

Influence of Nose Bluntness on Hypersonic Boundary-Layer Transition Onset in a Mach 8 Digital Wind Tunnel

Mateus Schuabb* and Lian Duan†
The Ohio State University, Columbus, OH 43210, USA

Anton Scholten‡
National Institute of Aerospace, Hampton, VA 23666, USA

Pedro Paredes§ and Meelan M. Choudhari¶
NASA Langley Research Center, Hampton, VA 23681, USA

Direct numerical simulations (DNS) are used to investigate laminar-turbulent transition over a 7-degree half-angle blunt cone with 5.2 mm nose radius at Mach 8, extending previous studies on the effects of free-stream acoustic disturbances on hypersonic boundary-layer receptivity and transition. Similar to the earlier work, the tunnel-like freestream acoustic disturbances are modeled from the acoustic radiation from the nozzle-wall turbulent boundary layer from a precursor DNS of the full-scale axisymmetric nozzle of the Sandia HWT-8. The onset of transition is clearly identified through surface quantities, including wall-shear stress, pressure, and heat flux. The transition region is also characterized by a rapid increase in spanwise-averaged, root-mean-square value of fluctuations in these quantities, accompanied by a broadening of both frequency and azimuthal wavenumber spectra. The onset of transition is considerably downstream in comparison with a nominally sharp cone at the same freestream conditions, highlighting the effect of nose bluntness. The peak N-factor for Mack's second mode disturbances at the transition onset location is very small, $N_T = 1.2$, similar to previous observations based on blunt cone experiments at hypersonic conditions. The paper also explores the physical mechanisms underlying the onset of transition despite such low modal amplification.

Nomenclature

C_f	=	skin friction coefficient, $C_f = \tau_w / (0.5\rho_\infty U_\infty^2)$, dimensionless
C_h	=	mean Stanton number, $C_h = q_w / (\rho_\infty U_\infty C_p (T_r - T_w))$, dimensionless
C_{p_w}	=	pressure coefficient, $C_{p_w} = (p - p_\infty) / (0.5\rho_\infty U_\infty^2)$, dimensionless
f	=	frequency, Hz
M	=	Mach number, dimensionless
m	=	azimuthal wavenumber, rad^{-1}
N	=	Logarithmic amplification factor
Pr	=	Prandtl number, $Pr = 0.71$, dimensionless
R	=	ideal gas constant, $R = 287$, J/(K·kg)
R_n	=	nose radius, m
Re_∞	=	freestream unit Reynolds number, $Re_\infty = \frac{\rho_\infty U_\infty}{\mu_\infty}$, m^{-1}
Re_{R_n}	=	nose Reynolds number, $Re_{R_n} = Re_\infty R_n$, dimensionless
T	=	temperature, K
T_r	=	recovery temperature, $T_r = T_\infty (1 + \sqrt{Pr} * \frac{\gamma-1}{2} M_\infty^2)$, K
U_∞	=	freestream velocity, m/s

*Graduate Student, Department of Mechanical and Aerospace Engineering

†Professor, Department of Mechanical and Aerospace Engineering, Associate Fellow, AIAA

‡Research Engineer, Computational AeroSciences Branch, NASA LaRC, Member, AIAA

§Aerospace Technologist, Computational AeroSciences Branch, M.S. 128, Associate Fellow, AIAA

¶Aerospace Technologist, Computational AeroSciences Branch, M.S. 128, Fellow, AIAA

a	=	speed of sound, m/s
p	=	pressure, Pa
r	=	radial coordinate
u	=	streamwise velocity, m/s
u_τ	=	friction velocity, $u_\tau \equiv \tau_w / \bar{\rho}_w$, m/s
v	=	spanwise velocity, m/s
w	=	wall-normal velocity, m/s
x, y, z	=	Cartesian coordinates, m
z_τ	=	viscous length, $z_\tau = \nu_w / u_\tau$, m
γ	=	specific heat ratio, $\gamma = C_p / C_v$, dimensionless
κ	=	thermal conductivity, $\kappa = \mu C_p / Pr$, W/(m·K)
μ	=	dynamic viscosity, kg/(m·s)
ν	=	kinematic viscosity, $\nu = \mu / \rho$, m ² ·s
ρ	=	density, kg/m ³
τ_w	=	wall shear stress, Pa

Subscripts

e	=	boundary-layer-edge variables
g	=	global
i	=	inflow station for the domain of direct numerical simulations
rms	=	root mean square
w	=	wall variables
$\infty, 1$	=	preshock freestream variables
$\infty, 2$	=	postshock freestream variables

Superscripts

$+$	=	inner wall units
$(\cdot)$	=	statistically averaged variables
$(\cdot)'$	=	perturbation from averaged variable

I. Introduction

Reducing uncertainty in boundary-layer transition (BLT) prediction remains a significant challenge for hypersonic vehicles. This is particularly important due to the considerable increase in surface heating due to transition. Classical stability theory has successfully explained transition onset on sharp axisymmetric cones via the exponential amplification of Mack’s second-mode instabilities [1]. However, as the nose radius increases, the resulting entropy layer can suppress these modal instabilities which initially results in a delay in transition, a trend supported by both experiments and transient growth [2, 3]. Further increase in nose bluntness leads to deviations between experimental observations and predictions from modal stability theory: the transition location stops moving downstream and instead advances upstream, referred to as transition reversal [2, 4–8]. All the datasets display two distinct regions referred to as “small bluntness”, where the transition location moves downstream with increased bluntness, and “large bluntness”, where the transition location rapidly moves upstream. Stetson [2] identified an intermediate regime for “moderately-large bluntness” values, wherein transition still moves downstream as the nose bluntness is increased, but at a significantly slower rate. In these cases, the frustum transition is measured at locations where the Mack’s second-mode (MM) waves are not sufficiently amplified, and the transition Reynolds number is lower than that for a sharp cone at the same flow conditions. In this regime, which approximately corresponds to $40,000 < Re_{R_n} < 500,000$, the MM N -factor decreases with increasing Re_{R_n} . Near $Re_{R_n} \approx 500,000$, the N -factor at the measured transition location is still nonzero but substantially smaller than the transition N -factor over a sharp cone. The physical mechanisms responsible for the low values of the MM N -factor at the transition location are unknown. This phenomenon suggests that modal growth alone cannot account for transition under all conditions. As a result, alternative frameworks involving nonmodal amplification mechanisms and forcing from freestream disturbances have been proposed as the cause of transition when strong bluntness effects are present. Recent work has shown that traveling disturbances may experience significant growth within the entropy layer [9–11] and trigger nonlinear processes that ultimately lead to breakdown [12]. As shown in the latter study, although the linearly amplified nonmodal disturbances are predominantly located outside the boundary layer, nonlinear interactions between oblique waves can generate stationary streaks that penetrate into and grow within the boundary layer. The subsequent breakdown of these streaks can then initiate transition. These insights indicate that the transition

process on blunt bodies is complex, influenced not only by baseflow characteristics but also by external forcing and transient growth pathways, particularly in conventional noisy hypersonic wind tunnels where most transition data are obtained.

