}} = \mathcal{J} \hat{\mathbf{U}}. \quad (4.4)$$

The eigenvalue problem is addressed using the implicitly restarted Arnoldi method, as implemented in ARPACK (Lehoucq, Sorensen & Yang 1998), and interfaced with the MUMPS library (Amestoy *et al.* 2011) for efficient linear algebra operations. The same discretisation method as in the flow solver is used except that the inviscid fluxes are calculated using a central scheme in smooth regions, while the modified Steger–Warming scheme is applied to compute inviscid fluxes at flow discontinuities detected by the modified Ducros sensor (Hendrickson, Kartha & Candler 2018), thereby minimising numerical dissipation. Boundary conditions are consistent with those implemented in the numerical simulations. In the turbulent regions, the linearised source term for turbulence in the Spalart–Allmaras model is incorporated as described by Crouch, Garbaruk & Magidov (2007), without resorting to simplifications such as the frozen eddy-viscosity approach (Carini *et al.* 2017). Furthermore, since the computational grids from the numerical simulations are directly applied to GSA, an examination of grid convergence for GSA is also presented in [Appendix B](#).

The present GSA solver has been effectively applied to both laminar and turbulent flow regimes (Hao *et al.* 2021, 2022; Hao 2023). For each case investigated in this study, we seek to compute 50 eigenvalues using the GSA solver. The eigenvalues consist of a real part λ_r and an imaginary part λ_i , which correspond to the growth rate and angular frequency of the perturbations, respectively. A flow is considered globally stable if λ_r is less than zero for all computed eigenvalues. Furthermore, the eigenmodes can be classified based on the imaginary parts: those with $\lambda_i = 0$ are termed stationary modes, indicating no oscillation in time, while those with $\lambda_i \neq 0$ are categorised as oscillatory modes.

The eigenvalue spectra for all three cases are presented in [figure 14\(a,c,e\)](#). To facilitate a direct comparison with experimental measurements, frequencies are expressed

dimensionally (f) and non-dimensionally (St_L), rather than as angular frequency (ω). An analysis of the spectra reveals consistently negative real parts across all eigenvalues in the three cases, thereby confirming global stability of the flow systems. A closer examination of the spectra shows that the growth rate of the least stable mode progressively increases with increasing nose bluntness. Moreover, a notable change in the characteristics of the least stable mode is observed: for cases RN0 and RN4, the least stable mode remains stationary; however, for case RN8, the least stable mode exhibits an oscillatory nature. Interestingly, this unsteady dominant mode in case RN8 possesses a damping rate very close to zero, classifying it as a marginally stable mode. The corresponding frequency of this mode is approximately $St_L = 0.112$ ($f = 2030$ Hz), which aligns remarkably well with the dominant frequency identified in the experimental results.

The eigenfunctions corresponding to the least stable modes in each case are illustrated in figure 14(b,d,f). For cases RN0 and RN4, where stationary modes dominate, the perturbation ρ' is predominantly concentrated within the shock-dominated regions – specifically the separation shock wave, detached bow shock wave and Mach shock wave reflection. A significant presence of perturbation is concurrently observed in the shear layer enveloping the separation bubble, whereas the recirculation zone within the bubble exhibits minimal modal activity. Interestingly, although the least stable mode in case RN8 is oscillatory, the spatial distribution of the perturbation ρ' shows some similarity to the stationary modes in cases RN0 and RN4, with peak amplitudes likewise localised within the shock wave structures. This consistent spatial localisation of least stable modes aligns closely with the shock-associated coherent structures identified via SPOD in prior experimental investigations, confirming the physical relevance of the computed eigenmodes.

