

HYPERSONIC TURBULENT BOUNDARY LAYER STATISTICS

by

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ABSTRACT

A physical understanding of supersonic and hypersonic turbulence is essential for the design of high-speed vehicles, as quantities such as aerodynamic drag and heat transfer directly influence vehicle performance and survivability. Non-intrusive flow diagnostic techniques, such as krypton tagging velocimetry (KTV) and acetone tagging velocimetry (ATV), enable measurement of these turbulent quantities without disturbing the flow field. Experiments were performed in the Stevens Shock Tunnel to obtain streamwise and wall-normal velocity fluctuation profiles in Mach 6 turbulent boundary layers over a hollow-cylinder-flare test article under cold and enthalpy-matched conditions. Surface measurements of heat transfer and wall shear stress were used to characterize boundary conditions and support validation of the optical diagnostics. A novel technique was developed to measure wall-normal velocity fluctuations using a single inclined laser line, enabling access to quantities that are difficult to capture with existing diagnostics.

The measured turbulence profiles are compared with experimental and computational studies from the literature to assess the validity of Morkovin's hypothesis, which relates compressible turbulent flows to their incompressible counterparts. The results show favorable agreement with theoretical and direct numerical simulation trends for streamwise fluctuations and capture the expected behavior of wall-normal fluctuations, where previous measurement techniques have exhibited discrepancies.

These findings provide new experimental evidence supporting Morkovin's hypothesis in hypersonic turbulent boundary layers and improve the characterization of turbulence structure in high-speed flows.

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To my friends and family.

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Chapter 1

Introduction

1.1 Introduction to High-Speed Turbulence

The dream of traveling from San Francisco to Madrid in just 45 minutes, or circumnavigating the globe along the equator in under four hours, captures the ultimate promise of hypersonic flight [10]. The development of hypersonic aircraft capable of flying at several times the speed of sound holds the revolutionary potential to transform commercial aviation into fast, long-range global transport, while also enabling reusable stages for access to low Earth orbit and advanced precision-strike weapons for national security [11]. However, more than seventy years after the sound barrier was first broken, attaining stable long-range manned flight at speeds greater than Mach 5 remains a staggering challenge that continues to outpace human determination [11].

The primary hurdle is the extreme gas environment that surrounds the vehicle. As these aircraft pierce the atmosphere, they are subjected to intense aerothermal loads and extreme temperatures driven by shock waves and air friction that continuously threaten their structural integrity and maneuverability [11]. For vehicles designed to cruise through the stratosphere at Mach 6 to 10, these severe conditions are complicated by a host of outstanding fluid mechanical problems. Chief among these are the profound complexities of the highly compressible turbulent flows that develop around the airframe and within propulsion systems [11]. Ultimately, advancing our measurement capabilities to better understand this high-speed turbulence is an essential piece of the puzzle to master the thermal and structural demands of hypersonic flight. As Anderson [12] notes, “aerodynamic heating dominates the design

of all hypersonic machinery, whether it be a flight vehicle, a ramjet engine to power such a vehicle, or a wind tunnel to test the vehicle.”

1.1.1 Heat Transfer, Skin Friction, and Shock-Wave Boundary-Layer Interaction

For flight in the hypersonic regime, the boundary-layer flow developing over the vehicle surface can be turbulent, which dictates both the aerodynamic drag and the intense aerothermal loads [13]. As a high-speed boundary layer transitions from a laminar, ordered state to a turbulent, chaotic state, local heating rates can jump by up to a factor of eight [14]. Turbulent skin friction contributes a significant portion of the total aerodynamic drag, while intense aerodynamic heating determines the maximum surface temperature of the vehicle structure. Together, these effects directly influence both the survivability and operational range of hypersonic vehicles [13–15]. Uncertainty in predicting turbulent skin friction and heating results in conservative thermal protection system design, increasing structural mass, and limiting vehicle performance [14, 16].

The physical complexity increases further in regions of shock-wave boundary-layer interaction (SWBLI). SWBLIs occur at control surface deflections, flared junctions, and engine inlets of hypersonic vehicles [17]. The pressure rise across an impinging shock imposes a severe adverse pressure gradient that often induces boundary layer separation [12]. This separation and subsequent reattachment generate intense localized heating and pressure loads [18]. Accurate prediction of separation extent and peak reattachment heating in SWBLI remains a major challenge for turbulence modeling and computational fluid dynamics (CFD) [18].

1.1.2 Modeling Dependence on Accurate Reynolds Stress Prediction

A predominant means of modeling high-speed flows is the use of CFD. Owing to the wide range of spatial and temporal scales in high-Reynolds-number wall-bounded turbulence, Direct Numerical Simulation (DNS) and Large Eddy Simulation (LES) remain computationally prohibitive for routine vehicle design [6, 19]. Due to these limitations, practical engineering analysis relies primarily on Reynolds-Averaged Navier–Stokes (RANS) and Reynolds Stress Models (RSM) [20]. However, the reliability of RANS predictions hinges on the accurate closure of the Reynolds stress tensor and the turbulent heat flux [21–23]. The development and validation of such closures require accurate experimental and numerical measurements of the individual Reynolds stress components in compressible boundary layers, a fundamental understanding of the physical mechanisms governing compressible turbulence, and how these mechanisms scale relative to their incompressible counterparts.

A limitation of current RANS methodologies is that many turbulence models were originally formulated for incompressible flows and were later extended to compressible flows through density-weighted, or Favre, averaging [17, 19]. These closures often rely on the Boussinesq eddy-viscosity hypothesis, which assumes that the Reynolds stresses can be related to the mean strain rate through a scalar turbulent viscosity. They also commonly assume a constant turbulent Prandtl number to model turbulent heat flux. Under severe hypersonic conditions, where strong compressibility, high-temperature effects, shock interactions, and nonequilibrium processes may be present, the validity of these assumptions can break down [17]. Accurate prediction of near-wall compressible flows requires resolving strong anisotropy in the normal stresses and modeling pressure–strain redistribution mechanisms [22, 24]. In phenomena like SWBLI, baseline two-equation models often fail to capture non-equilibrium

turbulence dynamics, leading to systematic errors in predicting separation extent and reattachment heating [17, 18].

1.1.3 Compressibility Effects on Turbulence

A defining feature of hypersonic turbulence is the strong nonlinear coupling between the kinematic and thermal fields [25]. In compressible boundary layers, density and pressure vary significantly and couple directly to the temperature field through the thermodynamic state relations [20]. Furthermore, the extreme viscous dissipation at hypersonic conditions ($M > 5$) can create temperatures high enough to excite the vibrational energy of molecules and cause chemical dissociation, introducing “real-gas” or high-temperature effects that significantly modify the thermodynamic behavior of the boundary layer [12].

At moderate supersonic speeds (generally below Mach 5), the dynamics of the boundary layer are commonly interpreted through Morkovin’s hypothesis [3]. This hypothesis posits that, provided the fluctuating Mach number remains small, the essential turbulence dynamics retain an incompressible character and compressibility effects enter primarily through mean density variation [1, 3, 26, 27]. Under this framework, Reynolds stress profiles in compressible boundary layers should collapse onto incompressible data when properly scaled by means of local density.

When the turbulent Mach number becomes sufficiently large, velocity fluctuations can become comparable to the local speed of sound, indicating a breakdown of incompressible-like similarity and the emergence of intrinsic compressibility effects [4, 28]. These intrinsic effects are characterized by finite fluid dilatation and include the frequent appearance of localized shock waves (“shocklets”) within the boundary layer, contributing to enhanced entropy production and dilatational activity. Yu et al. [28] state as a result, explicit terms that are absent in incompressible flows, such

as dilatation work and dilatational dissipation, become non-negligible components of the turbulent kinetic energy budget [1, 4, 28]. Furthermore, the extreme wall cooling required for hypersonic vehicles exacerbates density gradients near the wall, fundamentally modifying the inter-component energy transfer of Reynolds stresses and enhancing the coherency of near-wall turbulent structures [4, 27]. These mechanisms complicate both the prediction and scaling of hypersonic turbulence. Therefore, a central question in compressible turbulence research is whether high-speed turbulent boundary layers exhibit dynamic similarity with their incompressible counterparts. Addressing this question requires a theoretical framework describing the structure and scaling of wall-bounded turbulent flows. The foundations of such a framework originate in the development of classical boundary-layer theory in the early twentieth century.

1.2 Development of Compressible Turbulence Scaling

1.2.1 Early Boundary Layer Theory

Before the turn of the twentieth century, fluid dynamics was limited by the difficulty of solving the full Navier–Stokes equations for practical flows, making quantitative prediction of wall shear stress exceedingly challenging [29]. For a Newtonian fluid, the compressible Navier–Stokes continuity and momentum equations may be written following the notation of Kundu and Cohen [30] as

$$\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x_i} (\rho u_i) = 0, \quad (1.2.1)$$

and

$$\rho \frac{Du_j}{Dt} = \rho g_j + \frac{\partial}{\partial x_i} (T_{ij}), \quad (1.2.2)$$

respectively. Kundu and Cohen [30] discuss that Eq. 1.2.2, sometimes referred to as Cauchy's equation of motion and relates the body (ρg_j) and surface T_{ij} forces on the particle. By introducing the constitutive relation, the momentum equation can be written as

$$\rho \left(\frac{\partial u_j}{\partial t} + u_i \frac{\partial u_j}{\partial x_i} \right) = -\frac{\partial p}{\partial x_j} + \rho g_j + \frac{\partial}{\partial x_i} \left[\mu \left(\frac{\partial u_j}{\partial x_i} + \frac{\partial u_i}{\partial x_j} \right) + \left(\mu_v - \frac{2}{3}\mu \right) \frac{\partial u_m}{\partial x_m} \delta_{ij} \right]. \quad (1.2.3)$$

For incompressible flow, $\partial \rho / \partial t = 0$. The continuity equation (Eq. 1.2.1) is simplified to

$$\frac{\partial u_m}{\partial x_m} = 0 \quad (1.2.4)$$

This makes the dilatational term vanish and the governing momentum equation reduces to the incompressible Navier–Stokes form defined as

$$\rho \frac{Du_j}{Dt} = -\frac{\partial p}{\partial x_j} + \rho g_j + \frac{\partial}{\partial x_i} \left[\mu \left(\frac{\partial u_j}{\partial x_i} + \frac{\partial u_i}{\partial x_j} \right) \right]. \quad (1.2.5)$$

Although this form is simpler, the equations remain nonlinear partial differential equations and are still analytically intractable for most engineering flows. The modern study of wall-bounded flows is commonly traced to Prandtl's 1904 presentation at the Third International Congress of Mathematicians in Heidelberg, where he introduced the boundary-layer concept and laid the foundation for subsequent analyses of near-wall viscous motion. In this work, Prandtl proposed that the flow of fluids with small viscosity can be divided into two distinct regions [29]. Due to the no-slip condition at the wall, viscous frictional effects are predominantly confined to a very thin region immediately adjacent to a solid surface, known as the boundary layer. Outside of this thin layer, the flow behaves as an approximately inviscid freestream

that dictates the boundary conditions of the inner layer. This insight reduced the elliptic Navier–Stokes system to the parabolic boundary-layer equations, enabling quantitative calculation of skin-friction drag and analysis of flow separation. Most importantly, Prandtl’s boundary layer framework laid the foundation for the analysis of turbulent wall flows.

While Prandtl’s boundary-layer formulation enabled the analysis of viscous wall flows, it was originally developed for laminar conditions. In practical aerodynamic applications, however, boundary layers can become turbulent, introducing chaotic velocity fluctuations. Extending the laminar boundary-layer framework to turbulent flows introduces substantial mathematical complexity. Turbulent motion can be analyzed using the Reynolds decomposition, used in RANS and experimental analyses, in which the instantaneous velocity field is separated into mean and fluctuating components,

$$u_i = \bar{u}_i + u'_i, \quad (1.2.6)$$

with the property that the fluctuations have zero mean

$$\overline{u'_i} = 0. \quad (1.2.7)$$

Substituting this decomposition into the incompressible Navier–Stokes equations and performing a time average yields additional apparent stress terms of the form

$$S_{ij} = -\rho \overline{u'_i u'_j}, \quad (1.2.8)$$

which define the Reynolds stress tensor in incompressible flow [20]. These stresses represent momentum transport by turbulent fluctuations and introduce new unknowns

into the averaged governing equations, as the equations contain second-order mean products of non-zero fluctuating quantities. The resulting system contains more unknowns than equations, giving rise to the classical closure problem [26]. Because the Reynolds stresses depend on the local turbulent structure and mean flow gradients, early modeling efforts relied on empirical formulations to relate them to mean quantities. This need to close the equations motivated the development of similarity arguments and mixing-length models, including Prandtl’s mixing-length theory.

One of the most influential early similarity concepts was the Reynolds analogy, first proposed by Osborne Reynolds in 1874, which relates momentum transport (skin friction) to heat transfer within a turbulent shear layer [20]. By assuming that turbulent eddies transport thermal energy in a manner analogous to momentum, the analogy establishes a relationship between the mean temperature and mean velocity fields. Extensions of this idea, including the Crocco–Busemann relation, further connected velocity and temperature profiles under the assumptions of constant fluid properties and unity turbulent Prandtl number. Together, these early heat-transfer and velocity-temperature similarity relations motivated a broader search for scaling frameworks capable of collapsing turbulent boundary-layer behavior across the wall-normal direction.

A central result of classical wall-bounded turbulence theory is that turbulent boundary layers exhibit distinct inner and outer scaling regions. In the near-wall region, viscous stresses and the wall shear dominate the flow, so the mean velocity is normalized using the friction velocity, $u_\tau = \sqrt{\tau_w/\rho_w}$, and the viscous length scale, $\ell_w = \nu_w/u_\tau$. This defines the inner variables following White [20] as

$$u^+ = \frac{\bar{u}}{u_\tau}, \tag{1.2.9}$$

and

$$y^+ = \frac{yu_\tau}{\nu_w} = \frac{y}{\ell_w}. \quad (1.2.10)$$

In these coordinates, the mean velocity follows the linear relation $u^+ = y^+$ within the viscous sublayer and approaches the “logarithmic law of the wall”, defined by White [20] as

$$u^+ = \frac{1}{\kappa} \ln y^+ + B \quad (1.2.11)$$

in the overlap region, where κ is the von Kármán constant and B is an empirical offset. In the outer region, the “law of the wake” from Coles [31] utilizes outer scaling with nondimensional relations such as y/δ where δ is the boundary-layer thickness. The “law of the wall” and the “law of the wake” provide one of the earliest successful descriptions of turbulent shear flow, demonstrating that the near-wall and outer regions are governed by different characteristic scales.

These formulations implicitly assumed structural similarity of turbulent motion near the wall and were derived for incompressible flows. As flight speeds increase into the supersonic and hypersonic regimes, viscous dissipation produces large temperature gradients within the boundary layer. The resulting variations in density and viscosity alter both the mean velocity and temperature profiles, violating the assumptions underlying classical arguments.

1.2.2 Emergence of Compressibility Transformations

As flow velocities extend into the supersonic and hypersonic regimes, a significant portion of the freestream kinetic energy converts to internal energy through viscous dissipation within the boundary layer [12]. Because classical incompressible similarity variables assume constant fluid properties, mean velocity profiles in compressible boundary layers no longer collapse under standard inner-layer scaling. This break-

down of incompressible similarity motivated the development of transformation-based scaling approaches [20].

Van Driest (1951) addressed this issue by introducing a density-weighted velocity transformation that accounts for variable density near the wall [15]. Rather than modifying the turbulence structure itself, the transformation redefines the mean velocity as

$$\bar{U}_{VD} = \int_0^{\bar{u}} \sqrt{\frac{\bar{\rho}}{\rho_w}} d\bar{u}, \quad (1.2.12)$$

where ρ_w denotes the wall density. When plotted in standard wall coordinates, the transformed velocity, $\bar{U}_{VD}$, approximately recovers the log law observed in incompressible turbulent boundary layers. This result demonstrated that compressibility effects on the mean flow could be largely absorbed through appropriate density weighting. The van Driest transformed velocity as well as the analytical solutions for the viscous subregion ($\bar{u}_{VD} = y^+$ and the “law of the wall” ($u^+ = \frac{1}{\kappa} \ln y^+ + B$) are shown in Fig. 1.1 [1].