Complementing experimental and theoretical efforts, high-fidelity direct numerical simulations (DNS) have been employed to analyze the early stages of instability growth, particularly within the entropy layer on the nose region of blunt geometries. So far, investigations of boundary-layer receptivity in hypersonic flows have primarily relied on either stochastic inflow forcing over a broad range of frequencies and wavenumbers [13–16], or the prescription of idealized freestream disturbances modeled as planar acoustic waves [17, 18]. For instance, Balakumar and Chou [18] performed two-dimensional simulations of boundary-layer transition over a 7° half-angle straight cone with varying bluntness at Mach 10, using experimentally measured freestream disturbance spectra but restricting the inflow forcing to be planar acoustic waves. While this approach yielded reasonably accurate predictions of transition onset for sharp cones, it proved less effective for blunt configurations since the applied disturbances failed to replicate the spatial complexity of the three-dimensional acoustic field typically encountered in wind tunnel environments. In a related study, Balakumar et al. [19] analyzed the receptivity of a Mach 3.5 boundary layer over a sharp cone to obliquely incident acoustic waves described by ad hoc analytical expressions. However, because nonmodal amplification mechanisms involve the constructive interference of multiple eigenmodes, the applicability of these simplified disturbance representations to transition prediction for the flow over a blunt body remains limited and warrants further investigation.

To investigate the role of broadband freestream disturbances and their nonmodal amplification within the entropy layer on boundary-layer transition, the present study extends the work of Liu et al. [20] and Schuabb et al. [21], who combined DNS and nonmodal linear stability analysis to examine the receptivity of a Mach 8 boundary layer to acoustic disturbances radiated from the nozzle walls of a conventional wind tunnel. In their approach, a “tunnel-like” acoustic field, constructed based on the acoustic noise emanated from turbulent nozzle-wall boundary layers, was introduced upstream of the bow shock over a 7° half-angle cone with nose radius $R_n = 5.2$ mm. Their simulations revealed significant disturbance amplification in the entropy layer, closely aligned with the predictions of nonmodal growth theory [11], as evidenced by spectral signatures in wall pressure and heat flux, and schlieren-like contours that qualitatively matched experimental observations. Building on those earlier simulations, the current study investigates a longer cone geometry to capture the onset of the laminar-turbulent transition. The objective is to establish a clearer connection between the nonmodal amplification mechanisms and the onset of transition in the absence of significant modal growth.

The structure of this paper is as follows. Section II details the freestream flow conditions, computational setup, and the methodology used for the cone simulation. Section III presents the results of DNS investigating boundary-layer receptivity and transition over a 7° half-angle blunt cone with a nose radius of $R_n = 5.2$ mm at Mach 8 subjected to a tunnel-like acoustic disturbance field. Finally, section IV concludes the paper with a summary of key findings and future works.

II. Flow Conditions and DNS Methodology

DNS of boundary-layer instability and transition were conducted over a 7° half-angle cone with a blunt nose radius of $R_n = 5.2$ mm, with a usable axial length of 0.883 m. To simulate BLT over the cone model within a digital wind tunnel, a precursor DNS of the full-scale axisymmetric nozzle of the Sandia HWT-8 was first conducted at a nominal operational condition of total pressure of $P_0 \approx 4692$ kPa and total temperature of $T_0 \approx 617$ K. The flow conditions (including the unsteady acoustic disturbances) at the center of the nozzle, near the nozzle exit, were extracted, modeled and used for the primary cone simulation. A detailed description of the precursor DNS of the digital wind tunnel along with the methodology for extracting the freestream acoustic disturbances is given in Liu et al. [20] and Schuabb et al. [22]. Table 1 summarizes the extracted mean preshock and postshock freestream conditions used for the cone simulation. The preshock freestream variables are denoted by $(\cdot)_{\infty,1}$ and the postshock freestream variables are denoted by $(\cdot)_{\infty,2}$ and were calculated using the inviscid Taylor-Maccoll relation. The wall temperature is 298 K, corresponding to a wall-to-recovery temperature ratio of $T_w/T_r \approx 0.537$.

The full three-dimensional compressible Navier-Stokes equations in conservation form were solved numerically in cylindrical coordinates to simulate boundary-layer instability and transition over the blunt, circular cone. The fluid considered is nitrogen with an ideal gas constant of $R = 296.7$ J/(K·kg) and within the perfect gas regime. The constitutive relations for a Newtonian fluid were employed, wherein the viscous stress tensor is linearly proportional to the rate-of-strain tensor and the heat flux vector is linearly related to the temperature gradient in accordance with Fourier’s law. The coefficient of molecular viscosity μ was determined using Keyes’ law [23], while the thermal

Table 1 Preshock and postshock freestream conditions, wall temperature, and nose Reynolds number for the DNS of the Mach 8 blunt cone.

$M_{\infty,1}$	$p_{\infty,1}$ (Pa)	$T_{\infty,1}$ (K)	$M_{\infty,2}$	$p_{\infty,2}$ (Pa)	$T_{\infty,2}$ (K)	T_w (K)	T_w/T_r	$Re_{\infty} \times 10^6$ (1/m)	Re_{R_n}
8.085	448.80	43.90	7.07	1010.02	56.14	298	0.537	12.2	63440

conductivity, κ , was calculated based on the relation $\kappa = \mu C_p / Pr$, assuming a constant molecular Prandtl number of $Pr = 0.71$. A fifth-order weighted essentially nonoscillatory (WENO) scheme [24] was used for the spatial discretization of the inviscid fluxes while a fourth-order central difference scheme was used for the viscous fluxes. Time integration was performed with a third-order, low-storage, Runge-Kutta scheme [25].

The singularity at the polar axis, arising from the $1/r$ terms in the governing equations expressed in cylindrical coordinates, was addressed by offsetting the radial grid points by half a grid spacing, following the method proposed by Mohseni and Colonius [26]. This approach eliminates the need for less accurate one-sided finite difference schemes near the axis ($r = 0$) by mapping the flow domain from $(0, R) \times (0, 2\pi)$ to $(-R, R) \times (0, \pi)$, where R denotes the radial coordinate of the computational domain. The time-step restriction arising from the reduction in azimuthal grid spacing near the centerline (or pole) is mitigated by skipping every few grid points in the azimuthal direction in the computation of the azimuthal derivative [27]. The DNS of the cone was carried out using multiple overlapping computational domains, each responsible for a different portion of the streamwise extent. This domain decomposition allows for larger time steps in downstream regions and enables the use of more grid points in the azimuthal direction, as the azimuthal arc length increases with streamwise distance. A similar multidomain DNS approach was previously employed in the simulation of boundary layer transition over a cone [28, 29]. Additional details regarding the DNS methodology can be found in our prior work for a blunt cone [20] and for a sharp cone [22].