4.2. Resolvent analysis

The GSA results demonstrate that all three base flows behave as a noise amplifier. To investigate the response of these base flows to the perturbations, a small-amplitude forcing term is added to (4.2),

$$\frac{\partial \mathbf{U}'}{\partial t} = \mathcal{J}\mathbf{U}' + \mathcal{B}\mathbf{f}', \quad (4.5)$$

where the matrix $\mathcal{B}$ constrains the forcing to a localised region. We assume that the forcing and response are temporally harmonic,

$$\mathbf{f}' = \hat{\mathbf{f}}e^{i\omega t}, \quad \mathbf{U}' = \hat{\mathbf{U}}e^{i\omega t}, \quad (4.6)$$

where ω is the angular frequency. Substituting (4.6) into (4.5) yields

$$\hat{\mathbf{U}} = \mathcal{R}\mathcal{B}\hat{\mathbf{f}}, \quad \mathcal{R} = (i\omega\mathcal{I} - \mathcal{J}), \quad (4.7)$$

where, $\mathcal{R}$ and $\mathcal{I}$ are the resolvent and identity operators, respectively. The resolvent analysis aims to identify the forcing and response pair that exhibits the maximum energy amplification, or the optimal gain σ , which is defined as

$$\sigma^2(\omega) = \max_{\hat{\mathbf{f}}} \frac{\|\hat{\mathbf{U}}\|_E}{\|\mathcal{B}\hat{\mathbf{f}}\|_E}. \quad (4.8)$$

The Euclidean norm and Chu energy norm are employed to calculate the energies of both the forcing and the response. This optimisation problem can be reformulated as an eigenvalue problem, as demonstrated by Bugeat *et al.* (2019). To solve this eigenvalue