1.2.3 Morkovin’s Hypothesis

Following the success of mean-flow transformations, similar logic was extended to turbulence statistics. Standard inner-layer normalizations of Reynolds stresses and velocity fluctuations did not collapse compressible data, particularly at elevated Mach numbers [26]. In 1962, M. V. Morkovin formulated a hypothesis that has since become central to the interpretation of compressible wall-bounded turbulence [3]. He proposed that for moderate freestream Mach numbers (typically $M < 5$), turbulent fluctuations remain primarily incompressible in character. Morkovin did not argue that compressibility effects are negligible; rather, he suggested that the essential tur-

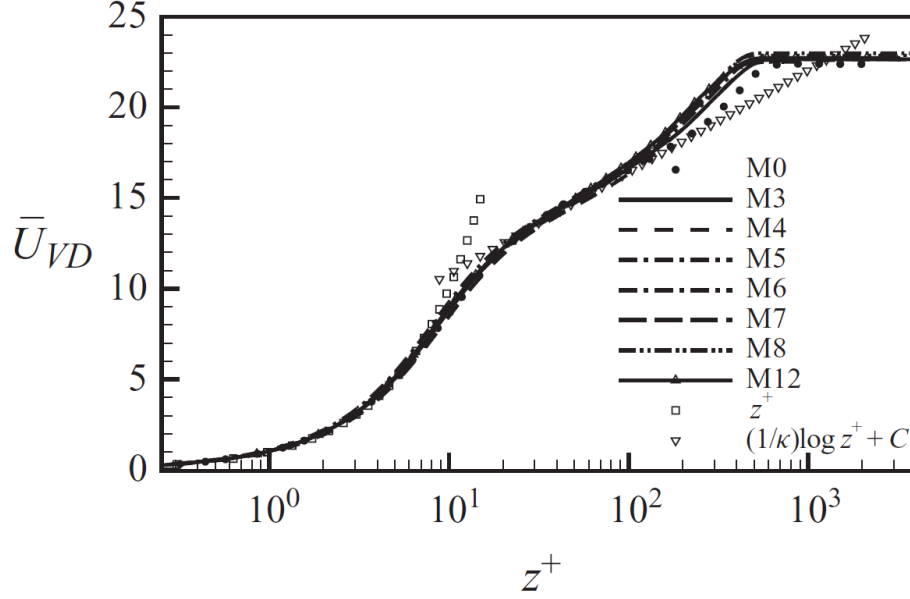


Figure 1.1: Van Driest transformed velocity for different freestream Mach number cases from Duan et al. [1]. The wall-normal coordinate z^+ corresponds to y^+ defined in the text. The analytical solutions for the viscous subregion ($\bar{u}_{VD} = y^+$ and the “law of the wall” ($u^+ = \frac{1}{\kappa} \ln y^+ + B$) are depicted as squares and upside-down triangles respectively.

bulence dynamics resemble those of incompressible turbulence when properly scaled by local mean properties. Under this interpretation, compressibility influences the flow primarily through variations in mean density and viscosity across the boundary layer, while the turbulence structure remains largely unchanged.

This framework rests on the weak compressibility assumption, which according to Dussauge et al. [32] requires the turbulent Mach number to remain small. Duan et al. [1] defines the turbulent Mach number (M_t) as

$$M_t = \frac{(\overline{u'_i u'_i})^{1/2}}{\bar{a}}, \quad (1.2.13)$$

where $(\overline{u'_i u'_i})^{1/2}$ is the root-mean-square (RMS) of the fluctuating velocity quantity and $\bar{a}$ is the mean sound speed. When $M_t \ll 1$, the turbulent field is dominated

by solenoidal (vortical) motions, and dynamic coupling between pressure fluctuations and dilatational modes remains weak [32, 33]. Smits and Dussauge [26] take $M_t \approx 0.3$ as the threshold at which compressibility effects become important for turbulent behavior. Below this threshold, density fluctuation energy is negligible relative to velocity fluctuations, and compressible turbulence statistics may be collapsed using density-weighted scaling such as

$$u'_{MS} = \frac{u'}{u_\tau \sqrt{\rho_w / \bar{\rho}}}. \quad (1.2.14)$$

Here, u'_{MS} is the Morkovin-scaled velocity fluctuation and $u_\tau = \sqrt{\tau_w / \rho_w}$ is the friction velocity. Here, the factor $\sqrt{\rho_w / \bar{\rho}}$ accounts for local density variation within the boundary layer. Under this formulation, compressible turbulence profiles were observed to collapse toward incompressible distributions across a range of Mach numbers up to Mach 5 [5]. Dussauge et al. [32] suggested based on these observations that compressibility might influence turbulent boundary layers primarily through mean thermodynamic variations rather than through fundamental changes in turbulence dynamics.

When M_t remains small, dilatational contributions to dissipation are limited and the assumption of a predominantly vortical, viscous-dominated cascade remains valid. However, as the freestream Mach number increases into the hypersonic regime, the assumptions underlying Morkovin's hypothesis become increasingly strained. Direct numerical simulations indicate that the turbulent Mach number grows within the boundary layer and may attain values approaching order unity in the outer region [1]. The influence of freestream Mach number on M_t throughout the boundary layer can be seen in Fig. 1.2. It can be seen that M_t surpasses the threshold for Mach 5 freestream, and continues to grow with increasing Mach number. As M_t increases,

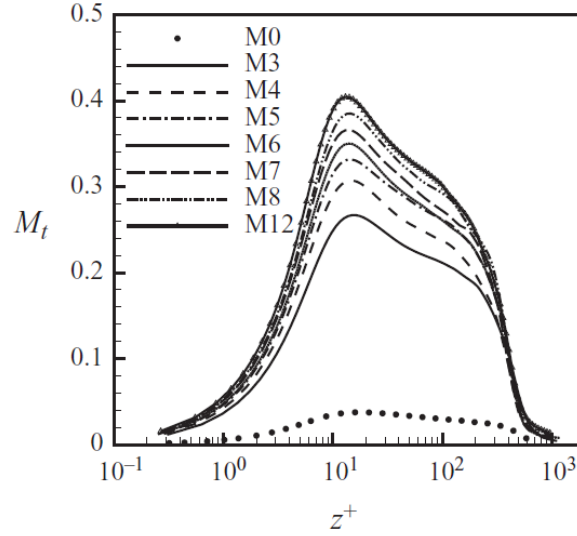


Figure 1.2: Simulation results from Duan et al. [1] of turbulence Mach number for various freestream Mach numbers

fluctuating velocities approach the local speed of sound, and intrinsic compressibility effects emerge. These include finite dilatation, stronger pressure–density coupling, and the formation of localized shocklets within turbulent boundary layers [26]. In the compressible turbulent kinetic energy budget, such effects introduce additional terms, including pressure-dilatation term, defined by Duan et al. [1] as

$$\Pi_d = \overline{p' \left(\frac{\partial u_i''}{\partial x_i} \right)}, \quad (1.2.15)$$

which directly couples pressure fluctuations with velocity divergence. Duan et al. [1] asserts as the dilatational kinetic energy increases and shocklets contribute to dissipation, the vortical dominance assumed under weak compressibility diminishes, thereby helping determine the conditions under which Morkovin’s hypothesis breaks down. Under these conditions, purely density-weighted incompressible scaling may no longer fully describe the turbulence dynamics. Figure 1.3 shows the increasing prevalence of Π_d as freestream Mach number increases. Determining the Mach number

range of validity of Morkovin’s hypothesis therefore requires detailed information about turbulence statistics in compressible boundary layers.

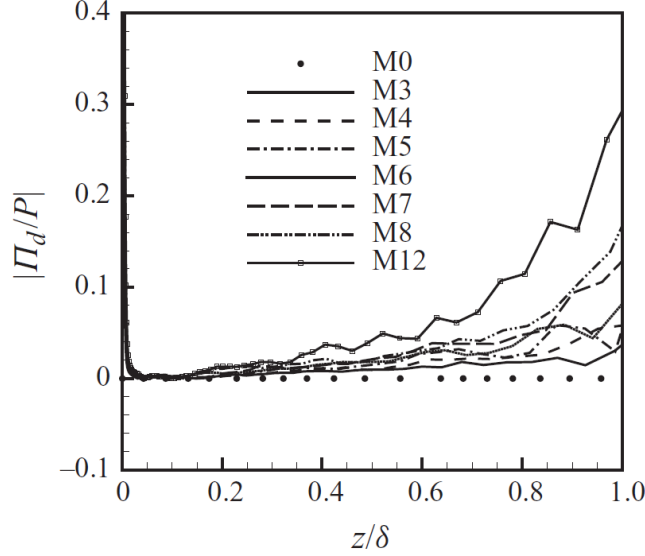


Figure 1.3: Pressure dilatation Π_d normalized by production $P = -\overline{\rho u_i'' w''} \left(\frac{\partial \tilde{u}_i}{\partial z} \right)$ from Duan et al. [1]

A direct consequence of this “weak compressibility” assumption is its application to the Kolmogorov energy cascade. In classical incompressible turbulence, kinetic energy is transferred from large, energy-containing eddies down through an inertial subrange, eventually reaching the smallest scales where it is dissipated by viscosity. As detailed by Smits and Dussauge [26], Morkovin’s hypothesis implies that this fundamental cascade mechanism remains largely intact in supersonic flows [26]. The primary modification required is that the Kolmogorov length and velocity scales must be redefined using the local mean density, mean viscosity, and mean dissipation rate per unit volume.

1.3 DNS of Compressible Boundary Layers

To overcome the inherent difficulties and uncertainties of experimentally measuring hypersonic boundary layers, researchers have increasingly relied on Direct Numerical Simulation (DNS). DNS solves the full unsteady, three-dimensional Navier–Stokes equations without employing any turbulence models or closure approximations. By directly resolving all spatial and temporal scales of turbulent motion, from the large energy-containing eddies down to the small dissipative Kolmogorov scales, DNS provides complete access to the instantaneous velocity, thermodynamic, and pressure fields [20, 22]. This makes it an invaluable tool for evaluating compressible similarity theories and extracting detailed turbulence statistics.

However, the application of DNS to high-speed wall-bounded flows is constrained by several computational and physical limitations. The number of computational grid points required to resolve all turbulent scales in DNS increases rapidly with the Reynolds number [13]. Therefore, most DNS databases for hypersonic boundary layers are strictly limited to low or moderate friction Reynolds numbers of $\text{Re}_\tau < 600$; however, Huang et al. was able to use long domains to allow for $\text{Re}_\tau \approx 1200$ [34]. Furthermore, to accurately capture near-wall structures as well as large-scale superstructures, computational domains must be large while maintaining extremely fine grid spacing near the wall to capture the most energetic and dissipative turbulence. Boundary layer simulations must also utilize long streamwise domains (often > 50 to 300 inflow boundary-layer thicknesses) to ensure the turbulence recovers from artificial inflow generation and achieves a fully developed equilibrium state [4].

Numerical algorithms in compressible DNS must also navigate a delicate balance when it comes to shock capturing. They must utilize schemes with low numerical dissipation to accurately preserve turbulent fluctuations, yet be robust enough

to capture localized flow discontinuities, such as shocklets, without introducing numerical oscillations [26]. Optimized high-order shock-capturing schemes, such as the seventh-order weighted essentially non-oscillatory (WENO) method, are frequently employed to compute the convective flux terms. To make these massive computations tractable, the majority of high-Mach-number DNS databases model the fluid as a perfect gas with constant specific heats, using Sutherland’s [1, 4, 22] or Keyes’ [19] laws for temperature-dependent viscosity. This assumption often neglects the high-temperature, real-gas effects, such as vibrational excitation and chemical dissociation, that characterize actual hypersonic flight.

Despite these computational constraints, extensive DNS databases covering Mach numbers from 2.5 up to 14 have been highly successful in testing Morkovin’s hypothesis and related compressibility transformations [4]. When Morkovin’s density-weighted scaling is applied to the DNS data, the streamwise turbulence intensities collapse well onto incompressible profiles across a wide range of Mach numbers and wall-cooling conditions [1, 27, 35]. DNS studies strongly support the validity of Morkovin’s hypothesis for the streamwise turbulence component even at hypersonic speeds. However, the wall-normal and spanwise turbulence intensities exhibit more complex behavior that depends strongly on both Mach number and wall cooling conditions. DNS results from Huang et al. [34] indicate that intense wall cooling and high Mach numbers increase the peak magnitude of the streamwise normal stress ($\rho \overline{u'u'}/\tau_w$), while simultaneously reducing the peak magnitudes of the spanwise and wall-normal turbulence intensities. The near-wall peaks of these components also shift away from the wall when expressed in classical inner scaling as the wall-cooling rate increases, although this shift can be largely accounted for using semi-local scaling coordinates. These trends highlight that while the streamwise turbulence component follows the scaling predicted by Morkovin’s hypothesis, the behavior of the wall-

normal and spanwise components remains more nuanced and continues to motivate further investigation.

A critical discrepancy in the literature is the failure of classical analytical curves and transformations to accurately predict the skin friction coefficient, C_f , under hypersonic, cold-wall conditions. While early theories were successful for adiabatic or moderately cooled walls, recent high-fidelity DNS databases have exposed their limitations. Huang et al. [34] systematically evaluated classical skin friction theories, specifically the van Driest II and Spalding & Chi transformations, against their DNS database of spatially evolving boundary layers up to Mach 14. They found that while these transformations work reasonably well for moderately cold walls ($0.3 \leq T_w/T_R \leq 1$), both classical theories fail to provide good predictions for highly-cooled walls ($T_w/T_R \leq 0.3$), where T_R is the recovery temperature and defined as

$$T_R = T_e \left[1 + \left(\frac{\gamma - 1}{2} \right) r M_e^2 \right], \quad (1.3.1)$$

where r is the turbulent recovery factor defined as

$$r = Pr^{1/3}. \quad (1.3.2)$$

Specifically, the van Driest II transformation overpredicts C_f by 10% to 20% under these extreme cold-wall conditions and the Spalding & Chi theory underpredicts C_f by 20% to 40%. There is a Reynolds number dependence in this discrepancy: the underprediction of the Spalding & Chi theory decreases gradually as the Reynolds number increases, indicating that boundary layers need a sufficient domain downstream of transition to relax into a fully turbulent state before the theory can be reliably applied. In efforts to correct the analytical curves, Zhao and Fu [36] assert

the fundamental defect of these older theories is that they rely on a linear transformation to eliminate the effects of Mach number and heat transfer on the momentum thickness Reynolds number (Re_θ). This linear approach does not completely absorb the severe density and viscosity variations present in hypersonic cold-wall boundary layers. To resolve this, a proposed redefinition of the Reynolds number based on a transformed momentum thickness (Re_θ^*) was presented. By integrating a new velocity transformation, referred to as the Generalized Reynolds Analogy, into the momentum thickness, their new general scaling law successfully collapses the C_f of highly cooled hypersonic boundary layers back onto the universal incompressible curve, drastically reducing the errors seen in the analysis of Huang et al. [34].

Despite the insight provided by DNS into the fundamental dynamics of compressible turbulence, the extreme computational cost associated with resolving high-Reynolds-number hypersonic boundary layers limits its applicability to realistic flight conditions, necessitating carefully designed ground-based experiments for validation. For this reason, experimental measurements remain essential for validating both numerical simulations and turbulence scaling theories in high-speed boundary layers.

1.4 Experimental Measurements of High-Speed Turbulence

While DNS provides detailed insight into turbulence dynamics, experimental measurements remain essential for validating numerical predictions and turbulence scaling theories. Obtaining accurate turbulence measurements in compressible boundary layers is particularly challenging due to the extremely small spatial and temporal scales of turbulent fluctuations and the severe thermodynamic conditions present in high-speed facilities. These flows exhibit high-frequency velocity fluctuations, strong temperature gradients, and large density variations across the boundary layer. De-

spite these difficulties, experimental measurements provide critical information about the structure and behavior of high-speed turbulent flows that cannot yet be fully reproduced numerically.

Experimental diagnostics are generally categorized into intrusive methods, which place a physical probe into the flow, and non-intrusive optical methods, which rely on lasers and cameras. Crucially, measurement bias and resolution limitations directly dictate the interpretation of compressible similarity theories and scaling laws; inaccurate measurements can obscure whether a scaling law is physically invalid or if the data itself is flawed [5]. Consequently, a wide range of experimental diagnostics have been developed to measure turbulent fluctuations in high-speed boundary layers, each with distinct advantages and limitations.

1.4.1 Intrusive Diagnostics

Intrusive techniques involve placing a physical probe directly into the flow field. While these methods have provided the foundational understanding of turbulence, their physical presence disturbs the high-speed flow they intend to measure. Nevertheless, surface-mounted intrusive sensors remain an important diagnostic tool in hypersonic facilities, providing measurements of boundary conditions such as heat transfer and wall shear stress that are essential for interpreting boundary-layer behavior.