The transitional and turbulent boundary layer over a circular cone subjected to “tunnel-like” freestream acoustic disturbances is investigated through direct numerical simulation (DNS). The computational domain is configured such that the inflow boundary lies upstream of the cone’s leading-edge shock. At this boundary, the preshock freestream condition given in Table 1 is superimposed with “tunnel-like” acoustic disturbances to replicate the disturbance environment of a conventional hypersonic wind tunnel. These disturbances are generated with a precursor DNS to simulate the acoustic field within a digital wind tunnel configured without the test article (i.e., an empty wind tunnel). Fluctuation data from the freestream region of this simulation are subsequently extracted, reconstructed, and imposed at the inflow of the DNS with the cone. No-slip and isothermal boundary conditions are imposed on the wall. To minimize acoustic reflections, the grid is gradually stretched over the last 100 axial points in each subdomain and unsteady non-reflecting boundary conditions are applied at the outlet boundaries [30]. Periodic boundary conditions are enforced in the azimuthal direction, with the computational domain spanning the full 360 degrees of the geometry.

Figure 1 illustrates the DNS configuration employed in this study. The colored contour planes show the pressure while the grayscale planes show numerical Schlieren based on the density gradient. Detailed information regarding the grid dimensions, domain extents, and resolutions for each DNS subdomain is provided in Table 2. Where N_x , N_{θ} , and N_r denote the number of grid points in the streamwise, azimuthal, and radial directions, respectively. The range of x_g specifies the Cartesian streamwise extent of each subdomain. The streamwise wall-parallel grid spacing is represented by Δx^+ , while $(r\Delta\theta)^+$ and Δz_n denote the azimuthal and wall-normal grid spacing at the wall, respectively. The cone’s leading edge is located at $x_g = -5.2$ mm. All quantities with a superscript “+” are normalized by the viscous length scale $z_{\tau} \approx 43.2$ μm evaluated at $x = 0.4$ m for consistency across the DNS domains.

III. Results

Figure 2 provides a three-dimensional view of the instantaneous wall shear stress (color) and the divergence of the velocity (grayscale) within a meridional plane at $\phi = 45^\circ$ that visualizes the broadband freestream acoustic disturbances. In the upstream region of the cone, the values of shear stress remain very close to the laminar solution, agreeing with our previous works [20, 21]. Further downstream, streamwise elongated streaks start appearing together with wave packets of high-frequency, wave trains with narrow azimuthal extent that travel approximately along the flow direction. A line plot of wall shear stress distribution along the $\phi \approx 306^\circ$ generatrix of the cone is also shown in comparison with the 2D laminar baseflow to visualize the local shear stress close to the streaks and streamwise-modulated wave

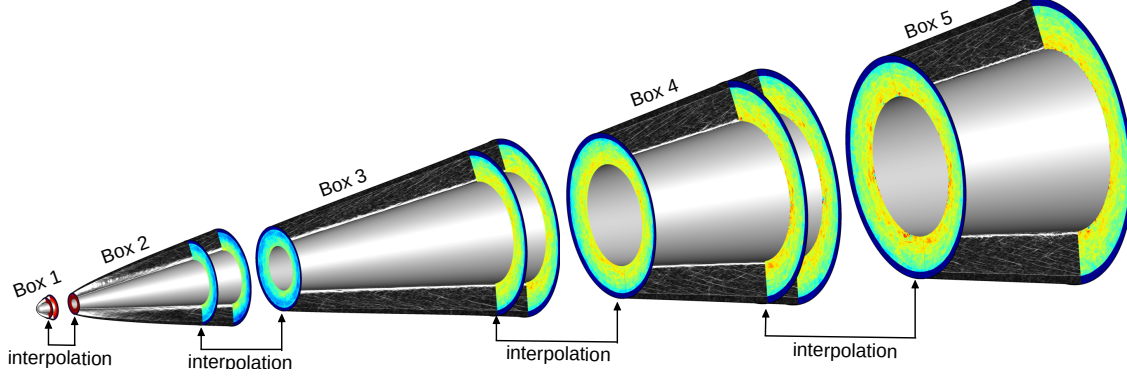


Fig. 1 Computational domain and simulation setup for the updated DNS of Mach 8 flow over a blunt 7-deg half-angle circular cone with a blunt nose radius of $R_n = 5.2$ mm. The contour plots show pressure (colored) and the numerical Schlieren based on the density gradient (grayscale).

Table 2 Domain size and grid resolution for the cone DNS runs.

Case	$N_x \times N_\theta \times N_r$	x_g (m)	Δx^+	$(r\Delta\theta)^+$	(Δz_n^+)
Box 1	$2000 \times 512 \times 1440$	$[-0.0052, 0.0105]$	$[0.22, 0.22]$	$[0.011, 1.80]$	$[0.081, 0.081]$
Box 2	$2000 \times 512 \times 1440$	$[0.00887, 0.202]$	$[0.22, 4.53]$	$[1.80, 7.86]$	$[0.081, 0.29]$
Box 3	$1700 \times 1024 \times 1440$	$[0.183, 0.495]$	$[4.53, 4.53]$	$[3.94, 9.02]$	$[0.29, 0.46]$
Box 4	$1200 \times 2048 \times 1440$	$[0.475, 0.689]$	$[4.53, 4.53]$	$[4.51, 6.38]$	$[0.46, 0.61]$
Box 5	$1200 \times 2048 \times 1440$	$[0.670, 0.883]$	$[4.53, 4.53]$	$[6.21, 8.07]$	$[0.58, 0.75]$

packet. In the meridional plane (grayscale), close to the end of the domain, it is possible to see the acoustic radiation from the transitional/turbulent boundary layer, similar to the acoustic radiation from the nozzle wall of a “noisy” wind tunnel. The line plot reveals a local maximum in the wall shear stress, at $x_g \approx 0.61$ m, due to the streaks, which are indicated by the arrow number 1 in both plots to emphasize that the local maximum is indeed due to the streak. Strong, short-scale oscillations are also visible directly downstream of this maximum. These oscillations are the footprint of the high-frequency, streamwise-modulated wave packets. At downstream locations, these oscillations increase in amplitude, highlighted by arrow number 2, and eventually, the flow breaks down to turbulence.

Figure 3 shows the spatio-temporal evolution of the wall shear stress (left column) and the wall pressure (right column) along a constant axial location at different times. For the wall shear stress in Fig. 3a, the slight increase in wall shear stress at $x \approx 0.5$ m in the topmost plot represents the streaks, similar to Fig. 2 pointed by the arrows. These streaks do not oscillate in time and instead increase rapidly in amplitude as they are convected downstream at a speed of around 500 and 600 m/s. When this local increase reaches $x \approx 0.6$ m, strong oscillations appear immediately downstream of it, and eventually breaking down to turbulence. Analogous stationary streaks that resulted from nonlinear interactions between nonmodal disturbances [12] had been found to play an important role in triggering the onset of transition. The spatio-temporal evolution of the wall pressure signal in Fig. 3b shows the axial signature of the azimuthally compact wave trains near the end of the domain, however, the boundary-layer streaks are not visible because the streaks do not generate significant pressure disturbances.