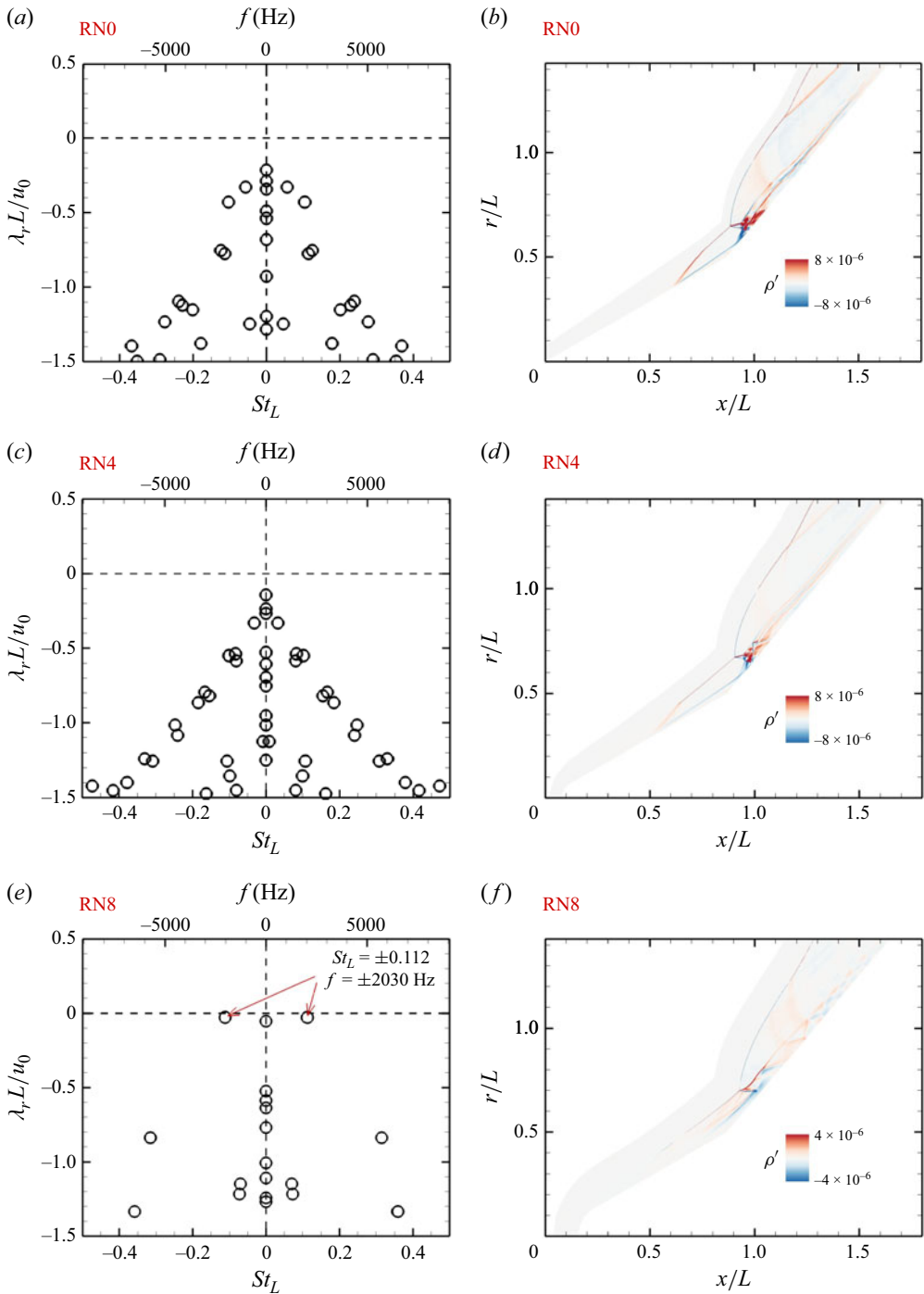


Figure 14. Eigenvalues spectra and real parts of the density perturbation of the least stable mode for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

problem, we utilise the power iteration method, which is well-suited for identifying the dominant eigenvalue associated with maximum energy amplification. The discretisation methods used in this process remain consistent with those implemented in GSA, ensuring coherence in the numerical framework.

The resolvent analysis is conducted based on the base flow over a wide frequency range to investigate its response to the environmental noise. Notably, guided by the scale decoupling assumption, the RANS simulations effectively capture low-frequency, large-scale structures – such as shock-wave motion, while high-frequency, small-scale vortical structures are modelled using turbulence models. Consequently, this preliminary resolvent analysis focuses exclusively on the low-frequency range in the current study. In the resolvent analysis, perturbations were applied throughout the entire computational domain using a methodology similar to that employed by Sartor *et al.* (2015*a,b*). Figure 15(*a,c,e*) illustrates the variation of the optimal gains with frequency. The results reveal that for cases RN0 and RN4, the maximum optimal gain occurs at the ultralow frequency, with a gradual decrease as frequency increases. This behaviour aligns with trends observed in the experimental results and has been reported in previous studies of laminar separation (Bugeat *et al.* 2022). Strikingly, case RN8 exhibits fundamentally different dynamics: after an initial decline, the optimal gain surges to a sharp peak at $St_L = 0.112$. To elucidate the low-frequency flow characteristics observed in RN0 and RN4, figures 15(*b*) and 15(*d*) present the optimal response modes at $St_L = 0.001$, corresponding to the red points marked in figures 15(*a*) and 15(*c*). The response modes in the two cases resemble the corresponding least stable modes predicted by GSA in figures 14(*b*) and 14(*d*). To quantitatively compare the modes in resolvent analysis and GSA, the projection of them is calculated, which is defined as

$$\xi = \frac{|\|\hat{\mathbf{U}}_{GSA}, \hat{\mathbf{U}}_{IO}\|_E|}{\|\hat{\mathbf{U}}_{GSA}\|_E \times \|\hat{\mathbf{U}}_{IO}\|_E}, \quad (4.9)$$