Hot-Wire Anemometry (HWA) has historically been the foundational tool for measuring turbulence structure and statistics in boundary layers. The technique measures velocity fluctuations by detecting the heat transfer rate from a heated, fine wire, which is often as thin as 2.5 to 5 μm , placed directly in the flow. HWA offers excellent frequency response (often extending to hundreds of kHz) and is capable of capturing small-scale, high-frequency turbulent velocity fluctuations. Because of its high temporal and spatial resolution, it has been historically crucial for resolving the

structural features of boundary layers and establishing early incompressible similarity laws.

However, the application of HWA in high-speed, compressible flows faces severe limitations. In supersonic and hypersonic regimes, hot-wires exhibit a complex mixed-mode sensitivity, meaning their output responds simultaneously to both mass-flux fluctuations and total temperature fluctuations. This introduces significant calibration challenges [5, 37]. Additionally, HWA is a highly intrusive technique. The probes themselves can induce shock waves or flow perturbations. In some hypersonic testing environments, extreme aerodynamic heating, high dynamic pressures, and particulate impacts lead to frequent wire breakage [18]. Spatial filtering is also a considerable source of error; if the wire length is larger than the local turbulent eddy scales (which shrink drastically due to near-wall density gradients in cold-wall hypersonic flows), the measured turbulence intensities are artificially attenuated [6]. These limitations restrict the applicability of HWA in high-enthalpy hypersonic facilities.

Klebanoff utilized HWA to create the first full set of incompressible turbulent statistics in a low-speed, zero-pressure-gradient boundary layer [37]. Using 2.5 μm platinum hot-wires in both a single wire and “X-wire” configuration, Klebanoff resolved streamwise, wall-normal and spanwise velocity fluctuations (u', v', w' respectively), as well as Reynolds shear stresses ($\overline{uv}$). To overcome the resolution limitations with the wires, he thickened the boundary layer by tripping the flow upstream of the measurement location on the flat plate. This forced the boundary layer to grow to 76 mm by the time it reached the HWA probes. Although tripping the flow made the HWA measurements possible, it also created “non-equilibrium turbulence” effects in the flow. Smits [38] showed that the boundary layer in Klebanoff’s experiment was still adjusting back from the step-change in roughness which caused his measured turbulent stresses in the outer layer to fall roughly 25% below the modern consensus

established by recent DNS and high-resolution optical measurements.

1.5 Surface-based Diagnostics

Surface-based diagnostics provide another method of measurement. These include wall shear sensors, pressure transducers, thermocouples, and heat transfer gauges. Mounted flush to the model surface, these sensors measure key surface quantities such as wall shear stress, surface pressure fluctuations, and local convective heat flux. They are particularly useful for detecting the boundary layer state, such as identifying the onset and end of laminar-to-turbulent transition by recording the sharp spikes in heat transfer or pressure fluctuations as turbulence forms and grows. The primary strengths of surface-based diagnostics are their robustness and high bandwidth. Unlike fragile hot-wires, some flush-mounted sensors can survive the severe aerothermal loads and dynamic stresses of high-enthalpy hypersonic shock tunnels without failing [18]. They provide highly reliable measurements of boundary conditions that are very useful for validating CFD codes. The major limitation of these measurements is that they only provide information at the wall, they cannot resolve velocity fluctuations within the boundary layer profile, nor can they provide any information regarding the spatial structure or energy cascade of the turbulence.

Wall heat flux measurements in reflected shock facilities are often made according to 1-D semi-infinite solid simplifications due to short test durations. A range of measurement techniques following this assumption have been evaluated [39]. The means of measuring heat flux on this time-scale include temperature sensitive paints (TSP) [40, 41], thin film sensors [18, 42, 43], and coaxial thermocouples [44, 45]. Temperature sensitive paint is relatively easy to implement as the only materials needed are a coat or a substrate film of paint, a light source for thermal quenching and a

high-speed camera. TSP measurements have been successful in obtaining heat-flux data in hypersonic shock tunnel facilities [46]. Thin film gauges and coaxial thermocouples are the state-of-the-art in current surface measurement techniques. Discrete sensors such as coaxial thermocouples and thin film arrays with small surface areas benefit from high frequency response to capture transient flow physics. Facilities, such as CUBRC, instrument numerous in-house thin film sensors to their models, citing the sensor's small sensing element size and insulating substrate as factors that make them ideal for resolving spatial and temporal heat transfer of test article surfaces [43]. CUBRC also used off-the-shelf Type-E coaxial thermocouples from Medtherm Corporation as these thermocouples are specialized in measuring surface temperature changes on hypersonic ground test articles. The thermocouples suffer from lower signal-to-noise ratio (SNR) at low heat transfer rates, but are considered robust.

Measurements of wall shear stress are also highly coveted by the hypersonic community. Recent advancements by Ahmic Aerospace in coaxial shear stress gauge construction have reduced uncertainty due to sensor misalignment and pressure driven moments on the sensing element [47, 48]. These sensors function using an embedded cantilevered beam positioned flush with the test article to capture only shear forces along the article wall. Strain gauges at the base of the beam capture the deflection as flow passes over a floating element attached to the top of the beam. Temporal response and minimum shear stress to exceed the sensor noise floor must be known for application in ground-test facilities. Ahmic wall shear stress gauges have been tested for flows with shear stress values for blowdown and in flight testing with second-long response times, with a range of 1-340 Pa with reported skin friction uncertainties less than 10% [48]. Therefore, it remains uncertain whether these gauges can capture transient turbulent phenomena within the short test windows typical of shock tunnels. While intrusive diagnostics and surface sensors have provided valuable measurements

of boundary layer properties, their limitations in high-speed environments provide motivation for alternative techniques.

1.5.1 Particle-Based Optical Diagnostics

In response to the limitations of intrusive measurement techniques, a variety of non-intrusive optical diagnostics have been developed to investigate turbulent structures in high-speed boundary layers. Laser Doppler Velocimetry (LDV) is a foundational particle-based optical technique that measures the velocity of a flow at a single point. LDV achieves this by evaluating the Doppler shift of laser light scattered by particles crossing an interference fringe pattern generated by intersecting laser beams. This technique provides high accuracy and good temporal resolution without the physical intrusion of a probe. It circumvents the mixed-mode sensitivity uncertainties associated with hot-wire anemometry in compressible flows. However, LDV requires the flow to be seeded with tracer particles. In high-speed flows, fluid inertia, compressibility, and non-continuum slip effects can cause these particles to lag behind the actual turbulent fluid motion, artificially low-pass filtering the velocity measurements [35]. The LDV study by DeGraaff and Eaton [6] is widely considered a modern standard for incompressible flat-plate turbulent boundary layers, effectively replacing older, flawed benchmarks like the HWA analysis by Klebanoff in 1955. Measuring the Reynolds shear stress accurately is difficult with standard LDV systems. DeGraaff and Eaton developed a custom-built, two-component, high-resolution Laser Doppler Anemometer (LDA). Their system formed a measurement volume using three beams, each with a different frequency shift. By using a Fourier transform, their system was able to extract both the streamwise and wall-normal velocity information simultaneously from a single Doppler burst. DeGraaff and Eaton performed their measurements in a wind tunnel enclosed in a pressure vessel, allowing them to capture data across

a range of momentum thickness Reynolds numbers, from $Re_\theta = 1430$ up to 31,000. Because of their high-fidelity LDA system and the wide range of Reynolds numbers tested, the study provided strong experimental evidence regarding turbulence scaling, showing that the near-wall peak of the streamwise velocity fluctuations increases with Reynolds number.

Particle-based techniques such as Particle Image Velocimetry (PIV) provide high-quality measurements of both mean and fluctuating streamwise velocity [5, 7, 16, 49]. Unlike single-point diagnostics such as hot-wire anemometry (HWA) or Laser Doppler Velocimetry (LDV), planar PIV provides instantaneous two-dimensional velocity measurements (or three-dimensional measurements via stereo-PIV) over an entire field of view. This is accomplished by seeding the flow with tracer particles and tracking their displacement between sequential image frames. Cross-correlation algorithms are often used to determine particle displacement between frames, yielding the macroscopic velocity vector field.

The ability to obtain full-field velocity measurements is particularly advantageous because it enables the computation of spatial correlations, identification of large-scale coherent structures, and direct estimation of Reynolds stresses throughout the boundary layer. However, the accuracy of PIV measurements depends strongly on particle dynamics. A key parameter governing particle response is the Stokes number (St), which is the ratio of particle response time to a characteristic flow time scale and defined by Samimy and Lele [50] as

$$St = \frac{\tau_p}{\tau_f} \ll 1 \approx 0.1 \quad (1.5.1)$$

where

$$\tau_p = \frac{d_p^2 \rho_p}{18\mu} (1 + 2.76Kn_p), \quad (1.5.2)$$

$$Kn_p = \frac{M}{Re_p} \sqrt{\frac{\gamma\pi}{2}}, \quad (1.5.3)$$

and

$$\tau_f = \frac{\delta}{U_\infty}. \quad (1.5.4)$$

Samimy and Lele [50] identified the Stokes number as a fundamental source of measurement error, stating that “the slip between particle and fluid imposes a fundamental limit on the accuracy of optical measurements such as LDV and PIV.”

Brooks et al. [2] examined the relative error in wall-normal Morkovin-scaled fluctuating velocity between nine PIV data sets and the canonical hot-wire measurements of Klebanoff [37] at $y/\delta = 0.2$, where δ is the boundary-layer thickness. Previous studies by Lowe et al. [51], Williams et al. [35], and Brooks et al. [52] demonstrated that particle lag reduces the magnitude of the wall-normal velocity fluctuations, v' . As shown in Fig. 1.4, the discrepancy between the wall-normal fluctuations measured by PIV and the Klebanoff HWA data are approximately 10% for $St \approx 0.25$, but exceeds 50% for $St \approx 1$.

Additionally, energy at higher wavenumbers often lies beyond the measurable range of PIV due to the combined effects of particle lag and finite spatial resolution, both of which act as low-pass filters on the velocity signal. Because the spectral distribution of the wall-normal fluctuating velocity is relatively flatter than that of the streamwise component, it is more susceptible to attenuation from these filtering effects [5].

To further illustrate the impact of particle lag on wall-normal velocity measurements, Fig. 1.5 compares experimental streamwise and wall-normal fluctuations measured by Williams et al. [5] with incompressible LDV data from DeGraaff and Eaton [6] and DNS results from Zhang et al. [4] for nominally similar turbulent bound-

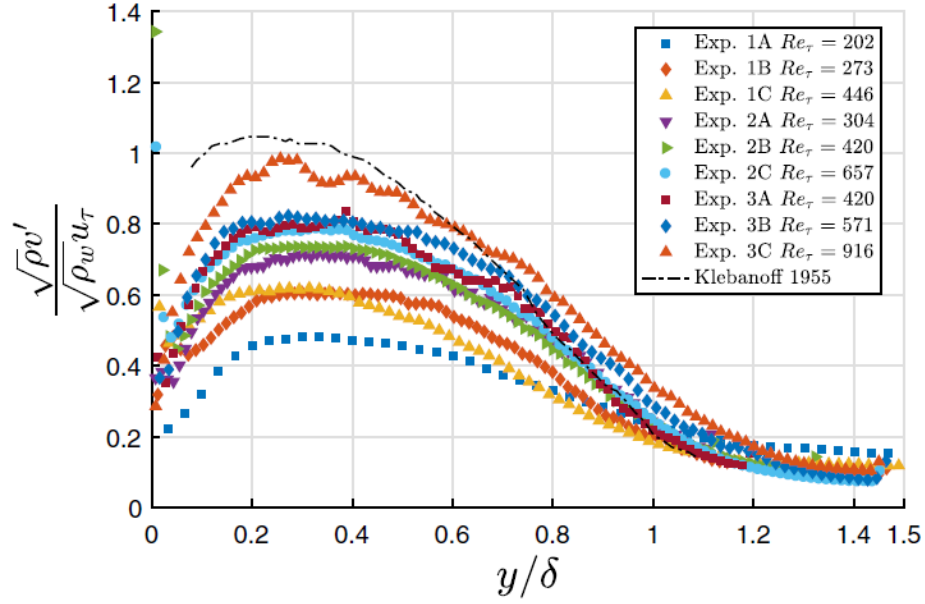


Figure 1.4: PIV Morkovin-scaled wall-normal fluctuations from Brooks et al. [2] in hypersonic conditions, with varying Stokes number and Re_τ compared to the incompressible HWA experiment of Klebanoff [37].

Reprinted from J. M. Brooks, A. K. Gupta, M. S. Smith, and E. C. Marineau, “Particle image velocimetry measurements of Mach 3 turbulent boundary layers at low Reynolds numbers,” *Experiments in Fluids*, vol. 59, p. 83, 2018, with permission from Springer.

ary layers. While the PIV-measured streamwise fluctuations follow trends consistent with both DNS and LDV data, the wall-normal fluctuations are noticeably lower in magnitude.

This behavior is also observed in other PIV data sets obtained under similar flow conditions, as shown in Fig. 1.6. Measurements of streamwise and wall-normal velocity fluctuations from Williams et al. [5], Brooks et al. [2], and Tichenor et al. [7] are compared with DNS and LDV results. Among these studies, the best agreement—although still somewhat underpredicted—occurs in cases with higher friction Reynolds numbers (Re_τ) and lower Stokes numbers. As the friction velocity u_τ increases, the Stokes number decreases, allowing the tracer particles to more faithfully follow the fluid motion.

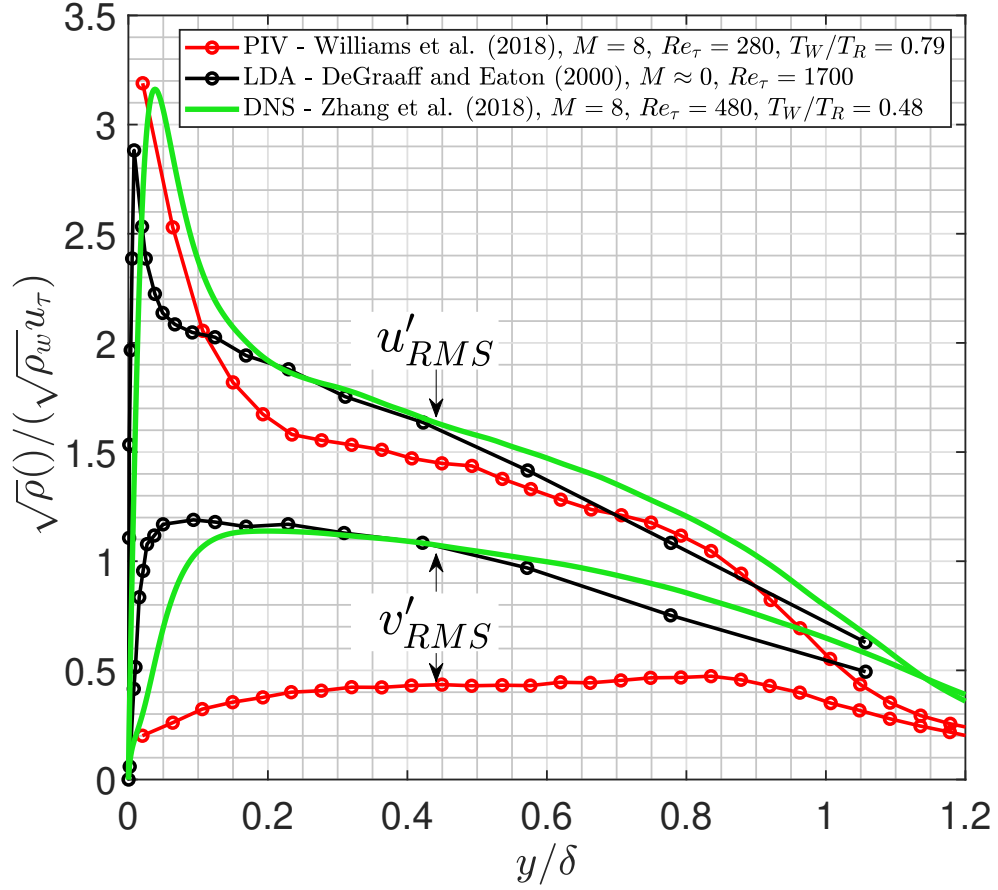


Figure 1.5: Streamwise and wall-normal fluctuations scaled per Morkovin’s hypothesis [3] in compressible, hypersonic turbulent boundary layers from DNS [4] simulation and PIV experiment [5] and compared to incompressible LDA [6] experiment.