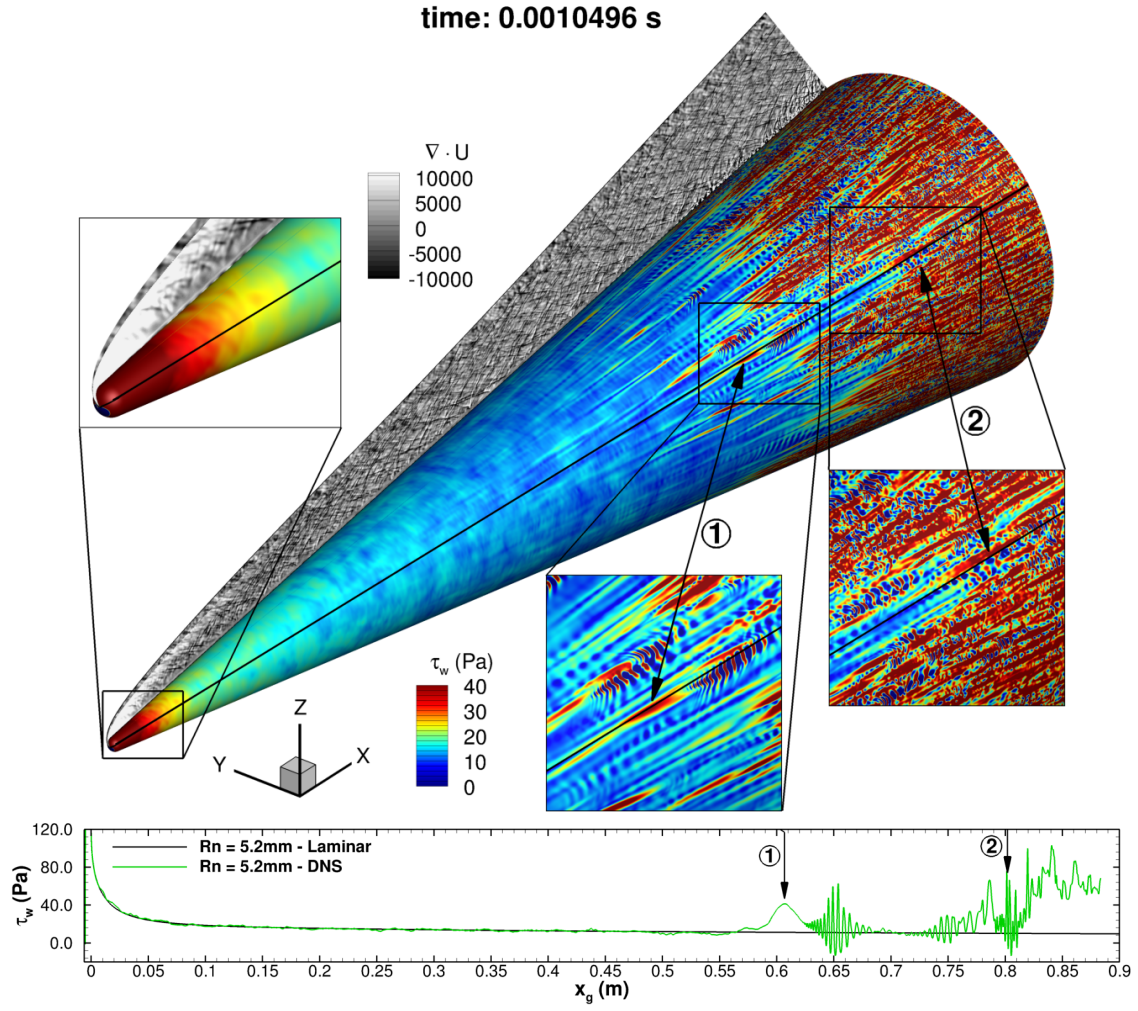


Fig. 2 Instantaneous wall shear stress (color scale) and divergence of the velocity (grayscale) at $t = 1.0496$ ms. The plot line was extracted from a row of axial locations along $\phi \approx 306$ deg meridian represented by the solid black line.

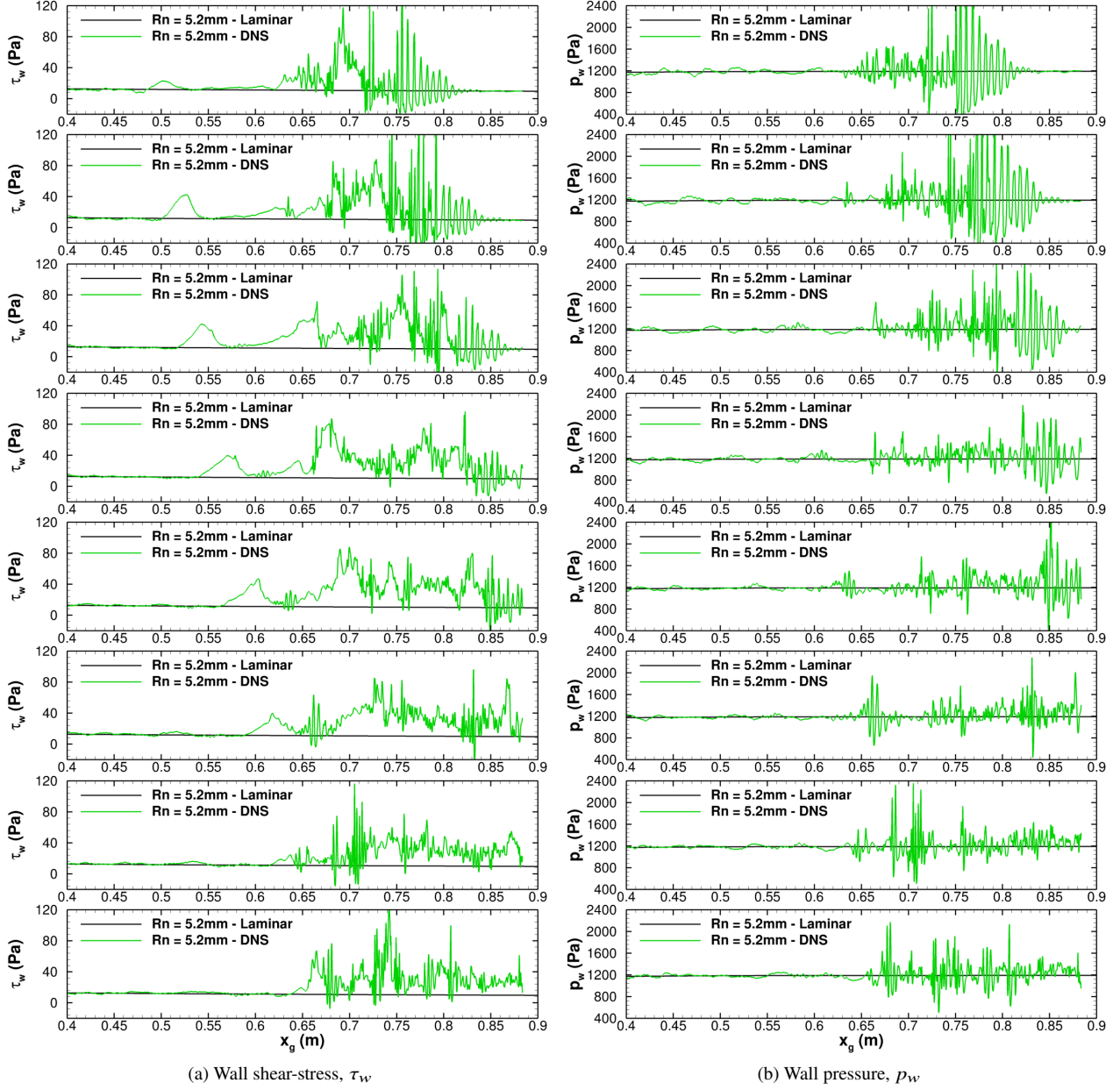


Fig. 3 Spatio-temporal evolution of (a) wall shear stress and (b) wall pressure along $\phi \approx 280$ deg meridian. The rows correspond to the times $t = [0.8256, 1.0496]$ ms with a time step of 0.032 ms.