where $\hat{\mathbf{U}}_{GSA}$ and $\hat{\mathbf{U}}_{IO}$ denote the least stable GSA mode and optimal resolvent mode, respectively. The notation $\|\cdot\|_E$ represents the Chu norm (Chu 1965). Thus, the projection coefficient $\xi = 0$ signifies the two modes are orthogonal, while $\xi = 1$ indicates a perfect alignment between the two modes. The calculation shows that $\xi = 0.875$ and 0.915 for cases RN0 and RN4. This result reveals a strong alignment between the least stable GSA mode and the optimal resolvent mode for cases RN0 and RN4. It indicates that their broadband, low frequency response is a modal amplifier phenomenon. The globally stable flow preferentially amplifies the low frequency content of the environmental noise, with maximum gain occurring at the lowest frequencies where the system's dynamics are governed by its least damped stationary mode. This is consistent with established results for laminar separation (Bugeat *et al.* 2022). For the peak frequency identified in case RN8, the optimal response pattern in figure 15(*f*) exhibits nearly perfect agreement with the least stable global mode predicted by GSA in figure 14(*f*), with the projection coefficient $\xi = 0.991$. This confirms that the sharp spectral peak in RN8 originates from a true modal resonance. Here, the flow possesses a marginally stable oscillatory global mode that selectively and powerfully amplifies disturbances at its natural frequency (~ 2 kHz). This contrasts with the broadband response in RN0/RN4, which arises from the modal amplification of noise by a stationary global mode.