The particle response characteristics in hypersonic flows were examined by Williams et al. [35] who analyzed particle motion as tracers passing through shock waves corresponding to 10° and 20° flow at Mach 7.6. The measured particle frequency responses were approximately 36 kHz and 77.2 kHz, respectively, to allow for the particle to faithfully follow the flow dynamics. Particle trajectories were then simulated using the quasi-steady drag formulation of Loth et al. [53] to account for the finite effects of Reynolds, Mach, and Knudsen number on particle drag. These results demonstrated that inertial, compressibility, and rarefaction effects can significantly degrade particle flow imitation in high-speed flows, particularly in the near-wall

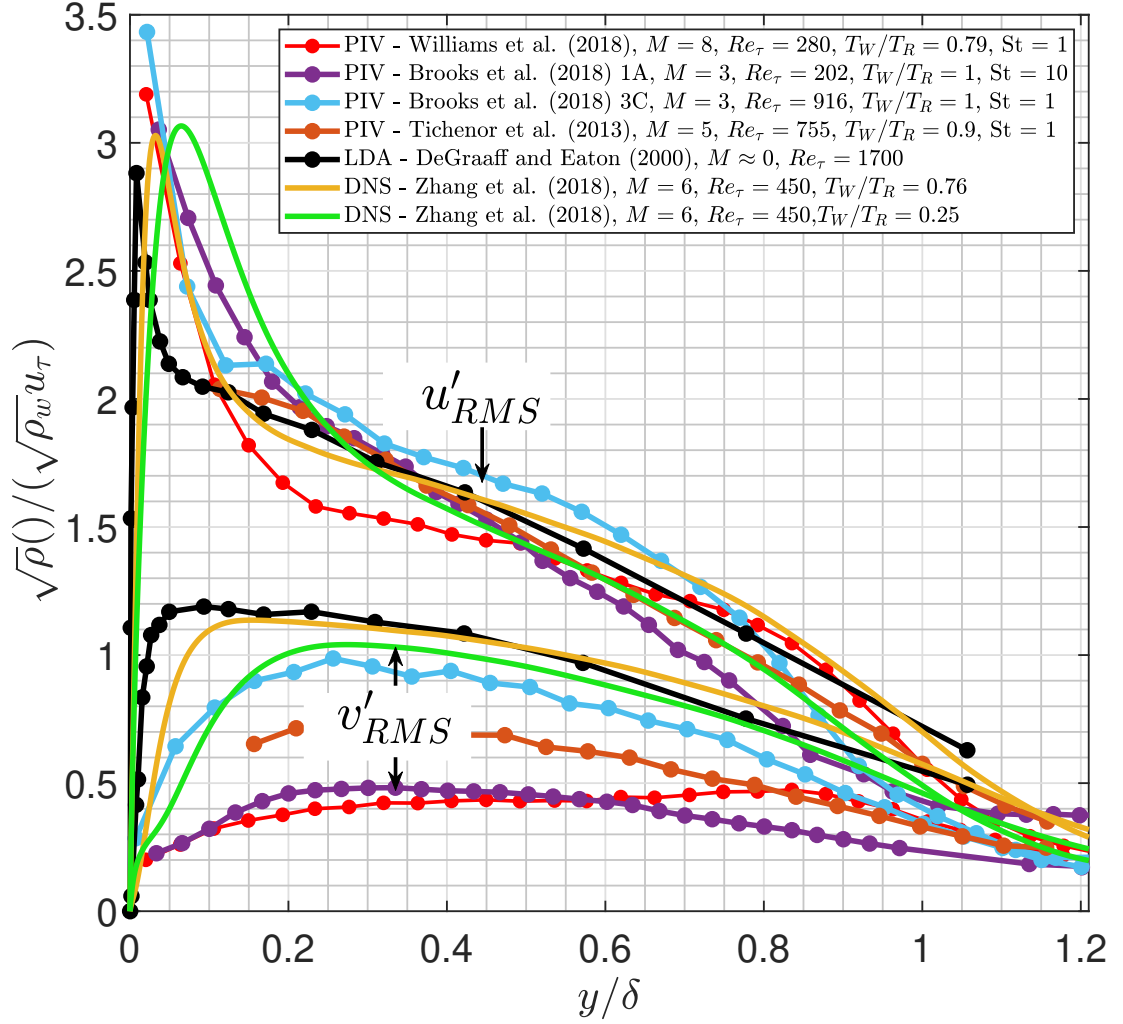


Figure 1.6: Profiles of the density-weighted streamwise and wall-normal velocity fluctuations, $\sqrt{\rho} u'_{RMS} / (\sqrt{\rho_w} u_\tau)$ and $\sqrt{\rho} v'_{RMS} / (\sqrt{\rho_w} u_\tau)$, are shown as a function of the normalized wall-normal coordinate y/δ . Experimental datasets from PIV measurements by Williams et al. [5], Brooks et al. [2], and Tichenor et al. [7], as well as incompressible LDA measurements by DeGraaff and Eaton [6], are compared with direct numerical simulation (DNS) results from Zhang et al. (2018) at two wall-to-recovery temperature ratios. Experimental measurements are indicated by symbols while DNS results are shown as solid lines. The arrows indicate the regions corresponding to the peak streamwise and wall-normal velocity fluctuations.

region where rarefaction becomes important. As a result, tracer particles may not faithfully follow the smallest turbulent motions, contributing to the attenuation of measured wall-normal velocity fluctuations.

Because of these particle-response limitations, researchers have increasingly turned to Molecular Tagging Velocimetry (MTV) to measure velocity fields in high-speed flows without relying on tracer particles [54]. MTV utilizes a molecular tracer seeded into the test gas of a ground-test facility, where it can be optically excited and subsequently emit light, enabling the motion of the tagged molecules to be tracked. Tagging velocimetry serves as a non-intrusive diagnostic technique particularly suitable for the hypersonic flow regime, providing off-surface, time-of-flight velocity profile measurements. By using high-repetition lasers and high-speed cameras, it can interrogate freestream flows and flows over canonical geometries in ground-test facilities. Time-resolved velocimetry data are then used to verify tunnel functionality, quantify turbulence quantities, and describe unsteady boundary layer profiles [55]. Designers of hypersonic vehicles can use tagging velocimetry results to verify computational fluid dynamics (CFD) predictions of flow structures and visualize flow profiles over test articles in regions where CFD encounters difficulty, such as separation bubbles [18].

Popular tracers used for tagging velocimetry include krypton [54–63], oxygen [64], acetone [65, 66], biacetyl [67], hydroxyl [68–71], iodine [72, 73], nitric oxide [74], nitrogen [75–78], and argon [79] as well as many others [80]. Krypton is chemically stable, noncorrosive, and nontoxic. Of the nonradioactive, commercially available noble gases, it has the second lowest ionization energy [81]. Excitation via (2+1) resonance-enhanced-multiphoton-ionization (REMPI) with or without a 769.5 nm CW diode (as measured in air) has been shown to tag krypton atoms. Tagged krypton emits photons from the long-lasting radiative cascade resulting from recombination processes, and re-excitation from metastable state when a second laser is used. A typical setup of a tagging velocimetry setup using a krypton tracer can be seen in Fig. 1.7. In the freestream of the California Institute of Technology (Caltech) T5

piston-driven reflected-shock tunnel, krypton tagging velocimetry (KTV), via (2+1) REMPI at 216.7 nm with a 769.5 nm diode was demonstrated by Shekhtman et al. [82] to resolve freestream velocity measurements. These measurements showed good agreement to DPLR simulated nozzle exit conditions. Making measurements in a high-enthalpy facility, with high background noise, demonstrates KTV's robustness compared to similar methods. Single-laser KTV at 212.6 nm was demonstrated in the Stevens Shock Tube by Mustafa et al. [83]. A successive image capturing technique was used to capture write and read lines over a sectioned, sharp-edged hollow cylinder. The beam was issued tangentially to the cylinder to increase near-wall resolution. Images were mapped to wall-normal values and compared to compressible laminar boundary layer similarity solutions. The generation of the excitation wavelength for KTV requires a tunable laser, such as a dye laser or an optical parametric oscillator (OPO) to picometer precision of the excitation line.

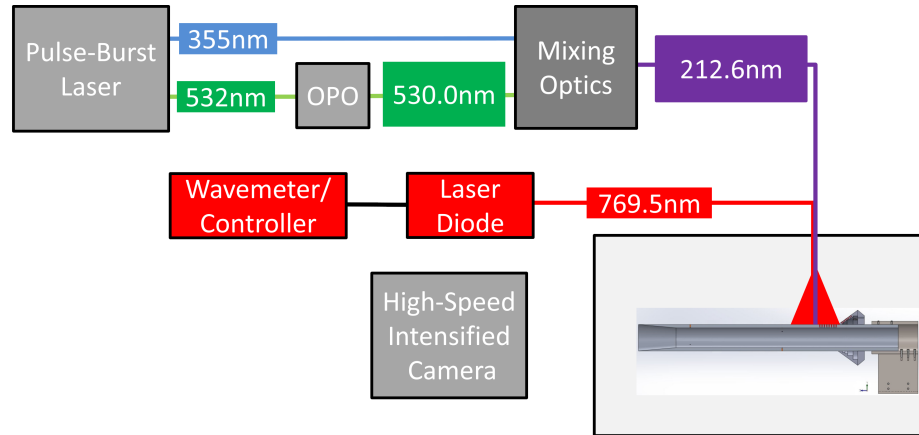


Figure 1.7: KTV experiment schematic

For another tagging tracer, acetone, the excitation spectrum is accessible by the fourth harmonic of an Nd:YAG laser (266 nm) [84, 85]. Acetone tagging at 266 nm was demonstrated by Lempert et al. [65, 66] in an acetone-seeded nitrogen free jet and by Gragston et al. [86] in the turbulent boundary layer and the freestream of

the University of Tennessee Space Institute (UTSI) Mach 4 Ludwig Tube filled with acetone-seeded air. As discussed by Lempert et al. [66], acetone-induced fluorescence has a decay time of $\approx 1 \mu\text{s}$, while phosphorescence has a decay time of the order of $1000 \mu\text{s}$. Lempert et al. [65, 66] observe lifetime of high signal-to-noise ratio (SNR) acetone fluorescence for 500 ns, which justifies their interframe time delay of 230 ns. Tran et al. [87] report higher lifetimes in liquid acetone if the dissolved oxygen in a liquid acetone solution is purged prior to use, resulting in signal lasting up to $4 \mu\text{s}$ in acetone-seeded nitrogen gas. Gragston et al. [86] report lifetimes on the order of microseconds for acetone seeded in oxygen. Fratantonio et al. [85] state that the optimal excitation wavelength for fluorescence in acetone is at 308 nm with a lifetime of $50 \mu\text{s}$. Gragston and Smith [86] report the existence of significant, measurable fluorescence signal at 266 nm. While acetone is 30% less massive than krypton, acetone will pyrolyze at temperatures exceeding 573 K [88]. Pyrolysis of the acetone reduces the fluorescence signal. According to Capelin et al. [89], acetone will pyrolyze noticeably at 650 K and completely by 850 K. Therefore, acetone can be used for shock tunnel experiments for which the nozzle reservoir does not exceed 573 K or if the timescale of acetone pyrolysis is greater than that of the tunnel test time, which may well be the case for 1% acetone/99% air gas mixture [90].

1.6 Scope and Outline

Despite extensive numerical simulations and experimental investigations, reliable measurements of turbulence statistics in hypersonic boundary layers remain limited, particularly for wall-normal velocity fluctuations. To further the understanding of hypersonic turbulent boundary layers, and to provide experimental measurements in support of Morkovin’s hypothesis, non-intrusive flow diagnostics were performed on a

Mach 6 turbulent boundary-layer flow over a hollow-cylinder flare (HCF). ATV was performed in the streamwise direction to obtain both mean flow and fluctuations, u and u' respectively. KTV was performed in the streamwise direction to obtain both mean flow and fluctuations, u and u' as well. A new technique for obtaining wall-normal velocity fluctuations, v' , was also used. By utilizing a single laser line oriented at a small angle over the test article, fluctuations were measured using perturbations from an edge fit line. The run conditions for the facility were: an “enthalpy-matched” condition to match flight conditions at Mach 6 and 20 km altitude with $Re_\tau = 780$ and $T_w/T_r = 0.17$ used for streamwise and wall-normal KTV; a “cold-flow” condition to increase Reynolds number and test time with $Re_\tau = 408$ and $T_w/T_r = 0.65$ used for streamwise ATV and $Re_\tau = 450$ and $T_w/T_r = 0.54$ for both streamwise and wall-normal KTV. Both the streamwise and wall-normal fluctuation data were non-dimensionalized using Morkovin scaling and compared with the literature. The analysis shows that the streamwise experimental values agree favorably with the theoretical and DNS values in the literature. Additionally, the wall-normal fluctuation values match the DNS trends, unlike alternative velocimetry techniques.

1.7 Research Approach

The following chapters provide a detailed report on the experimental setup including: the facility (Chapter 2) and optical setup (Chapter 3); streamwise tagging using acetone (Chapter 4) and krypton (Chapter 5) tracers; and wall-normal fluctuation tagging using krypton tracers (Chapter 6); turbulent statistics and discussion with data measured during the experiment campaigns (Chapter 7). Conclusions and future work can be found in Chapter 8

Chapter 2

Experimental Facility: The Stevens Shock Tunnel

This chapter describes the Stevens Shock Tunnel (SST) facility used to perform the experiments presented in this dissertation. The SST is a reflected-shock tunnel designed to generate short-duration hypersonic flows representative of Mach 6 free-flight conditions. The run conditions for the facility were: an enthalpy-matched condition to match flight conditions at Mach 6 and 20 km altitude with $Re_\tau = 780$ and $T_w/T_r = 0.17$ and a cold-flow condition to increase Reynolds number and test time with $Re_\tau = 450$ and $T_w/T_r = 0.54$. This chapter summarizes the facility design, operational parameters, and the modifications implemented during the commissioning and shakedown phases that enabled the turbulence measurements described in later chapters.

2.1 Introduction to the Stevens Shock Tunnel

The Stevens Shock Tunnel is a facility designed to replicate Mach 6 free-flight flow conditions with an enthalpy of 1.5 MJ/kg and a unit Reynolds number $0.35 - 8.1 \times 10^6 \text{ m}^{-1}$ for at least 4 ms using helium as the driver gas. In addition, a cold-flow condition uses an N_2 driver, providing a longer 15 ms test time and an increased Reynolds number of up to $32 \times 10^6 \text{ 1/m}$. This reduces the reservoir enthalpy but provides sufficient heating such that the flow in the freestream does not condense. As determined by the standards and by hydrotesting, the maximum allowable pressure of the straight portions of the driver and driven sections is 16.6 MPa (2400 psig) and the maximum allowable pressure of the elbow is 10.4 MPa (1500 psig). In Fig. 2.1, the current configuration of the tunnel is displayed. The driver section is 5.0 m long,

and the driven section is 11.07 m long. Five transducer stations are used to provide a shock-speed measurement. Pressure transducers are flush with the pipe. The nozzle reservoir pressure is measured 1 inch (0.0254 m) from the shock tube end. The nozzle exit diameter is 16 inches (406 mm) and the test rhombus is approximately 13 inches (330 mm) high by 5 ft (1.5 m) long. The dump tanks can hold a volume of 8.78 m³ (2320 gal) at a vacuum pressure of 350 mtorr and have a maximum rated pressure of 40 psig. More structural details of the facility can be found in Shekhtman et al. [91].

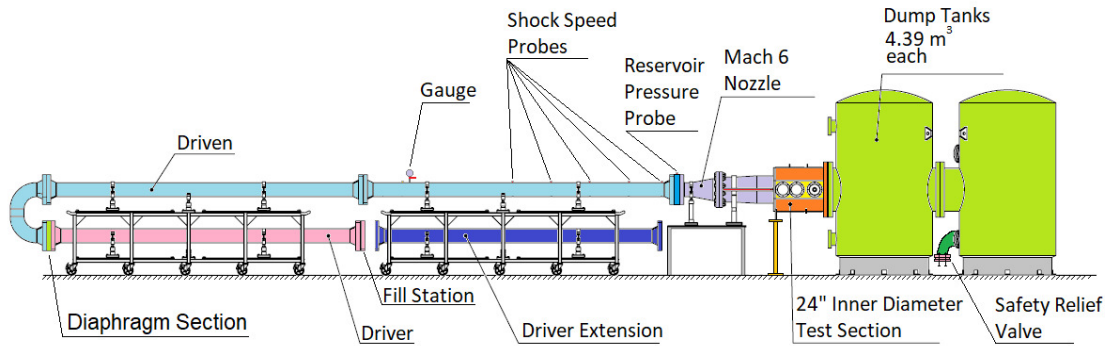


Figure 2.1: Schematic of Stevens Shock Tunnel (SST).