Figure 4 shows the contours of instantaneous wall pressure P_w , wall shear stress τ_w , and wall heat flux q_w on the surface of a nearly sharp (top half) and a blunt (bottom half) cone. The sharp cone results are from the work of Schuabb et al. [22], which computed the DNS of a 0.56-m-long, 7° half-angle cone under the same freestream conditions as the current paper but with a smaller nose radius of $R_n = 0.5$ mm (the blunt cone results are for the present study). Each contour visualization is accompanied by a representative line plot extracted along the center of this contour plot (i.e., along $y_g = 0$). The sharp cone shows the emergence and growth of quasi-two-dimensional structures, indicating the dominance of the second-mode instability, which develops strong azimuthal modulation downstream, presumably due to a secondary instability. Eventually, the sharp cone boundary layer breaks down to turbulence. For the blunt cone, while the results have not reached a turbulent flow by the end of the domain, it is evident that the onset of transition has shifted considerably downstream. Moreover, contrary to the sharp cone, no quasi-two-dimensional rollers are visible in the upstream region. But we can see local streaks in the wall shear stress and wall heat flux contours.

Figure 5 shows the streamwise evolution of the spanwise average of the instantaneous values of wall pressure coefficient (C_{p_w}), the skin-friction coefficient (C_f), and the Stanton number (C_h) at $t = 0.9984$ ms, in comparison with spanwise- and time-averaged results of a sharp cone ($R_n = 0.5$ mm) from Schuabb et al. [22]. The shaded grey region represents the envelope of the spanwise variation of the instantaneous solution of the blunt cone. The 2D laminar baseflow calculated with the same DNS solver using the axisymmetric Navier-Stokes equations, and the C_f and C_h computed from the van Driest II theory [31] are included for comparison purposes. The pressure coefficient, the skin-friction coefficient, and the Stanton number are defined as

$$C_{p_w} = \frac{P_w - P_{\infty,2}}{0.5\rho_{\infty,2}U_{\infty,2}^2}, \quad C_f = \frac{\tau_w}{0.5\rho_{\infty,2}U_{\infty,2}^2}, \quad C_h = \frac{q_w}{\rho_{\infty,2}U_{\infty,2}C_p(T_r - T_w)}. \quad (1)$$

For the theoretical turbulent skin friction coefficient of a cone, a cone rule based on the von Karman momentum integral [32] was used:

$$C_{f,Cone} = C_{f,FlatPlate} \times G, \quad (2)$$

where $1.087 < G = (2 + n)^{n/(n+1)} < 1.176$ (or $1/8 < n < 1/4$). In this work, the value of n was chosen to be approximately $1/7$ with $G = 1.1$. The theoretical turbulent Stanton number was then calculated according to the Reynolds analogy as

$$C_{h,cone} = \frac{C_{f,cone}}{2Pr^{2/3}}. \quad (3)$$

From Fig. 5, the wall pressure coefficient values start to deviate from the laminar solution at $x \approx 0.32$ m for the sharp cone and ≈ 0.7 m for the blunt cone. However, for both the skin-friction coefficient and the Stanton number, the deviations start earlier at $x \approx 0.20$ m for the sharp and $x \approx 0.56$ m for the blunt cones. Moreover, it is possible to see that the rate of increase of τ_w and q_w is lower for the blunt cone than it is for the sharp cone, agreeing with previous experimental/numerical results [2, 33, 34].

Linear stability analysis based on the parabolized stability equations is performed to calculate the linear amplification of Mack's second-mode waves for the sharp and blunt cones. Figures 6a and 6c show the streamwise spectral content of the N -factor based on wall pressure while Figs. 6b and 6d show the N -factor curves for selected frequencies for both cones. The stabilizing influence of nose bluntness on Mack's second mode disturbances is clearly seen from these figures. The dominant MM frequency decreases from around 300 kHz for the sharp cone to approximately 150 kHz for the blunt cone. Defining the transition location as the intersection of the tangent line to the rising curve of the skin-friction (C_f) and heat-transfer (C_h) coefficients with the laminar baseflow curve, as schematically shown by the blue dash-dotted lines in Figs. 5b and 5c, the transition locations for the sharp and blunt cones are $x \approx 0.22$ m and $x \approx 0.56$ m, respectively. At these respective locations, increasing bluntness leads to a reduction in the N -factor, decreasing from about 6.2 in the sharp case to 1.2 in the blunt case. This observed trend aligns with previous experimental and numerical results at Mach 6, 9, and 10 [2–7, 18].

Furthermore, Fig. 7 presents the root-mean-square (r.m.s.) distributions of wall pressure, wall shear stress, and wall heat flux. The r.m.s. values are averaged over time and spanwise direction for the sharp cone, while for the blunt cone, they are only in the spanwise direction. For the sharp cone, the r.m.s. values peak near $x \approx 0.28$ m, marking the transition region. In contrast, for the blunt cone, the r.m.s. values continue to increase toward the downstream end of the domain, indicating that the peak, and thus, the full transition process has not yet been captured. This suggests that an extended computational domain is needed for the blunt case. Additionally, the gradual rise in the r.m.s. levels for the blunt cone indicates a slower transition process, consistent with Fig. 5.

Figure 8a shows the power spectral density (PSD) of wall heat flux fluctuations normalized by their mean values at $x = 0.4$ m as a function of the streamwise coordinate. The black line represents the neutral boundary of the second