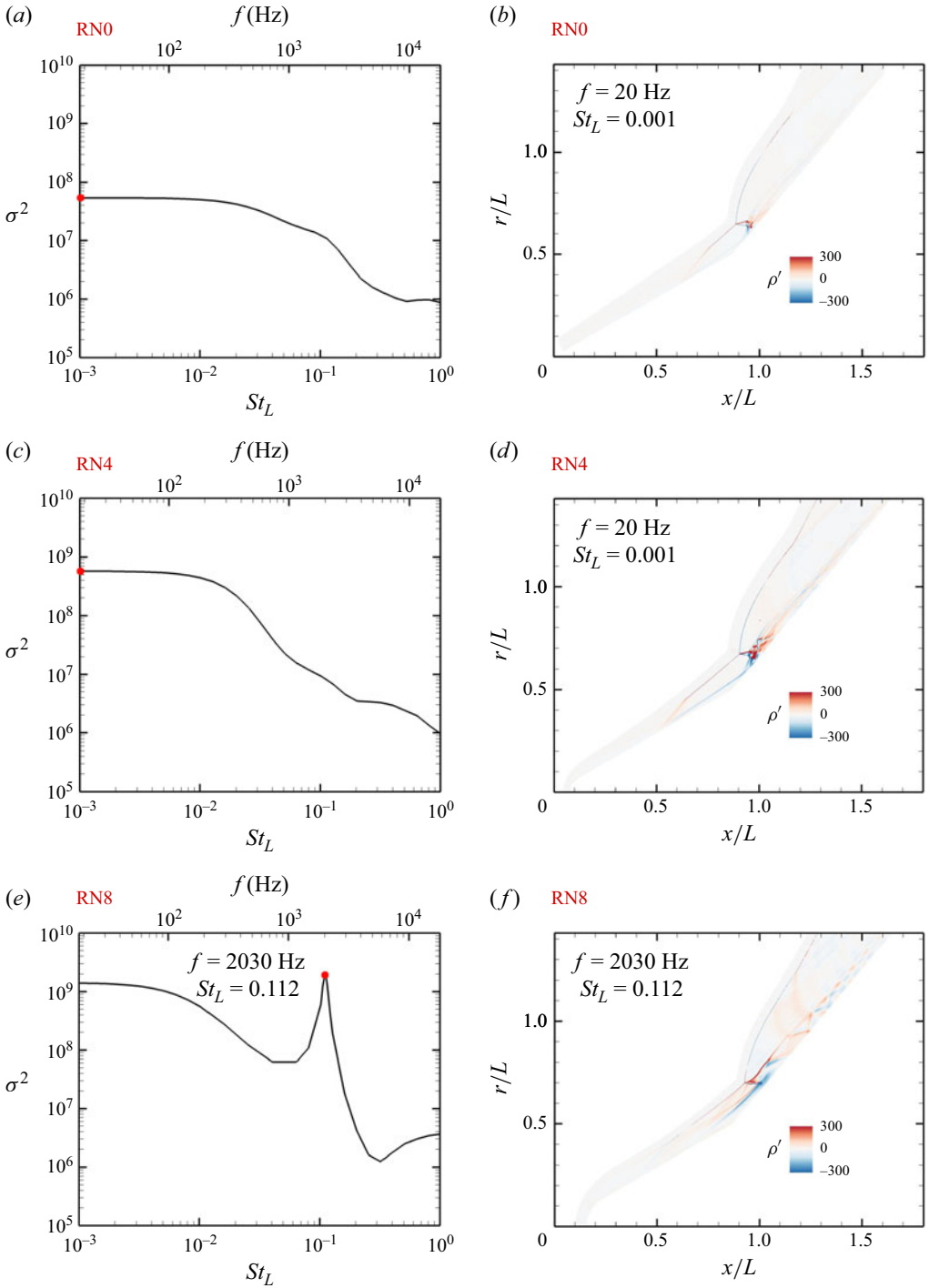


Figure 15. Optimal gain distributions and real parts of the density perturbations for the selected mode at specific frequency for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

5. Conclusion

This study investigates hypersonic double-cone flows at Mach 6 using combined experimental and numerical approaches. Three different blunt noses with radii of 0, 4, and 8 mm were employed to investigate their influence on the flow characteristics. Experimental investigations used high-speed schlieren photography and pressure measurements via sensors and PSP technique to capture both the general flow pattern and unsteady dynamics of the double-cone flows. Guided by the experimental observations, numerical simulations successfully captured the key features of the double-cone flows by implementing a fixed transition location. A detailed comparative analysis of the aerodynamic properties and flow physics for the three configurations is provided in the current study.

The results reveal that, under the 30° – 50° cone angle conditions in the double-cone flows, boundary layer separation consistently occurs at the cone junction across all three cases, accompanied by a detached bow shock wave located ahead of the second cone. Notably, the separation bubble size exhibits a monotonic increase with increasing nose bluntness. For the smaller bluntness cases (RN0 and RN4), the shear layer within the separation bubble on the first cone displays distinctly laminar characteristics. In contrast, case RN8 exhibits the development of vortical structures within the shear layer on the first cone, indicating the onset of shear-layer instability. Significant differences are also identified in the shock wave system beneath the TP. Cases RN0 and RN4 exhibit strong Mach shock wave reflections, resulting from the interaction between a robust reattachment shock wave and the TSW below the TP. Conversely, case RN8 demonstrates relatively weaker regular shock wave reflection. The combined experimental and numerical results reveal differences in separation topology: in cases RN0 and RN4, the cone-junction-induced separation between the two cones and the second separation triggered by the impingement of a strong oblique shock wave beneath the TP TP3 on the boundary layer are spatially isolated. In case RN8, however, these two separation regions coalesce, forming a single large-scale coupled separation bubble. The wall surface pressure distributions within the separation bubble regions also vary notably. Cases RN0 and RN4 exhibit pronounced pressure plateaus, with differences in plateau pressure levels explained by classical free-interaction theory. In case RN8, the pressure plateau within the separation bubble is significantly shorter, followed by a gradual pressure rise on the first cone – an observation closely linked to the onset of shear-layer transition process.

Experimental characterisation of unsteady dynamics through high-speed schlieren imaging identifies two primary types of unsteadiness across all cases. The first type involves low-frequency oscillations, characterised by the motions of the separation shock wave, detached bow shock wave and the shear layer enveloping the separation bubble. The second type of unsteadiness originates from the shear-layer instability on the second cone; here, intense shear between the supersonic jet and the mainstream flow generates large-scale vortical structures oscillating at very high frequencies, approximately 60–80 kHz. A novel observation unique to case RN8 is the experimental identification of a distinct low-frequency spectral peak at approximately 2 kHz. This dominant low-frequency motion has been further corroborated by unsteady pressure measurements. Moreover, SPOD analysis of the unsteady wall-pressure fields reveals that the low-frequency oscillations in all three cases exhibited a predominantly axisymmetric nature, indicating that the oscillations are primarily along the streamwise direction.