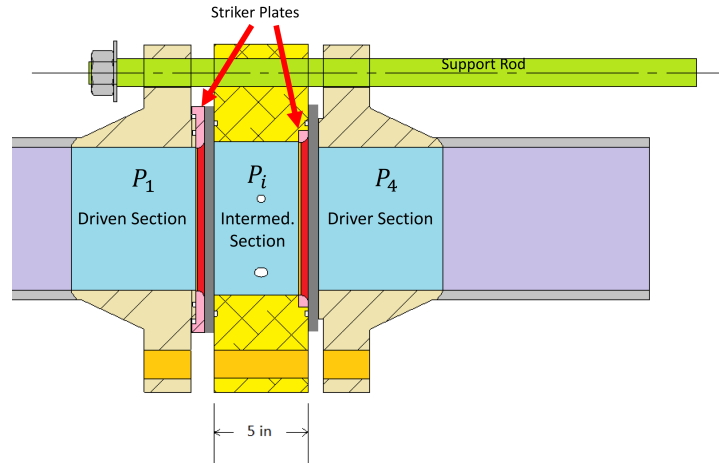


Figure 2.2: Double Diaphragm Section in the Bottom Diaphragm Station, before the Elbow, as depicted in Fig. 2.1.

Table 2.1: Max Pressure Differential vs. Diaphragm Thickness for Temper O aluminum. For all diaphragms, pressure acts on a diameter of 194 mm (7.625 in).

Grade	Thickness (in)	ΔP (psi)	Scored 0.010~in, ΔP (psi)	Stamped Pressure (psi)	Stamped ΔP (psi)
1100-O	0.020	103	N/A	N/A	N/A
1100-O	0.032	171	N/A	N/A	N/A
1100-O	0.050	280	N/A	500	140
1100-O	0.090	495	377	1000	140
1100-O	0.125	665	550	1000	140
6061-O	0.125	956	670	1025	140
1100-O	0.190	1045	750	N/A	N/A

2.1.1 Firing Mechanisms and Diaphragm Material

The pressure rupture of either a single or double diaphragm initiates a run in the Stevens Shock Tunnel. Single-diaphragm rupture is used for prototyping runs. As shown in Fig. 2.2, a double diaphragm mechanism [92] is used to conduct tests at specified pressures and pressure ratios.

Typical diaphragms consist of standard square 12 in $\times$ 12 in 1100-O aluminum of different thicknesses ranging from 0.010 to 0.190 inches. Burst pressures for such diaphragms with fixed-end boundary conditions are listed in Table 2.1. We observed that soft Temper O aluminum and cold-rolled steel were more effective in petalling than harder materials, which tended to shatter during rupture. Clamping of the diaphragm is achieved by clamping the section between the two flanges with twelve 1-3/8-6 Grade 5 bolts torqued to 340-420 lbf·ft and concentric serrations on the Class 900 Flanges that prevent the slippage of diaphragm material relative to the flange. Fixed ends constrain a diaphragm to thin out at the center instead of the ends. Flange serrations are critical to this effect. According to Gaydon et al. [92, Ch. 5], friction generated by bolt torque is generally insufficient as a sole means to enforce the fixed-ends constraint. A striker plate, as shown in Fig. 2.2, is used to form a bend radius at

the edge of each diaphragm as it bulges into a dome via applied pressure. Diaphragm scoring is used to encourage petalling of diaphragms and to improve pressure recovery P_R/P_4 . Scoring to a depth of 0.010 in is accomplished via CNC machining.

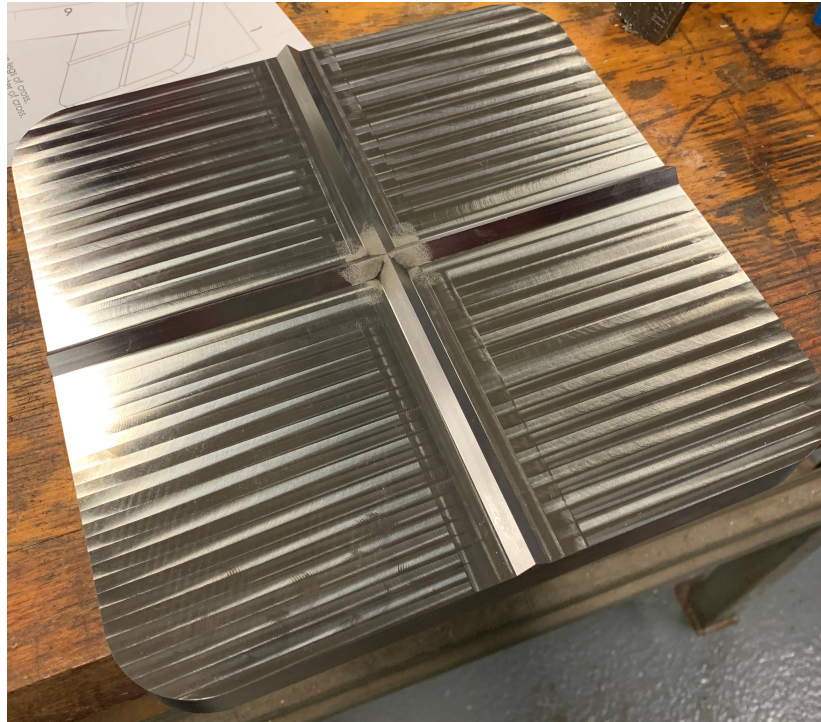
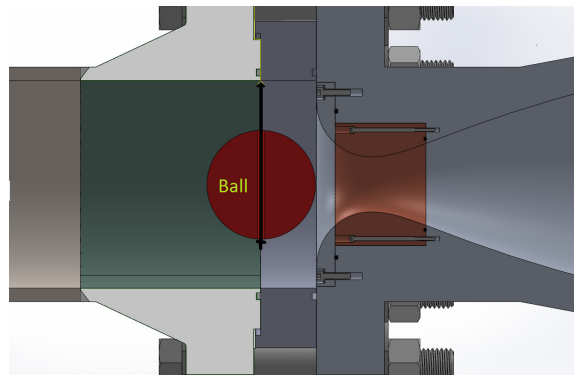


Figure 2.3: Diaphragm stamper used for all diaphragm thicknesses.

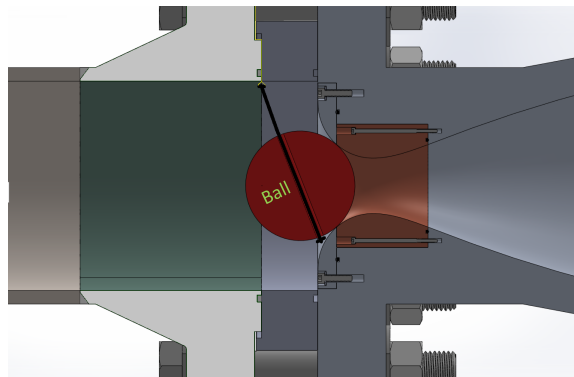
The diaphragm stamper was developed to replace the CNC scoring process. In addition to being time consuming, CNC scoring melted material into the ridges of the diaphragms, which reduced repeatability. Small pieces of material lodged in the ridges also introduced noise during runs. The diaphragm stamper was designed to interface with an in-house hydraulic press and to stamp diaphragms to prescribed burst pressures. An adapter chuck was fitted to the press cylinder and housed magnets to hold the stamp in place. The stamp consisted of a 1 in base with four sharp blades extending radially from the center to the edge of the plate over a span of 4.25 in. A2 tool steel was selected for the assembly to maintain blade sharpness and improve repeatability. An image of the stamp is shown in Fig. 2.3. Based on burst-test trials,

the stamping pattern was selected as a “plus” orientation for 1100-O diaphragms, with the grain direction aligned vertically, and as a “cross” orientation for 6061-O sheets with the grain direction aligned vertically.

2.2 Ball Centerbody Design



(a) Ball Start Position



(b) Ball Final Position

Figure 2.4: Schematic of Ball Plugging the Nozzle Converging Section. The ball moves a distance of 35.8 mm (1.41 in) from positions (a) to (b).

Premature arrival of driver gas in the reservoir region would manifest as large oscillations in Pitot-pressure traces, a bump in the center of the Pitot-pressure profile, and potentially, altering the test-gas composition. The causes of this early, nonuniform reservoir contamination likely include bifurcation. According to Davies [93] and

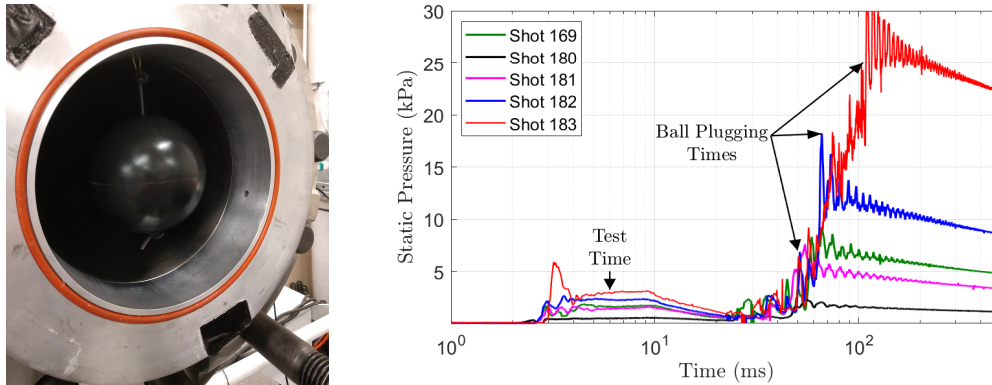


Figure 2.5: Ball Centerbody. *Left:* Picture of Ball Centerbody. *Right:* Static Pressure Traces for Multiple Shots Showing Test Intervals and Ball-Plugging Times.

Davies et al. [94], when helium gas drives nitrogen test gas, bifurcation occurs when the incident shock Mach number is greater than 3.5, which is near the run condition for generating the tunnel's design enthalpy. Wall-jetting due to bifurcation forms a finger-like structure that contaminates the centerline of the nozzle exit flow [95, 96]. Following the strategy of Holden et al. and CUBRC [97, 98] for the CUBRC Shock Tunnel, a simplified implementation of a centerbody apparatus was installed in the reservoir section of the Stevens Shock Tunnel to terminate nozzle flow once the useful test gas slug has been depleted. As shown in Fig. 2.4, a spherical neoprene ball (referred to as the ball centerbody) was placed 16 mm (0.63 inches) from the reservoir region. During a shock tunnel run, driven and driver gas push the ball centerbody approximately 36 mm (1.415 in) towards the nozzle converging section.

The ball centerbody fulfills three major functions. It (1) prevents diaphragm debris from entering the nozzle, (2) ends the test time, and (3) prevents dump-tank overpressurization. If the ball centerbody properly seals during a run, nozzle recoil and structural loading are significantly reduced; and potentially, driver gas remaining in the tube can be piped out and reused. The ball centerbody design proved capable of reliably sealing the nozzle throat and ending the test time, as seen in Fig. 2.5.

When the ball centerbody is used in conjunction with the diaphragm station placed upstream below before the elbow, pressures traces like those obtained during Shot 169 become routine in the SST, as seen in Fig. 2.6. Compared with Shot 179, which was conducted under the same run conditions, Shot 169 does not exhibit the drop in static pressure or the Pitot-pressure oscillations observed at 7 ms, effectively increasing the usable test time by approximately 1 ms.

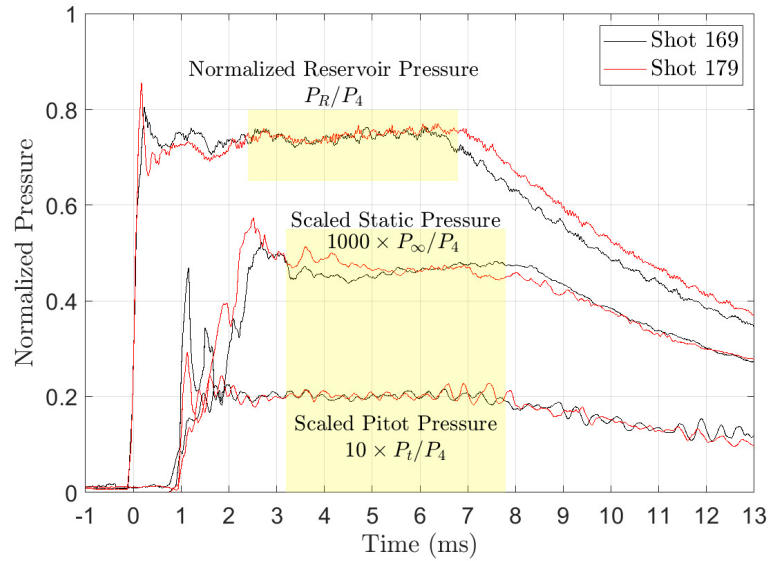


Figure 2.6: Shock Tunnel Pressure Traces for Shot 169 (with ball centerbody) and Shot 179 (without ball centerbody). Traces include reservoir pressure, freestream static pressure, and freestream Pitot pressure normalized by driver pressure P_4 and scaled by the factors of 1, 1000, and 10, respectively. Test time is denoted by yellow highlighted regions.

2.3 Revised Centerbody Design

A nozzle centerbody was installed in the nozzle reservoir region of the Stevens Shock Tunnel. The centerbody is shown in Fig. 2.7 with individual parts labeled in Fig. 2.8. Unlike the previous iteration of the centerbody which consisted of a suspended neoprene ball that swung into the nozzle converging section, the new centerbody is con-



Figure 2.7: Centerbody in the Stevens Shock Tunnel.

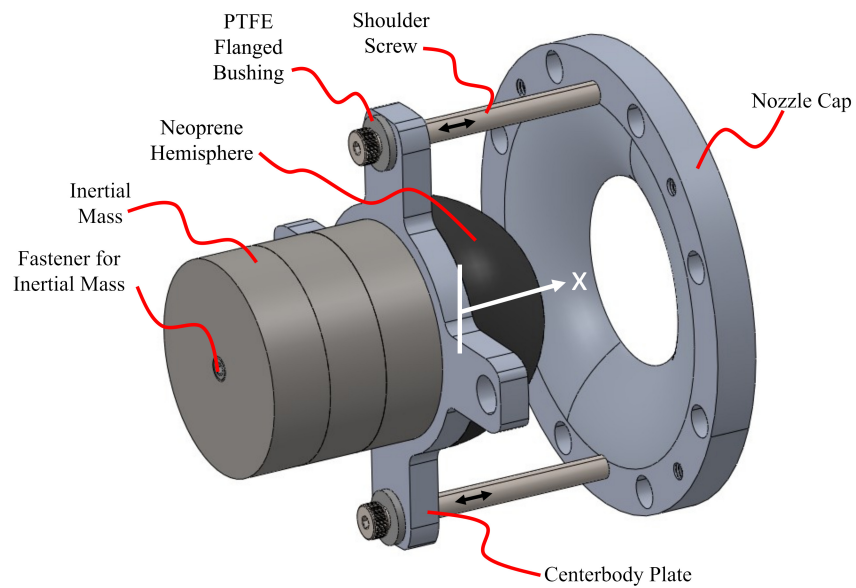


Figure 2.8: Centerbody assembly.