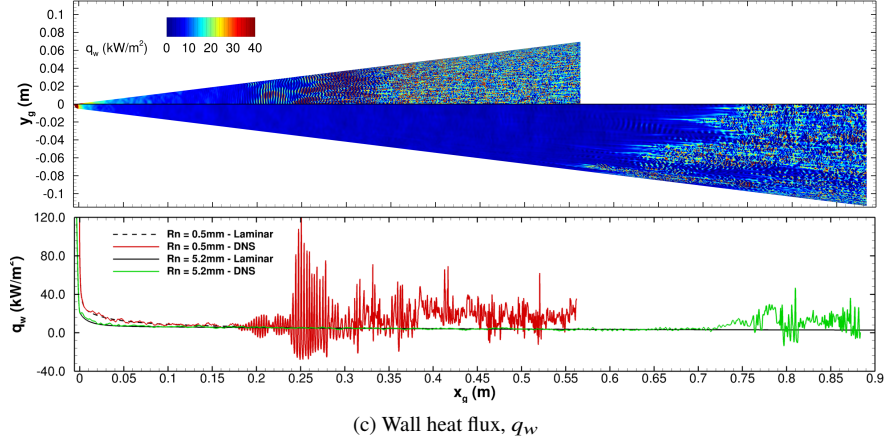
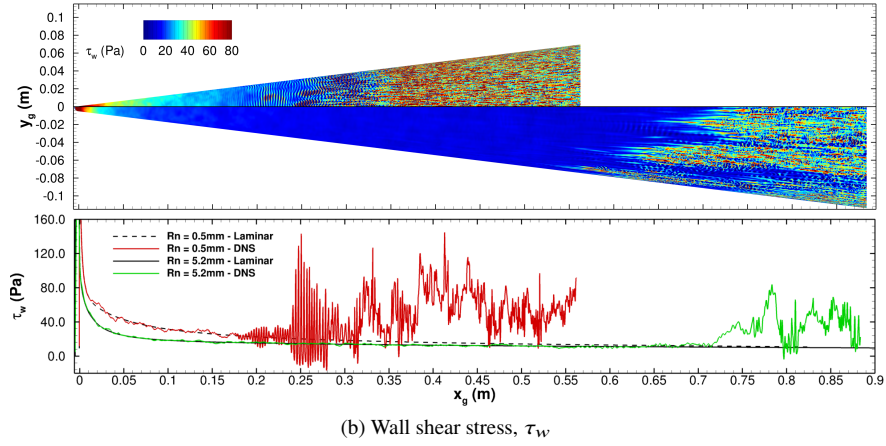
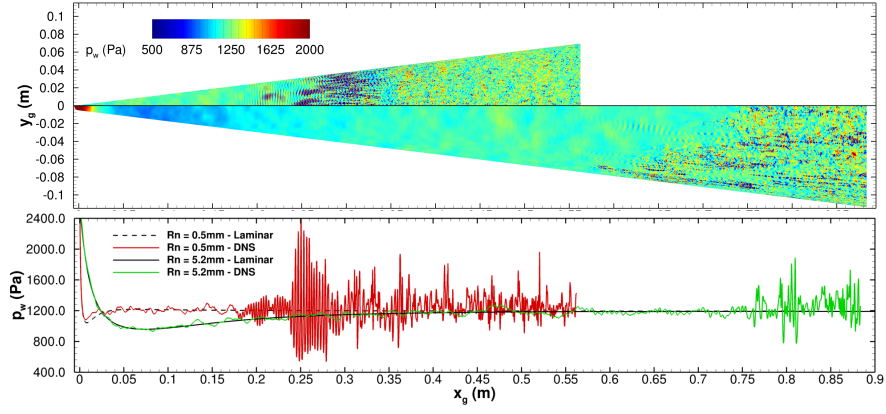


Fig. 4 Instantaneous wall pressure, wall shear stress, and wall heat flux at $t = 0.838$ ms for the sharp cone and $t = 1.0496$ ms for the blunt cone. The line plots are extracted along the center of the domain (i.e., along $y_g = 0$). The solid and dashed black lines represent laminar solutions of the corresponding geometry.

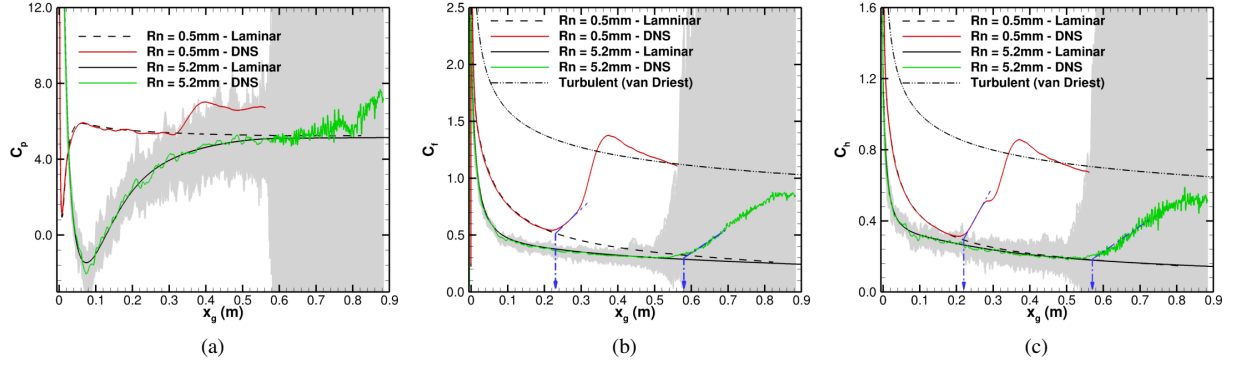


Fig. 5 Streamwise evolution of the wall (a) pressure coefficient C_{p_w} , (b) Skin-friction coefficient C_f , and (c) Stanton number C_h for the relatively sharp cone ($R_n = 0.5$ mm) and blunt cone ($R_n = 5.2$ mm) in comparison with corresponding laminar base flows. The RANS solutions for the blunt cone are also included in parts (b) and (c). DNS results for $R_n = 0.5$ mm are averaged in both time and span, whereas DNS results for the $R_n = 5.2$ mm case are only averaged in span.