Building upon the scale decoupling assumption and flow fields reconstructed from numerical simulations, GSA and resolvent analysis are performed to elucidate the origins of the observed low-frequency unsteadiness. The GSA results conclusively demonstrate

that all three flow configurations are globally stable, thereby classifying them as noise amplifiers. The GSA eigenvalue spectra reveal that the least stable mode is stationary for cases RN0 and RN4 but oscillatory for case RN8. The resolvent analysis yields excellent agreement with experiments. Specifically, the optimal distributions display the peak within the ultralow-frequency range for cases RN0 and RN4, while case RN8 possesses a distinct peak at the experimentally observed dominant frequency (~ 2 kHz). The spatial structure of the optimal response corresponding to the peak frequency in case RN8 presents a striking resemblance to the least stable mode predicted by GSA. This compelling structural congruence strongly suggests that the dominant experimental oscillation arises from a modal resonance mechanism, wherein the marginally stable mode is selectively amplified by the flow's response to the environmental noise.

Overall, this work provides a comparative study of double-cone flows under varying nose bluntness conditions at Mach 6. The unconventional approach of using a fixed transition location in the simulations – guided by experimental data – successfully captured the overall flow structures and enabled a reasonable explanation of the low-frequency unsteadiness via stability analysis based on the scale decoupling assumption. Nevertheless, inherent limitations of this method persist: the simulations using the RANS method inherently lack the resolution needed to accurately depict the intricate transition process as well as the delicate vortical structures that develop on the second cone. To address these complexities, future research should integrate high-fidelity direct numerical simulations, which can offer the necessary physical fidelity to elucidate these intricate phenomena.

Supplementary movies. Supplementary movies are available at <https://doi.org/10.1017/jfm.2026.11443>.

Acknowledgements. The authors would like to express their gratitude to the editors and reviewers for their constructive feedback and valuable time, which have enhanced the quality of the paper.

Funding. This work was funded by the National Natural Science Foundation of China (no. 12472239) and the Hong Kong Research Grants Council (no. 15204322).

Declaration of interests. The authors report no conflict of interest.

Appendix A. Estimation for the onset of laminar–turbulent transition

The schlieren sequences reveal distinct characteristics among the three cases (figure 5). For RN0 and RN4, the shear layer on the first cone remains visually smooth, suggesting predominantly laminar flow. In contrast, case RN8 exhibits the emergence of finescale, unsteady vortical structures (figure 5c), which are indicative of growing shear-layer instabilities and may signal the beginning of the transition process. To obtain a quantitative, albeit indirect, estimate for the streamwise location of the transition process onset (x_t), we analysed the normalised standard deviation (SD_{nor}) of the schlieren intensity. The SD_{nor} fields are shown in figures 16(a), 16(c) and 16(e). The normalisation ensures that the black and white areas in the schlieren images correspond to values of 0 and 1, respectively. The SD_{nor} contours reveal that the shear layer on the first cone has relatively low SD_{nor} values for cases RN0 and RN4, while a notable increase is observed in case RN8 in this region. From these contours, the SD_{nor} distributions along the shear layer (indicated by the purple dashed–dotted lines) were extracted and plotted in figures 16(b), 16(d) and 16(f). It was found that SD_{nor} distributions increase at a specific location across all three cases which is taken as an initial reference for estimating x_t . In cases RN0 and RN4, these reference locations are located approximately at $x/L = 0.86$ and 0.86 , respectively, close to the cone junction ($x_c/L = 0.866$). For case RN8, the reference location on the