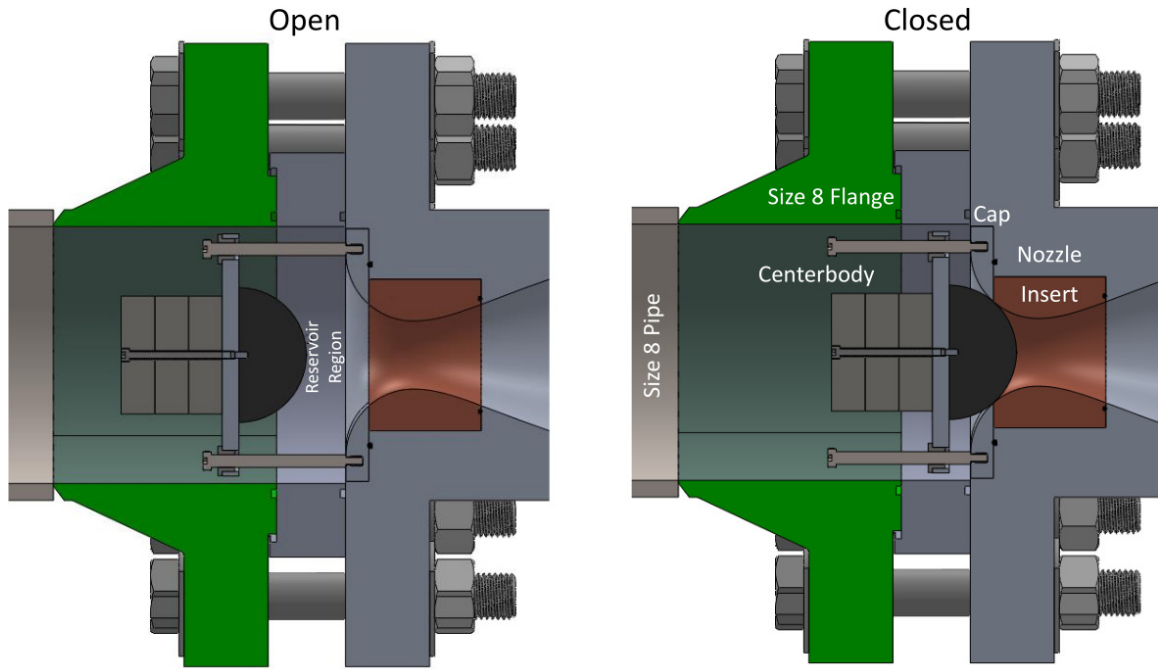


Figure 2.9: Side View of Centerbody Plugging the Nozzle.

strained to slide linearly toward the nozzle throat. As shown in Fig. 2.8, the centerbody assembly consists of a solid neoprene hemisphere mounted onto an aluminum plate with PTFE flanged bushings that slide along two 3/8-in diameter shoulder screws. The shoulder screws are screwed into a nozzle throat cap which clamps the secondary diaphragm in place. The neoprene half-ball functions as a plug for the nozzle. In the rear of the aluminum plate, an adjustable inertial mass is used to tune the closing time of the centerbody. The new design offers several advantages: (1) the shoulder screw design ensures the centerbody centered throughout the run; (2) mounting the centerbody to the nozzle throat cap enables quick turnaround time; (3) centerbody sealing time is adjustable and tunable; (4) the concept is more structurally robust; and (5) the neoprene hemisphere is less prone to bouncing off the nozzle surface due to its increased inertia. In Fig. 2.7, the centerbody prototype is shown in its closed state in the nozzle reservoir region of the Stevens Shock Tunnel.

Side-view schematics of the open and closed states of the sliding-plug centerbody are presented in Fig. 2.9. During an experiment, the centerbody obtains initial momentum relative to the shock tube from both the facility recoil and the incident shock wave to translate downstream towards the nozzle converging section. The inertia of the centerbody is set such that the centerbody reaches the surface of the nozzle after the useful test time of the facility. To estimate the time required for the centerbody to seal the nozzle throat, a simple momentum-based model was developed based on the impulse imparted by the incident shock. Using Fig. 2.10, the closing time of the centerbody is estimated to be

$$t_{close} = \frac{\left[\Delta x - u_t \left(\frac{L_{Driven}}{u_s} \right) \right]}{\left(\frac{P_2 A}{u_s} \right) \left(\frac{r_{ball} + t_{plate} + L_{mass}}{m_{ball} + m_{plate} + \rho_{steel} A_{mass} L_{mass}} \right)}. \quad (2.3.1)$$

The numerator of Eq. 2.3.1 represents the remaining stroke ($< \Delta x$) of the centerbody after the tunnel has recoiled at velocity $u_t = (m_{Driven,gas} + m_{Driver,gas})u_2/m_{tunnel,pipes}$ during the time it takes for the incident shock to arrive L_{Driven}/u_s at the centerbody from the diaphragm station. The denominator of Eq. 2.3.1 is the velocity of the centerbody resulting from momentum transfer from the incident shock passing over the body. Gas processed by the incident shock moving at speed u_s results in pressure P_2 and velocity u_2 relative to the tunnel. The front-facing projected area of the centerbody is A . The residence time of the shock wave on the centerbody is estimated as $(r_{ball} + t_{plate} + L_{mass})/(u_s)$. The impulse from the shock therefore calculated as $P_2 A (r_{ball} + t_{plate} + L_{mass})/(u_s)$. Dividing the impulse by the mass of the centerbody, the velocity of the centerbody is obtained. Eq. 2.3.1 can be used to estimate the length of required inertial mass L_{mass} . Once at the surface, the centerbody forms a plug seal, terminating the nozzle flow. For shots 300, 303, and 306, the mass of the

centerbody is 5 kg, and the estimated plugging time is ≈ 50 ms.

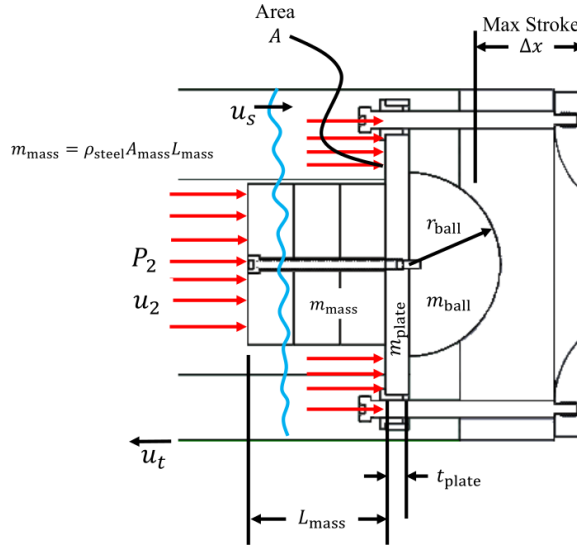


Figure 2.10: Sketch of Nozzle Centerbody for Plugging Time Model.

2.3.1 Centerbody Results and Discussion

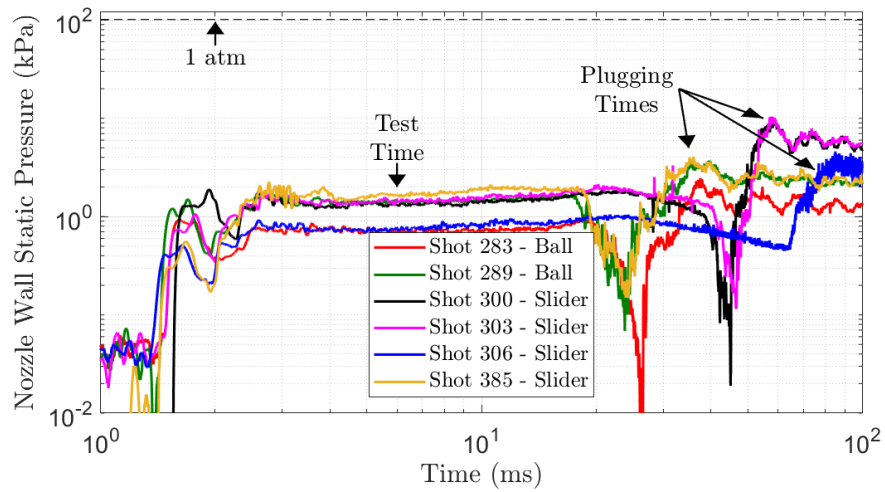


Figure 2.11: Nozzle wall static pressure traces showing plugging times for N_2 -driven shots in the Stevens Shock Tunnel using a sliding nozzle centerbody.

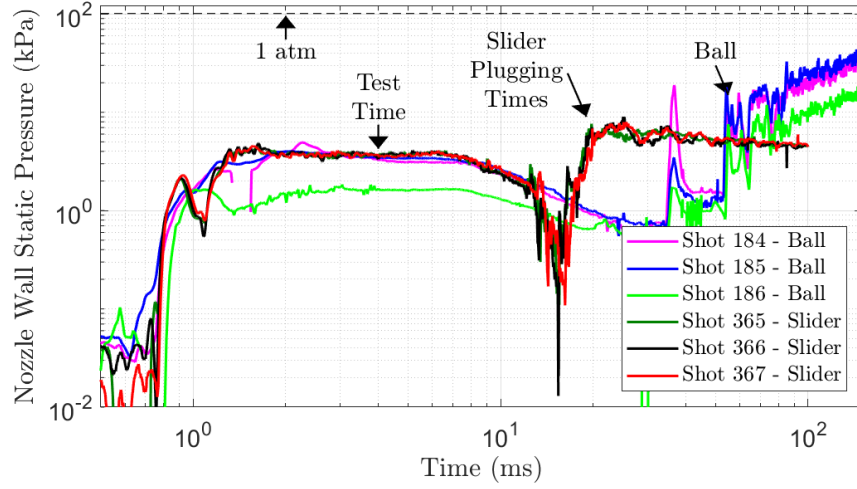


Figure 2.12: Nozzle wall static pressure traces showing plugging times for He-driven shots in the Stevens Shock Tunnel using a sliding nozzle centerbody.

Pressure traces from testing performed in the Stevens Shock Tunnel are presented to compare the performance of the previous and updated centerbody designs. As shown in Figs. 2.11 and 2.12, the nozzle wall static pressure traces of various experiments depict the usable test time and plugging times of the nozzle centerbody. Smooth traces in the test time indicate clean flow which both centerbody designs achieve consistently. The centerbody plugging time is indicated by the rising edge of the right-hand pressure plateau. In Fig. 2.11, it is observed that the ball centerbody closes at approximately the same time for every run when using N_2 as a driver gas. This phenomenon is concerning because shots with higher driver pressure and the same pressure ratio will increase the pressure acting on the ball centerbody, which should decrease the ball plugging time. We are currently only able to speculate on the performance of the ball centerbody. We believe that there may be cases where the ball centerbody does not immediately seat itself in the nozzle throat.

The updated sliding centerbody translates as a result of an impulse from the incident shock. Its rear inertial mass allows for customization of its mass to control

the plugging time at the nozzle throat. In shots 300, 303, and 306 with the sliding centerbody, the mass of the centerbody was ≈ 5 kg. Shots 300 and 303 have a higher driver pressure than shot 306. The lower pressure shot, which featured a decreased driven slug pressure, took longer for the slider to plug the nozzle throat, which was expected due to decreased impulse acting on the body. An analysis was performed to obtain the weight of the inertial mass to emulate the plugging time of the original ball centerbody for the nitrogen driver gas condition. After testing various weights, it was determined that a mass of approximately 2 kg yielded almost identical run results between the two centerbodies with similar run conditions. Shots 289 and 385 in Fig. 2.11 show the successful replication of the plugging behavior. For helium driver experiments, the mass of the centerbody was decreased to 1 kg to have a plugging time of ≈ 20 ms as shown in Fig. 2.12. The sliding centerbody plugged without bouncing off the nozzle throat and experienced consistent plugging behavior for each experiment run.

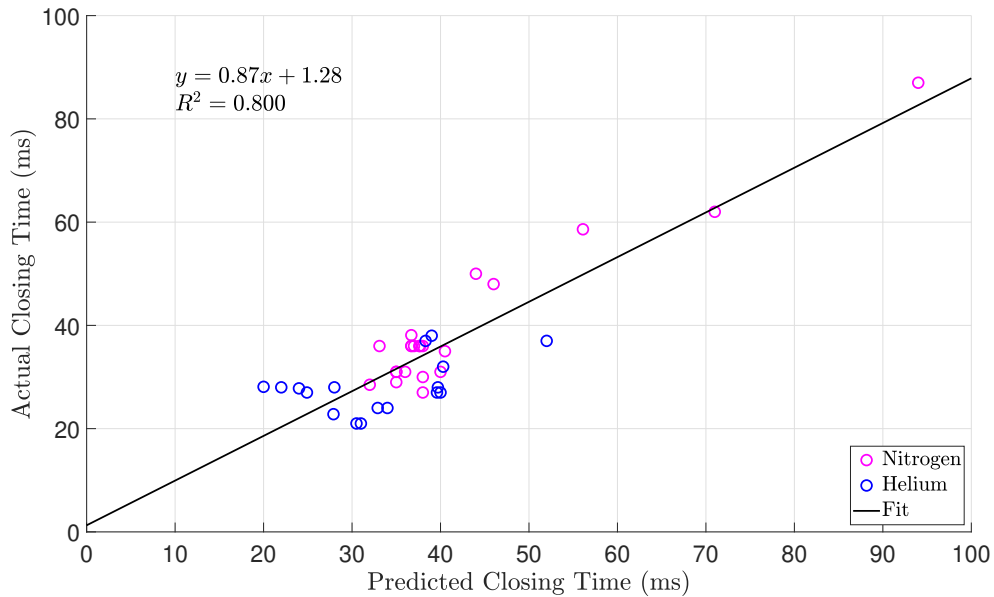


Figure 2.13: Closing time comparisons between Eq. 2.3.1 and experimental closing time. Combined line of best fit plotted with equation and R^2 value.

Plugging times for both the N₂ and He shots were determined using the nozzle wall static pressure traces shown in Figs. 2.11 and 2.12. These values were compared to the calculated closing time values using Eq. 2.3.1. This is shown in Fig. 2.13. A linear regression was performed on the data to validate the linearity between the calculated and measured closing times. The predicted closing time shows good correlation with the measured closing time, with a linear fit yielding $R^2 = 0.80$. This indicates that the equation can give a good closing time estimate. However, the fitted slope of 0.87 suggests a systematic over-prediction at longer closing times, while the clustered short-time cases show appreciable shot-to-shot scatter. Therefore, the model is useful for estimating closure-time trends, but individual predictions should be treated as approximate rather than exact.

2.4 Effect of Tunnel Cleanliness

The discussion on the location of the diaphragm station and the use of the ball centerbody as a debris blocker extends into shock tunnel cleanliness. Diaphragm ruptures and dump-tank dust clouds deposit material throughout the shock tunnel. Cleaning is essential to protect the integrity of the contact surface during a test and to reduce the amount of debris passing past the ball centerbody into the test section. Using the motorized cart presented in Fig. 2.14 to push out large debris and lay a line through the driven section in Fig. 2.1, a large foam bucket with microfiber cloths is pulled through, collecting dust and small debris. In Fig. 2.15, the effect of tunnel cleanliness manifests as the amount of diaphragm debris prevented from flying into the test section. Shot 128 is a control shot with no implemented cleaning strategy. Shot 179 is a shot that uses driven tube cleaning, without a nozzle centerbody (the ball centerbody). Shot 157 shows that the ball centerbody and tunnel cleaning are effective

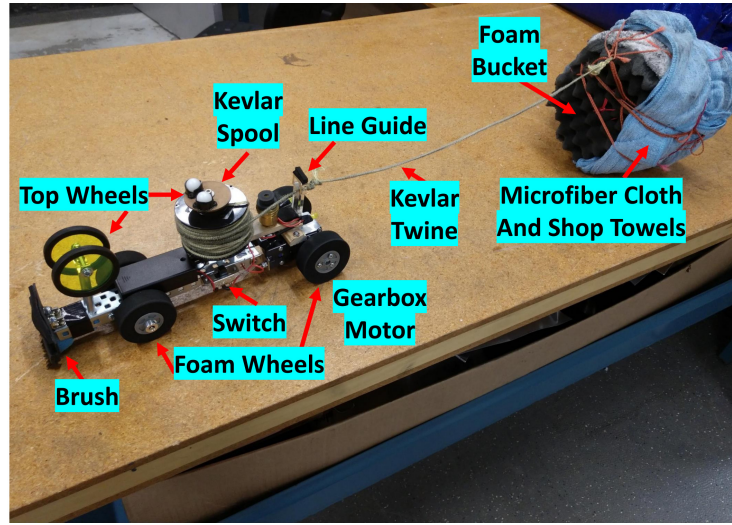


Figure 2.14: Motorized Twine-laying Cart and Microfiber Cloth bucket for Stevens Shock Tunnel.

in keeping the test section clean during the test time. It is noted that schlieren was unable to indicate the presence of small particulate in the test section. Images were obtained via a Phantom 2512 at 20 kHz framerate aimed at the test section window with a Nikon NIKKOR 24–85-mm f/2.8-4D. Maintaining tunnel cleanliness preserves the surface integrity of test articles and prevents premature transition to turbulence on test articles by diaphragm debris. This adds further sustainability to the test article surface (including coatings such as pressure-sensitive and temperature-sensitive paints [99]) and contributes to the integrity of turbulence experiments, as emphasized by the work of Jewell et al. [100].