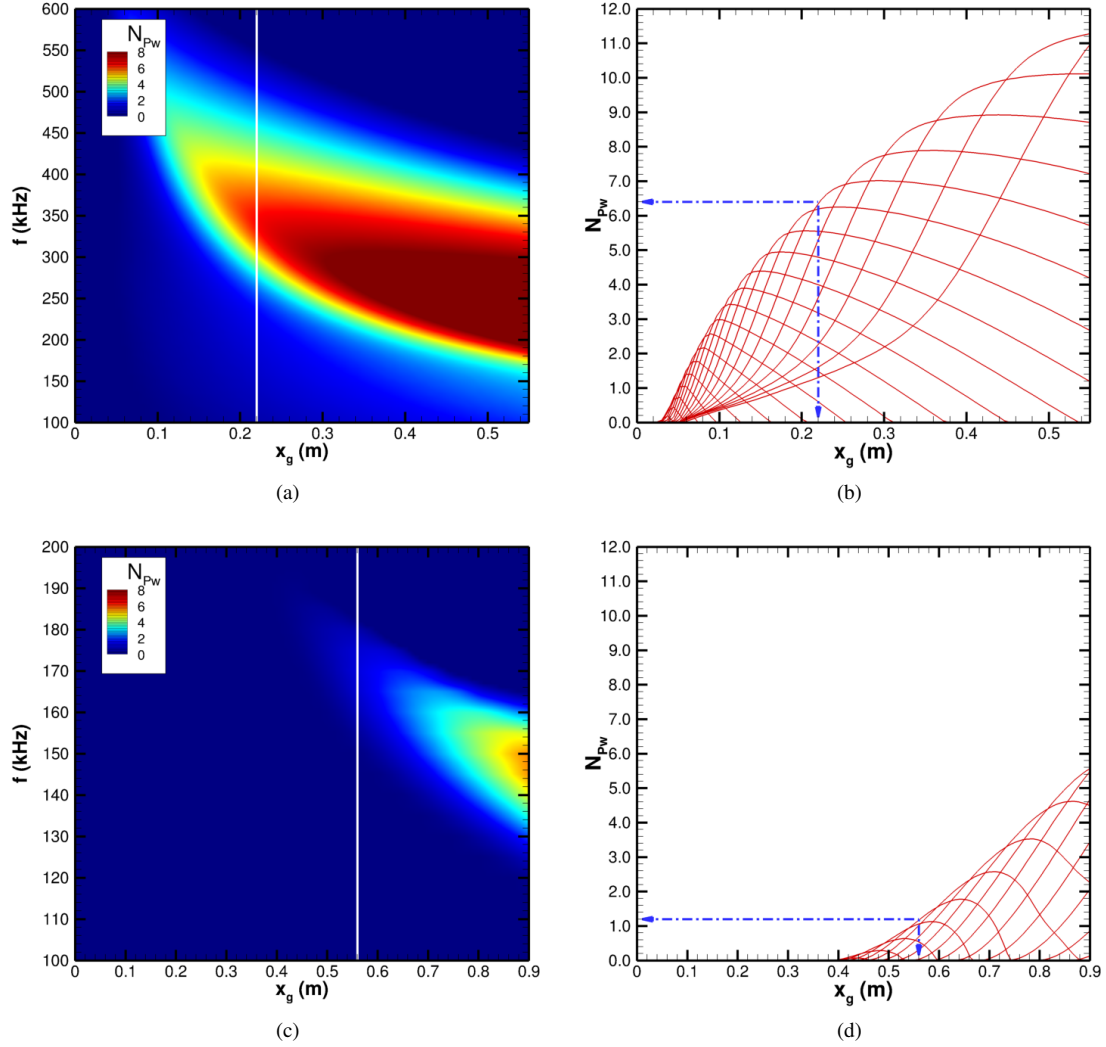


Fig. 6 N -factor of the wall pressure over the (a, b) $R_n = 0.5$ mm cone and the (c, d) $R_n = 5.2$ mm cone.

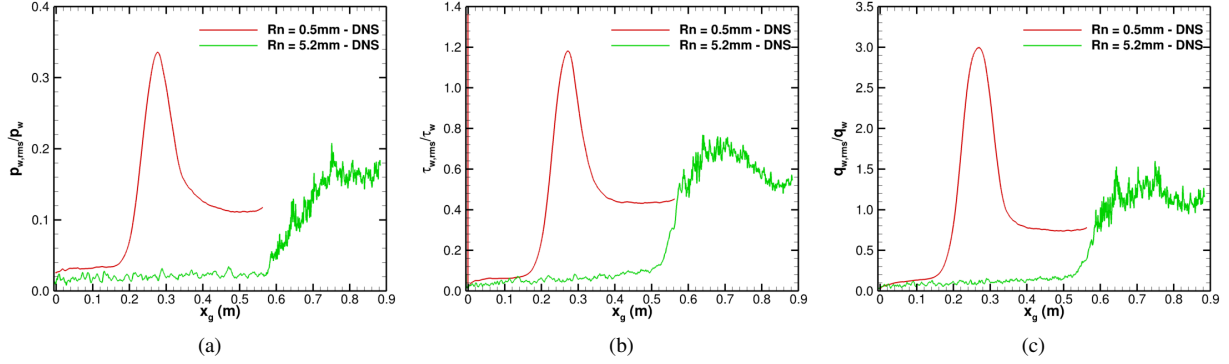


Fig. 7 Comparison of the streamwise evolution of (a) r.m.s. wall pressure, (b) r.m.s. skin-friction, and (c) r.m.s. wall heat flux against those of the DNS of the sharp cone by Schuabb et al. [22]. The r.m.s curves of the blunt cone are noisier than for the sharp cone as they were integrated over a longer time.

mode computed from the linear stability analysis, i.e., the region where the N -factor (Fig. 6c) becomes greater than zero. Similar to our previous work [21], a peak appears between 150 kHz and 200 kHz starting around $x \approx 0.20$ m. This peak persists at downstream locations, with nearly a constant value until $x \approx 0.52$ m, where the spectrum undergoes a rapid increase in amplitude as well as frequency broadening. Previous work employed nonmodal stability analysis to identify the dominant frequency associated with entropy-layer disturbances (around 175 kHz), which happens to be close to the second mode frequency [20]. During transition, interactions may occur between entropy-layer disturbances and the second mode. Figure 8b illustrates the azimuthal wavenumber spectrum of the same variables as a function of x . The spectrum is concentrated in small wavenumbers (large wavelengths) in most of the domain. Around $x \approx 0.5$, the spectrum rapidly gets broader, symptomatic of the onset of the laminar-turbulent transition.

Figure 9 shows the spatial evolution of numerical schlieren based on the instantaneous density gradient over an axial plane at $\phi \approx 280$ deg. The entropy layer and boundary layer edges as well as the sonic line of the 2D baseflow are also shown by the blue, red, and green lines respectively. Consistent with experimental observations in conventional wind tunnels for blunt straight cones [35, 36] and blunt ogive-cylinder configurations [37], the formation of inclined flow structures extending beyond the boundary-layer edge is also observed. The inclined structures observed here are similar to those documented in Refs. [20, 21]. These inclined structures are more visible for $x_g \geq 0.12$ m, between the edge of the entropy layer and the edge of the boundary layer. As the entropy layer is swallowed into the boundary layer, the structures become more apparent. Similar to our previous work [21], they exhibit minimal changes in size or shape as they are convected at a nearly constant speed close to the mean flow velocity near the boundary-layer edge, agreeing with a previous experimental study for a blunt cone in a conventional Mach 6 Ludwig tube [36]. Within the region extending from $x_g \approx 0.66$ m to $x_g \approx 0.69$ m, we can observe significant distortion of the inclined structures, and further downstream ($x_g \geq 0.76$ m), smaller-scale structures develop in the near-wall region, characterizing the breakdown process. During the breakdown, the boundary layer over the cone is seen to radiate acoustic disturbances to the free stream, similar to how a turbulent boundary layer behaves at the tunnel walls.