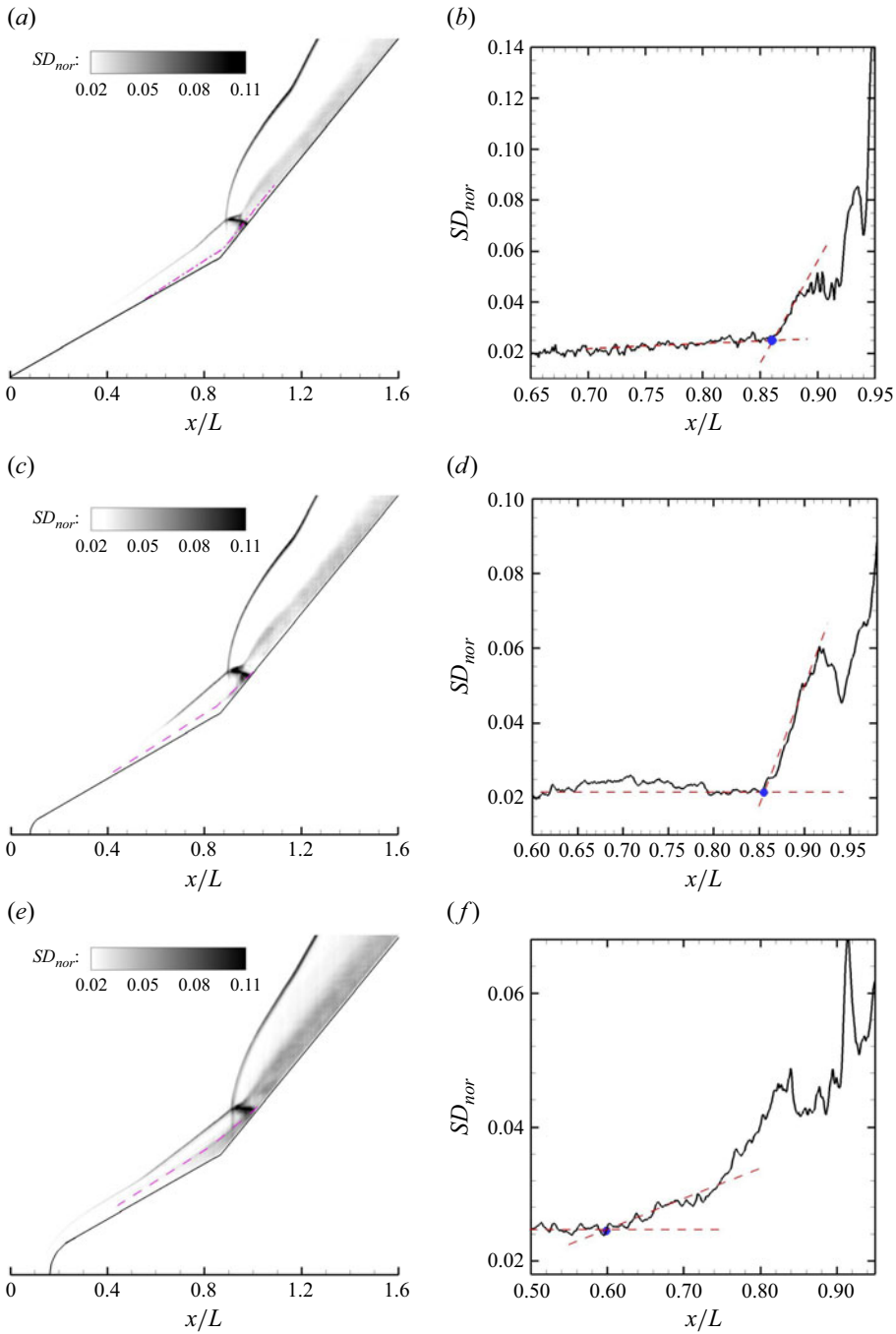


Figure 16. The standard deviation distributions derived from the schlieren sequences for cases (a,b) RN0, (c,d) RN4 and (e,f) RN8.