2.5 Conclusions

Steady Pitot and static pressure traces are presented for experiments over a range of unit Reynolds numbers ($0.35\text{--}8.1 \times 10^6$ 1/m) at Mach 5.8–6.0 and enthalpy 1.54 MJ/kg. Diaphragm debris was prevented from entering the test section via the nozzle centerbody. Tunnel cleanliness also played an important role in configuring the Stevens

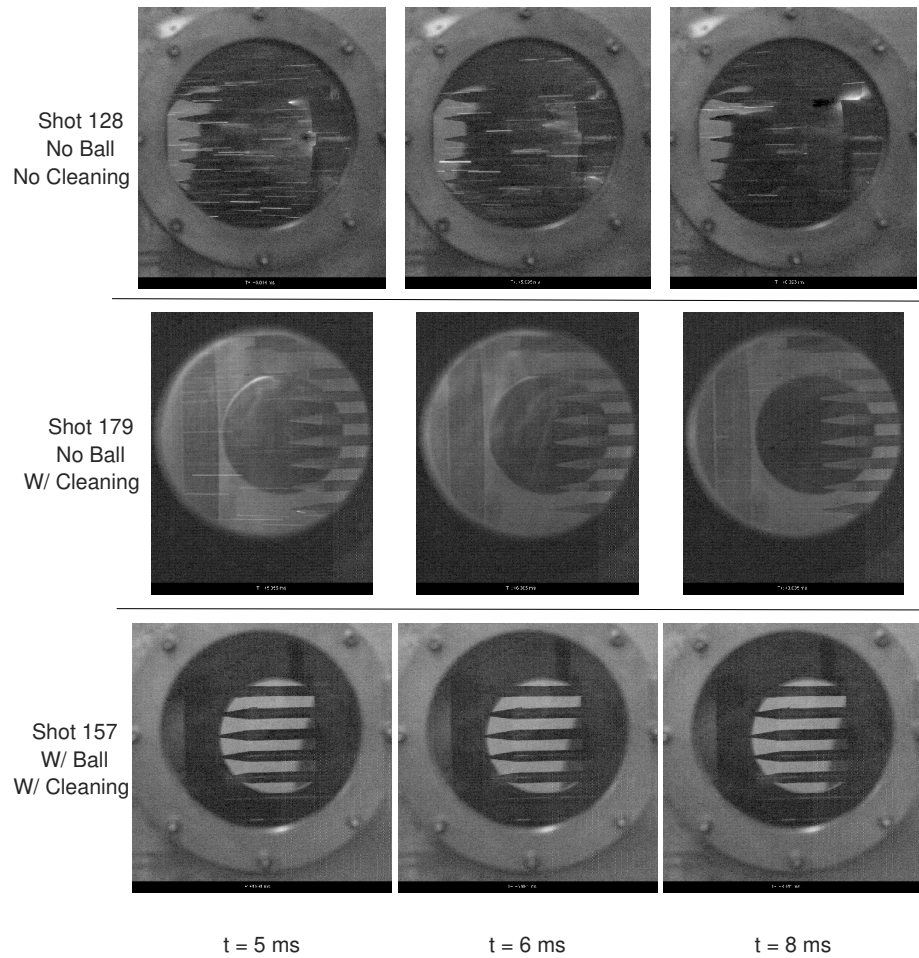


Figure 2.15: Effect of Tunnel Cleaning Procedures and Centerbody Apparatus on Debris Entering the Test Section.

Shock Tunnel for transition and turbulence experiments. The high-speed nozzle-flow termination has reduced facility recoil at the end of an experiment and prevents test-section and dump-tank over-pressurization. Plugging the flow following the useful test time creates an opportunity to reclaim the driver gas, which in the present experiments was helium. A cost/benefit analysis will be performed to assess the viability of the driver-gas reclamation process. We observed that the Pitot- and static-pressure traces were consistently steadier for a longer period of time when a centerbody was used. Using the t-test method, we determined there to be a statistically significant,

12% reduction in the Pitot-pressure fluctuations with the addition of the centerbody. We speculate that this could be an indication that the centerbody reduces driver-gas contamination due to bifurcation or test-gas slug depletion, or, alternatively, that the geometry change in the reservoir altered the propagation of noise in the facility. Investigation of the fundamental mechanisms behind this behavior is a topic of future work. The ball centerbody has performed well for approximately 75 shots in the SST with no observed wear at reservoir conditions typically above $T_R = 1600$ K; however, other researchers may require a centerbody that could withstand more demanding test conditions.

Chapter 3

Optical Setup

The measurements presented in this dissertation rely on high-repetition-rate laser diagnostics capable of resolving turbulent velocity fluctuations within the short test times of impulse facilities. This chapter describes the optical system developed to perform krypton tagging velocimetry (KTV) in the Stevens Shock Tunnel. The system consists of a burst-mode Nd:YAG laser, an optical parametric oscillator (OPO), a high-speed camera, and an intensified relay optic. A representative design of the complete system is seen in Fig. 3.1. The design considerations and operating principles of each component are described in the following sections.

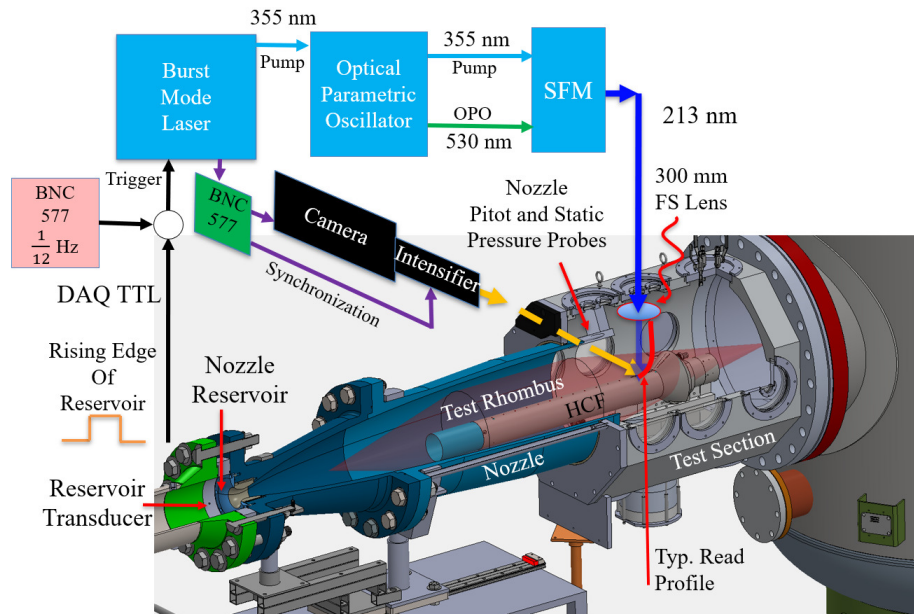


Figure 3.1: KTV Experiment Schematic

3.1 Burst Mode Laser

To obtain a statistically significant number of realizations during the short test time of the SST facility, a pulse-burst laser with a high-repetition rate was required. This generation of burst-mode laser systems by Spectral Energies was used concurrently in high-repetition rate laser diagnostic work including planar laser-induced fluorescence (PLIF) [101–104] as well as molecular tagging velocimetry with krypton and other tracers such as nitric oxide [8, 105, 106].

The QuasiModo 1200 utilizes seven Nd:YAG flashlamp amplifiers to create fundamental (1064.46 nm) beams at high power and high repetition rate. The laser system includes three Lithium Triborate (LBO) crystals capable of performing second harmonic generation (SHG), third harmonic generation (THG) and fourth harmonic generation (FHG) creating the harmonic Nd:YAG beams 532 nm, 355 nm and 266 nm respectively. The laser system has three seed laser modes: continuous wave (cw), nanosecond pulsed diode (ns) and giant mode in which the pulse width can be increased from the order of 10 ns to 1000 ns. The cw seed mode was used for all experiments in this work since the mode features much more narrowband linewidth relative to ns mode.

3.2 Optical Parametric Oscillator

Following the work of [54, 107], the krypton excitation line of 212.56 nm was chosen. Ideally, a burst mode laser would be able to reach that wavelength independently by means of a fifth harmonic, but the output wavelength is 212.89 nm which is too far away from the krypton excitation line. Therefore, an OPO was introduced to the system to finely-tune the output laser wavelength to the desired krypton excitation line. The OPO was originally documented in Grib et al. [8, 101]. A diagram of the

OPO system is shown in Fig. 3.2.

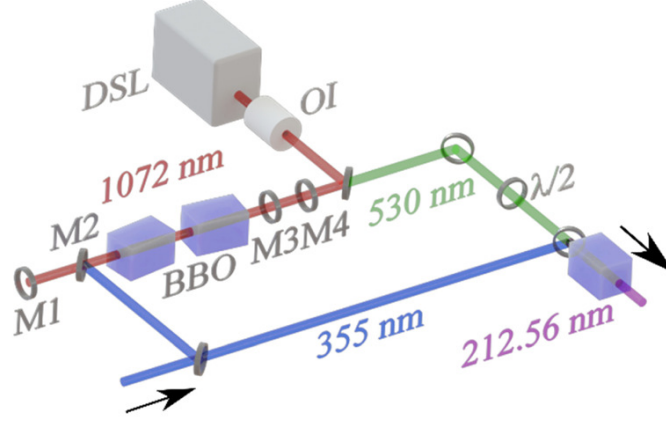


Figure 3.2: Diagram of OPO design from Grib et al. [8] to produce 212.56 nm light. A beamsplitter was used to split the 354.82 nm output and direct a portion into the OPO, while the remainder was used for mixing. M1 and M3, OPO cavity mirrors; M2, 45° turning mirror; M4, 354.82 nm double-pass reflector; $\lambda/2$, half-wave plate; DSL, diode seed laser; OI, optical isolator.

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In an OPO, energy conservation requires that the pump wavelength be divided into signal and idler wavelengths [108]:

$$\lambda_{signal} = \frac{1}{\frac{1}{\lambda_{KTV}} - \frac{1}{\lambda_{pump}}}, \quad (3.2.1)$$

in which 354.82 nm is the pump wavelength λ_{pump} and 212.56 nm is the desired KTV excitation wavelength λ_{KTV} . With knowledge of the signal wavelength, the idler wavelength of the OPO cavity can be determined

$$\lambda_{pump} = \frac{1}{\frac{1}{\lambda_{signal}} + \frac{1}{\lambda_{idler}}}. \quad (3.2.2)$$

Using the above equations, the OPO signal wavelength λ_{signal} and the idler wavelength λ_{idler} are determined to be 530.13 nm and 1072.93 nm respectively. During

an experiment, the λ_{idler} is actively managed.

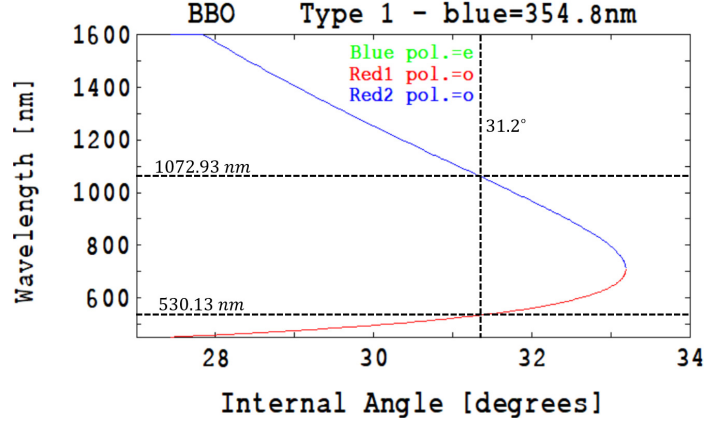


Figure 3.3: Output from the SNLO `Opoangles` widget used to determine the BBO cut angle for OPO phase matching. The pump wavelength, 354.82 nm, is *e*-polarized, while the signal and idler wavelengths, 530.13 nm and 1072.93 nm, are both *o*-polarized.

The OPO is pumped by the third harmonic of the burst-mode laser at 354.82 nm. A beam splitter was used to direct 66% of the power into the OPO cavity. The OPO optical cavity consists of two Type I Beta Barium Borate (BBO) crystals cut at 31.2° relative to the optical axis. The cut angle was determined using the SNLO "Opoangles" widget [109]. A plot of the internal angle and wavelength shown in Fig. 3.3 illustrates the overlap at the desired signal and idler wavelengths for the pump input. The beam is directed into the cavity with a 45° mirror (M2) that reflects 355 nm light but transmits signal and idler beams. The beam then passes through the phase-matched BBO crystals that generate the signal and idler beams. The beam is then partially (20%) reflected back into the cavity by a coupling mirror (M3) and fully reflected back through the cavity by a retroreflector (M4) to increase mixing efficiency. Additionally, the cavity was seeded with a beam that is tuned to the idler wavelength, resulting in a narrow-band 530.13 nm OPO signal output. This beam is then mixed with the residual 34% 354.82 nm and directed into another Type I BBO crystal cut at

70° relative to the optical axis to undergo sum-frequency generation (SFG) to output a 212.56 nm beam. The SNLO Opoangles widget for the SFG is shown in Fig. 3.4.

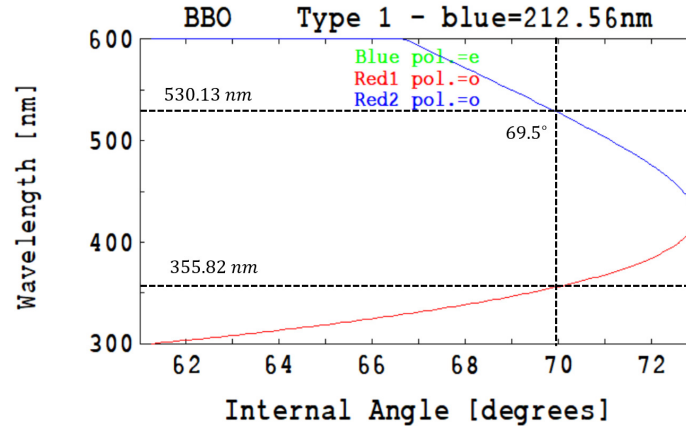


Figure 3.4: Output of SNLO "Opoangles" widget to determine SFG BBO cut angle.

Two additions to the original OPO design were implemented to increase the stability and efficiency of the UV output: 1) Installing BBO heaters (Geola DTC-B) into the OPO at both the cavity and SFG and 2) amplification of the idler seed by a THORLABS YDFA100P polarization maintaining Ytterbium-Doped Fiber Amplifier. The crystal ovens reduced the shot-to-shot instability of the OPO as the BBO crystals were kept at a constant temperature of 50° C for the OPO and 70° C for the SFG. Although the facility environment has a dedicated heat and humidity-control systems to keep it at a constant temperature year round, variations cause enough temperature change to de-phase-match the crystals. Maintaining a steady temperature of 50°C in the OPO ovens has reduced the probability of moisture eroding the BBO crystals.

Implementation of the seed amplifier has increased OPO seeding stability, helping to create a more narrow-band signal for better SFG efficiency. While tuning the OPO cavity, a spectrometer (Ocean Insights HR4 Pro) is used to monitor wavelength and relative beam intensity. A seeded OPO signal will appear narrow-band and at least double in intensity while recording approximately the same power as the un-

seeded case. Fig. 3.5 shows a side-by-side comparison of an unseeded and seeded signal. Introducing the seeded beam to the OPO cavity requires overlap of the beam and the 6 mm beam waist of the pump. After amplifying the seed, the SFG efficiency increased slightly, but more importantly reduced shot-to-shot energy variance. The OPO with all system components is shown in Fig. 3.6.