IV. Conclusions

This paper investigated the receptivity of a Mach 8 laminar boundary layer over a 7° half-angle blunt cone with nose radius $R_n = 5.2$ mm to broadband freestream acoustic disturbances based on a precursor digital wind tunnel simulation. Using direct numerical simulations (DNS), the study provides new insights into the interaction of freestream disturbances with the entropy layer and their role in boundary-layer transition. The three-dimensional visualizations and spatio-temporal analysis of the wall shear stress reveal the emergence of streamwise-elongated streaks and azimuthally-localized wave packets in the downstream region of the cone. These streaks grow in amplitude and ultimately break down, marking the onset of the laminar-turbulent transition. The observed behavior supports the transition scenario proposed by Paredes et al. [12], emphasizing the role of stationary streaks, arising from nonlinear interactions of linear nonmodal disturbances, as precursors to turbulence in hypersonic boundary layers. Comparing wall pressure, shear stress, and heat flux contours between a nominally sharp cone (data from Ref. [22]) and the present blunt

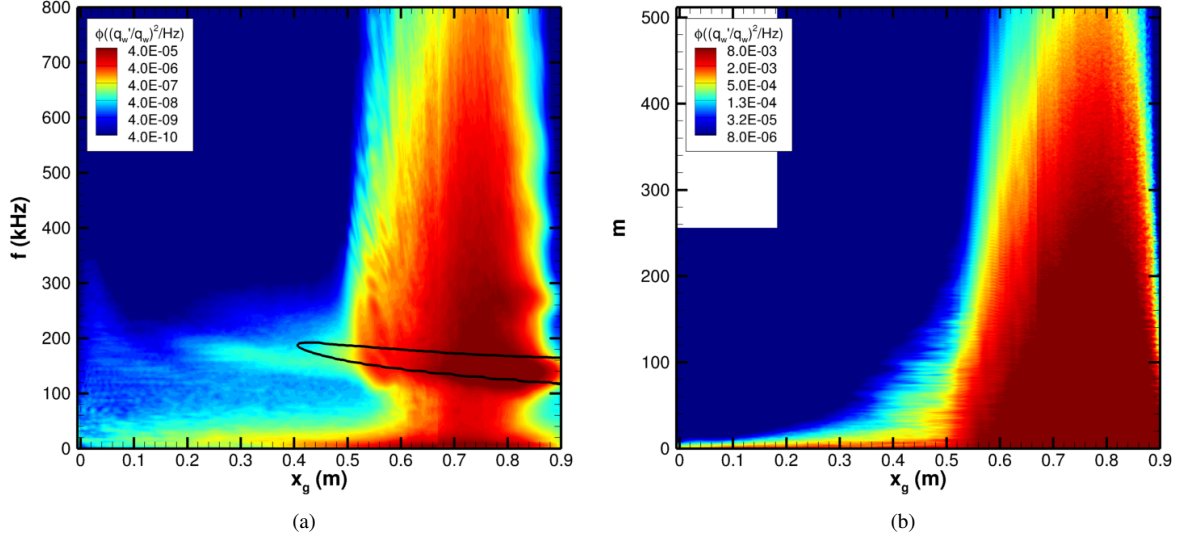


Fig. 8 Frequency (left) and azimuthal wavenumber (right) spectrum as a function of the streamwise coordinate of the wall heat flux (Note the logarithmic nature of the color map). The black line represents the neutral boundary computed from the linear stability analysis.

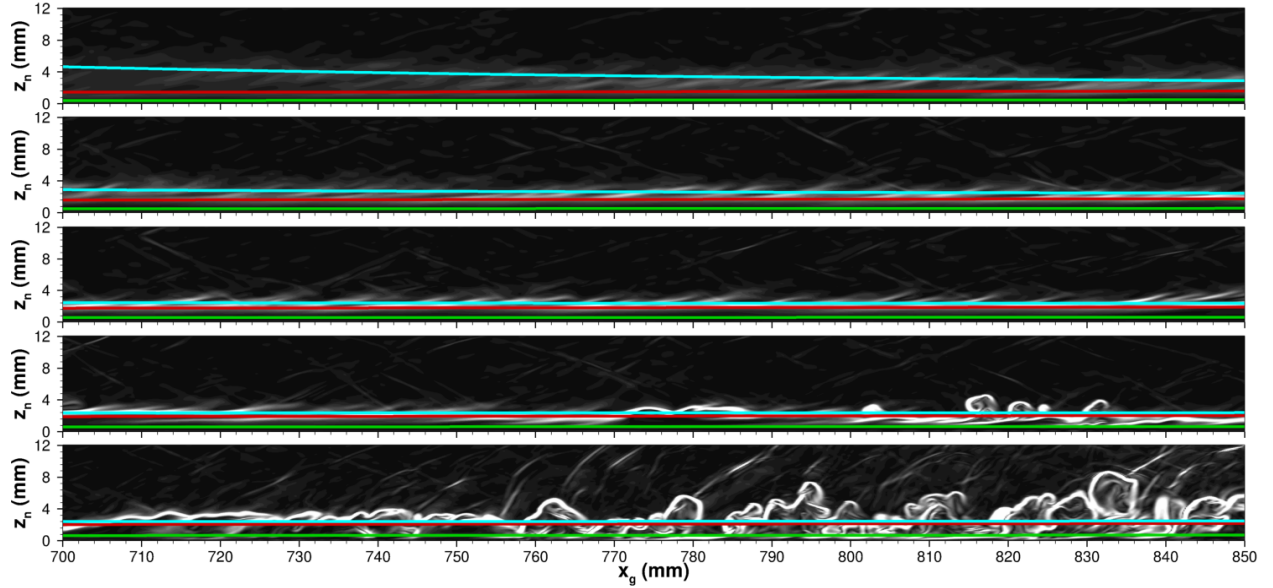


Fig. 9 Numerical (synthetic) schlieren based on the density gradient at $\phi = 280$ deg and $x_g = [0.10, 0.85]$ m at $t = 0.9984$ ms. The light blue, red, and green lines represent the entropy-layer edge, boundary-layer edge, and the sonic line, respectively.

cone reveals distinct transition mechanisms in the two cases. While the sharp cone shows early onset of second-mode instabilities followed by azimuthal modulation and breakdown to turbulence, the blunt cone exhibits delayed transition characterized by localized streaks and streamwise-modulated wave packets. The absence of quasi-two-dimensional spanwise rollers in the blunt case highlights the stabilizing influence of nose bluntness on second mode amplification. DNS results show that the boundary layer over the blunt cone remains laminar until $x \approx 0.56$ m, indicating significantly delayed departure from the laminar wall pressure, skin-friction coefficient, and Stanton number distributions. The growth rates of τ_w and q_w are notably lower for the blunt cone, indicating a more gradual transition process. Similar observations can be seen in the r.m.s. of wall pressure, shear stress, and heat flux.

The spectral analysis of wall heat flux fluctuations reveals the early emergence of a dominant frequency band between 150 kHz–200 kHz, which rapidly intensifies in amplitude and broadens to frequencies in excess of 800 kHz near $x \approx 0.52$ m, indicating transition onset. This behavior aligns closely with predictions from non-modal linear stability analysis and with previous works [20, 21]. Additionally, the azimuthal spectrum confirms the prevalence of small wavenumber disturbances throughout most of the domain, followed by a rapid broadening near the transition point, further highlighting the shift to turbulent flow. Numerical (synthetic) schlieren visualizations show inclined structures consistent with blunt cone experiments [36]. Such structures appear in the most upstream region near the edge of the entropy layer, and as they are convected downstream at speeds close to the mean flow near the boundary-layer edge, they become more evident. Downstream, they have significant distortion and, eventually, break down to smaller-scale structures characterizing the laminar-turbulent breakdown.

Future work will extend the size of the simulation’s domain downstream as the current simulation was not sufficient to capture the whole laminar-turbulent transition process over this blunt geometry.

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