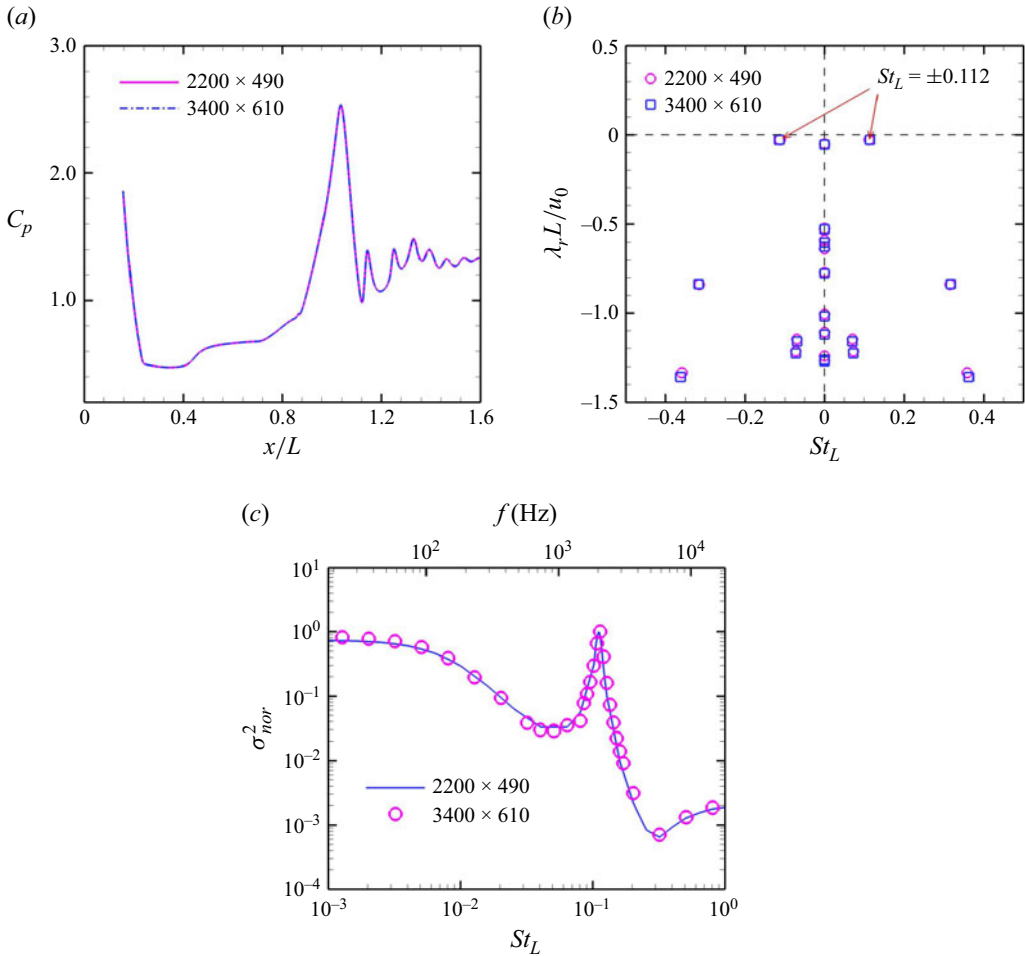


Figure 17. (a) pressure coefficient distributions, (b) eigenvalues spectra of GSA and (c) normalised optimal gain of resolvent analysis obtained using two computational grids.

first cone is approximately at $x/L = 0.60$. Guided by these experimental observations, the RANS simulations treat the flow upstream of x_t as effectively laminar (with turbulence production disabled), activating the full turbulence model downstream. Notably, the final values used in the simulations ($x_t/L = 0.86, 0.85, 0.59$ for RN0, RN4, RN8) were arrived at after iterative tuning around these initial estimates. The criterion was achieving the best overall agreement with the experimental mean flow features, as shown in [figure 5](#). Thus, the chosen x_t is not only consistent with the observed onset of shear-layer instability but is also validated by its ability to reproduce the key global flow structures measured in the experiments.

Appendix B. Validation of grid convergence

To assess grid convergence, we performed numerical simulations using two different grid resolutions – a coarse grid (2200×490) and a fine grid (3400×610). [Figure 17\(a\)](#) shows the C_p distributions for case RN8 obtained from both simulations. The close agreement between the two sets of results, evidenced by the nearly identical wall pressure

distributions, confirms good grid independence. Additionally, since the stability analysis uses the same computational grids, figures 17(b) and 17(c) compares the eigenvalue spectra and normalised optimal gain for case RN8 at both grid resolutions. The strong alignment of the results validates the adequacy of the coarse grid. Based on these thorough comparisons, we conclude that the coarse grid resolution employed in this study meets the necessary criteria for grid convergence.

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