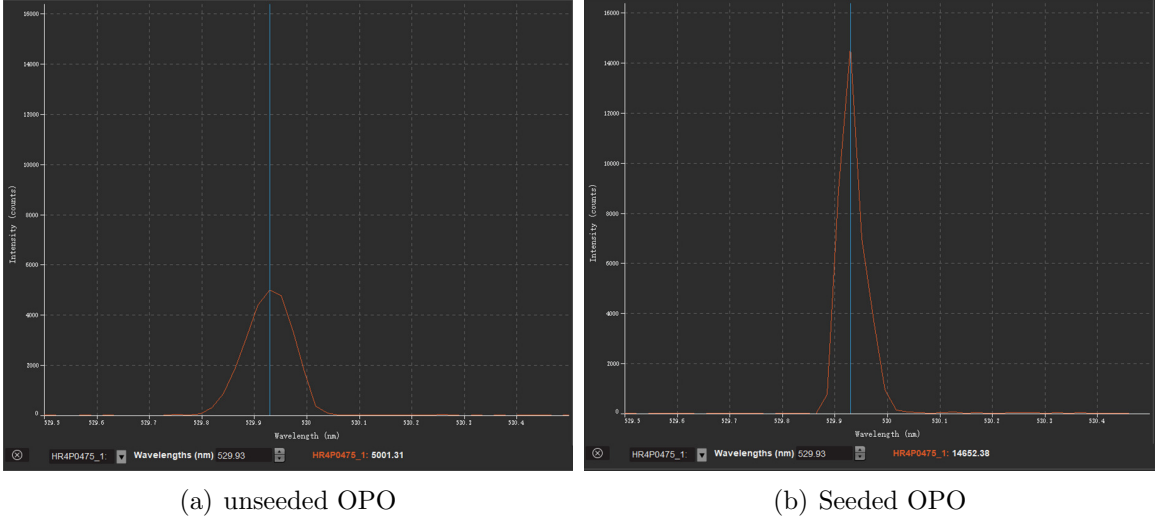


Figure 3.5: Comparison of OPO seeding

3.2.1 Polarization of OPO

Correct polarization of the OPO system is integral to having a highly efficient UV output beam. The polarization of the entire system is shown in Fig. 3.6. The polarization was checked using a 20° Wollaston prism throughout the setup, splitting s-polarization and p-polarization vertically. The third harmonic output of the burst-mode laser is initially p-polarized. A $\lambda/2$ wave-plate rotates the polarization to mostly ($\approx 80\%$) s-polarization before the beam is split and enters the OPO. The BBO crystals undergo Type I phase-matching in which the input pump beam is extraordinary to the optical axis and the output is ordinary to the optical axis. This process is

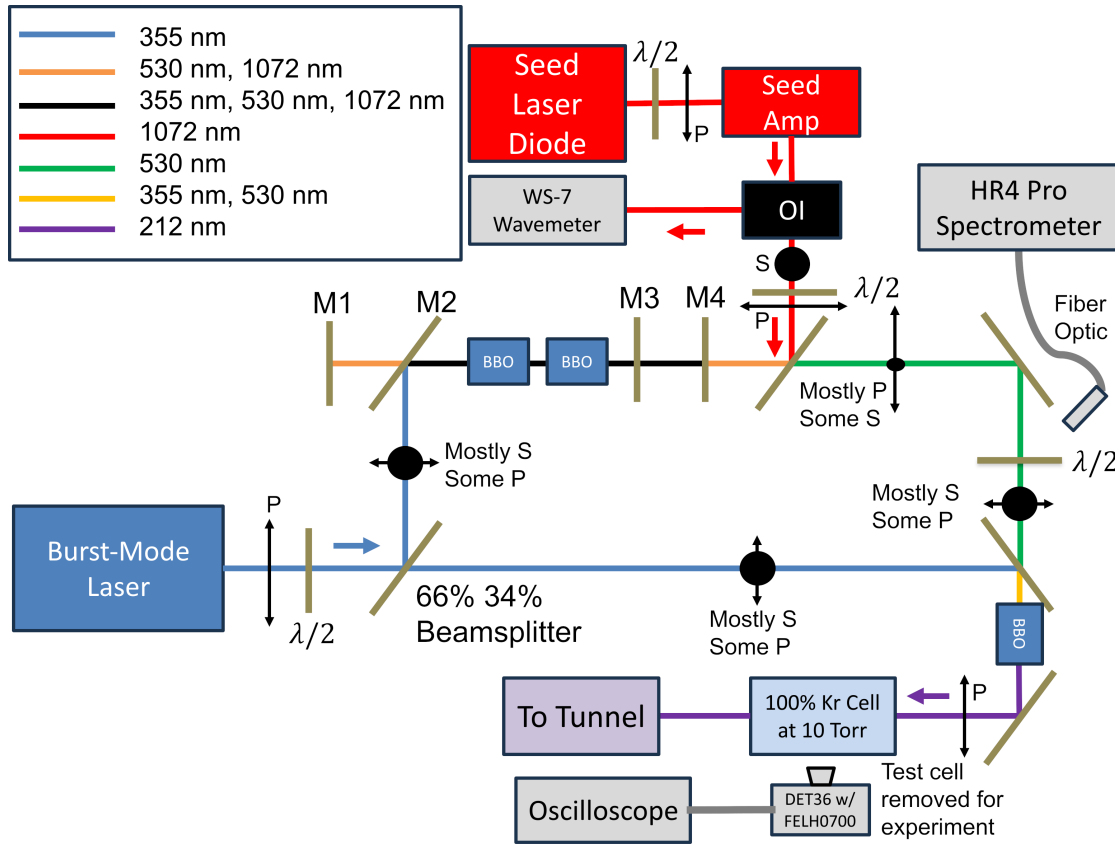


Figure 3.6: Schematic of the OPO system. Polarization of the beam is noted throughout the process. Diagnostic components include a wavemeter, spectrometer, and test cell with photodetector for seed wavelength and beam intensity monitoring.

shown in Fig. 3.7 This changes the polarization from mostly s-polarized to mostly p-polarized. The idler diode output is equally s-polarized and p-polarized. The first $\lambda/2$ wave-plate alters the beam to p-polarization before it enters an optical isolator. The isolator rotates the beam polarization by 90° to s-polarization. Another $\lambda/2$ wave-plate is used to rotate the beam back to p-polarization in which it is introduced to the OPO cavity output. The OPO signal reaches another $\lambda/2$ wave-plate in which it is rotated to mostly s-polarization. The signal and residual pump beams are recombined and directed into the SFG BBO along the axis ordinary to the optical axis and the output 212.56 nm beam is extraordinary to the optical axis. This process is

shown in Fig. 3.8.

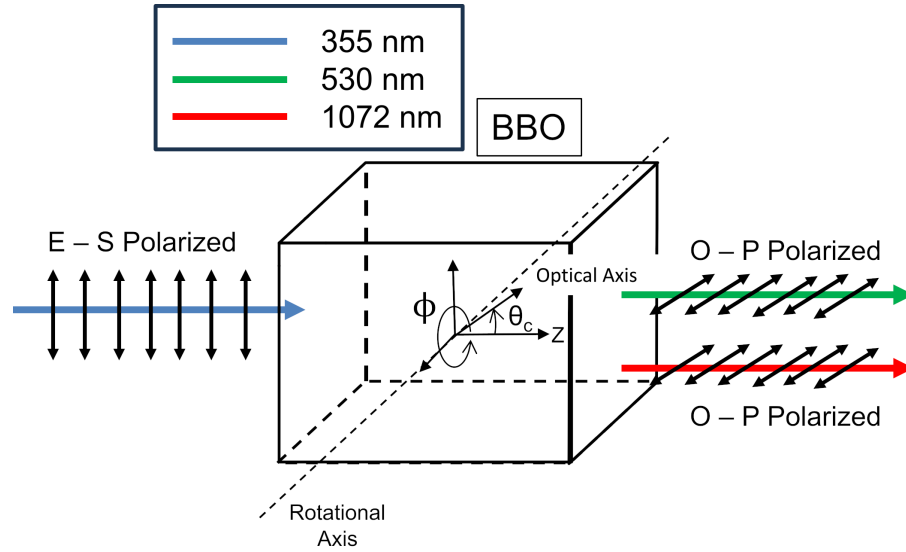


Figure 3.7: Polarization of Type 1 BBO in OPO cavity. Diagram is simplified to show one crystal, although two are used to phase-match signal. The cut-angles of the BBO crystals are relative to the optical axis. Kinematic mounts rotate the crystals about the rotation axis.

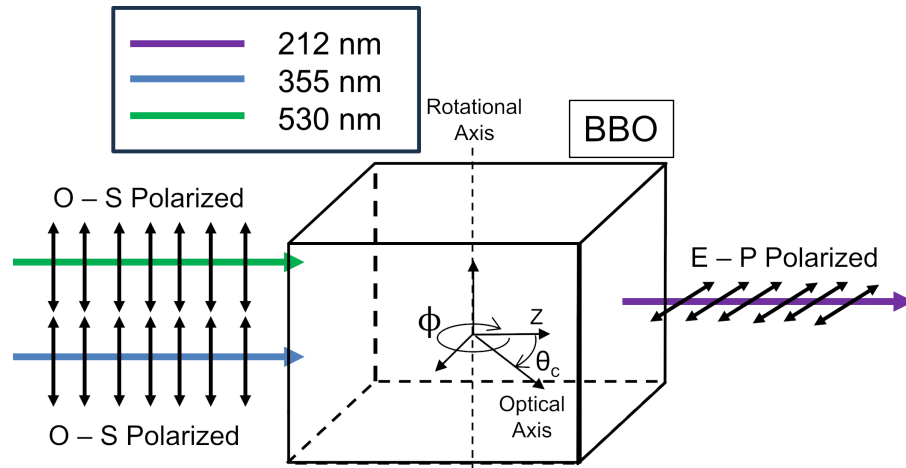


Figure 3.8: Polarization of Type 1 BBO for SFG. The cut-angles of the BBO crystals are relative to the optical axis. Kinematic mounts rotate the crystals about the rotation axis.

3.3 Imaging Equipment

The tagging velocimetry setups utilized two high-speed cameras, a Phantom TMX 7510 and Phantom v2512. Two high-speed cameras were used depending on the experimental configuration and signal intensity requirements. The requirements for capturing sufficient tagging data were the following: 1) high spectral response at emission wavelengths; 2) short exposure gate to reduce signal blurring and increase signal to noise ratio (SNR); 3) capability of reaching 100 kHz level capture frequencies; 4) maintaining high resolution for (3); Both the TMX 7510 and 2512 were capable of meeting most of these requirements as they both boasted high-repetition-rate, low exposure time image gating over a wide spectral response range. However, for the KTV tests, an image intensifier was required to further shorten image gate times and increase SNR. A Video Scope International (VSI) image intensifier (VS4-1845-P24) was introduced to meet these needs. The spectral response curves for the two cameras are shown in Fig. 3.9 and Fig. 3.10 which exhibit much lower quantum efficiency (the number of photoelectrons generated per incident photon at wavelength λ (QE)) in the near infrared (NIR) region where the krypton emission lines are. The intensifier was equipped with a P24 phosphor plate which has a higher QE than the cameras, seen in Fig. 3.11, and a gain dial to further increase the signal brightness. The output signal is at 520 nm which falls near the peak spectral response for both cameras. The gate on the intensifier can be set as low as 10 ns which is shorter than 95 ns and 195 ns of the TMX 7510 and v2512, respectively.

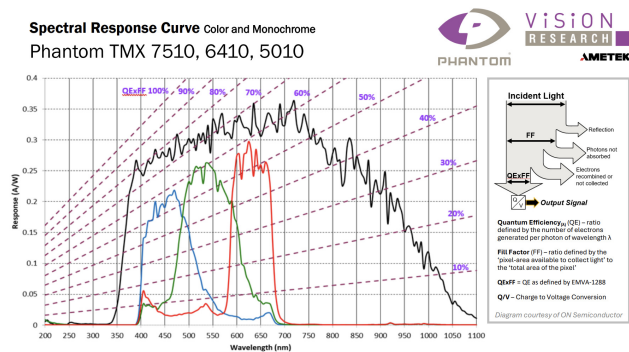


Figure 3.9: Phantom TMX 7510 spectral response curve.

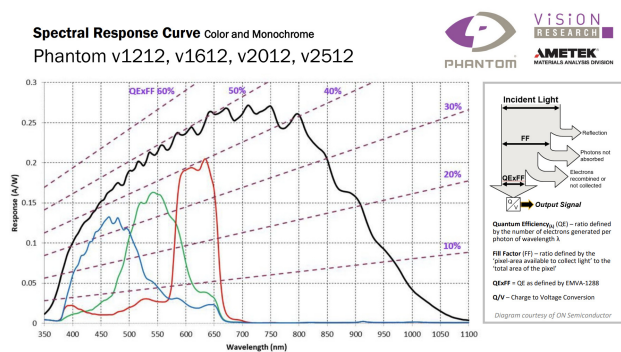


Figure 3.10: Phantom 2512 spectral response curve.

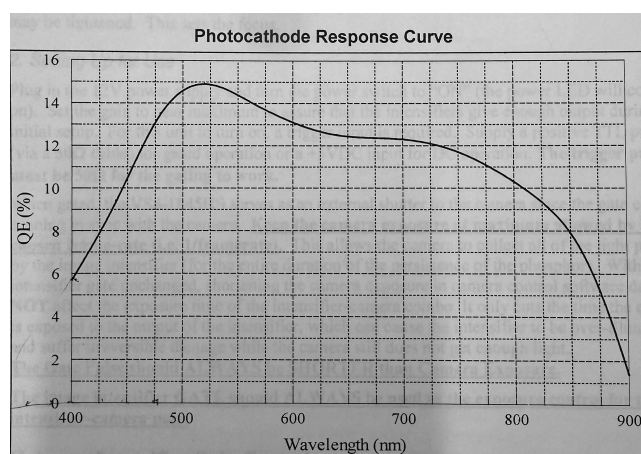


Figure 3.11: VideoScope intensifier spectral response curve.

3.4 Conclusions and Future Work

The inclusion of the burst mode laser system, OPO, Phantom high-speed camera, and image intensifier allowed for the work in future chapters to be completed. One drawback of an OPO is decreased beam quality compared to the output of the burst mode laser. Efforts were made to tune the fundamental wavelength to one that would not need an OPO. For the 5th harmonic of 212.56 nm, a fundamental wavelength of 1062.78 nm is required. As documented by Oishi et al. and Saito et al. [110–112], this technique is possible using Nd:YAG laser rods as it has a peak emission at the required wavelength. Fig. 3.12 shows the emission of Nd:YAG that is currently used in the system. There was hope that the Nd:YAG laser rods would still emit 1062.78 nm light. By modifying the temperature on thermoelectric cooler (TEC), the seed wavelength could be walked down to the desired wavelength. A THORLABS Ccs175 spectrometer and Ophir F110PF-DIF power meter were used to measure the power output and emission. Although the seed was successfully modified, the output intensity was unstable and generated negligible power.

Another attempt was made to generate emission at the Nd:YAG peak of 1073.50 nm, corresponding to the fundamental required to excite the 214.67 nm krypton transition [54, 107]. To improve seed stability, the original seed diode in the OPO was replaced with a Sacher TEC 500 Littman tunable laser. As discussed previously, the nominal idler wavelength of 1072.93 nm lies close to the desired 1073.50 nm and falls near the peak power output of the tunable diode. The seed was amplified using a Thorlabs YDFA100P polarization-maintaining ytterbium-doped fiber amplifier prior to injection into the acousto-optic modulator, which pulses the continuous beam at 100 kHz.

Although the laser system was operated at approximately 80% of its typical

power, amplified spontaneous emission (ASE) at 1064.46 nm remained significantly stronger than the desired 1073.78 nm seed. To attenuate the ASE component, a 1100 nm longpass filter (Thorlabs FELH1100) was mounted on a rotation stage and inserted into the burst-mode laser cavity. Rotation of the filter shifts the effective cutoff wavelength. Using the tunable diode, attenuation angles were determined for both wavelengths: 1073 nm was attenuated at approximately 19° , and 1064 nm at approximately 21° . The filter was therefore set to 20.5° to preferentially transmit 1073 nm while suppressing 1064 nm.

The placement of the filter was optimized to attenuate the 1064 nm emission prior to amplification. Two locations were evaluated: between the second flashlamp and retroreflector on the double-pass leg of the burst-mode laser, and after the double-pass before the third flashlamp. Placement in the double-pass region proved ineffective, as small beam steering of the seed was amplified through the cavity.

Power measurements after the first two amplifiers indicated that the 1073 nm seed was approximately two orders of magnitude weaker than the seeded 1064 nm signal. Although reduced emission strength at this wavelength is expected from the Nd:YAG gain spectrum, a factor of 100 was substantially lower than anticipated. After full amplification, the 1073 nm fundamental output was approximately 10 mJ per pulse at 100 kHz, compared to approximately 250 mJ per pulse at 1064 nm under comparable conditions.

Second harmonic generation (SHG) was then attempted using the 1073 nm seed. The SHG stage employs a LBO crystal in a non-critical phase-matching (NCPM) configuration ($\theta = 90^\circ$, $\phi = 0^\circ$), where conversion efficiency is governed by crystal temperature and increases nonlinearly with input energy. For 1064 nm SHG, optimal temperature is approximately 149°C . For 1073 nm, the predicted NCPM temperature is approximately 137°C and was experimentally observed near 138.5°C . However, the

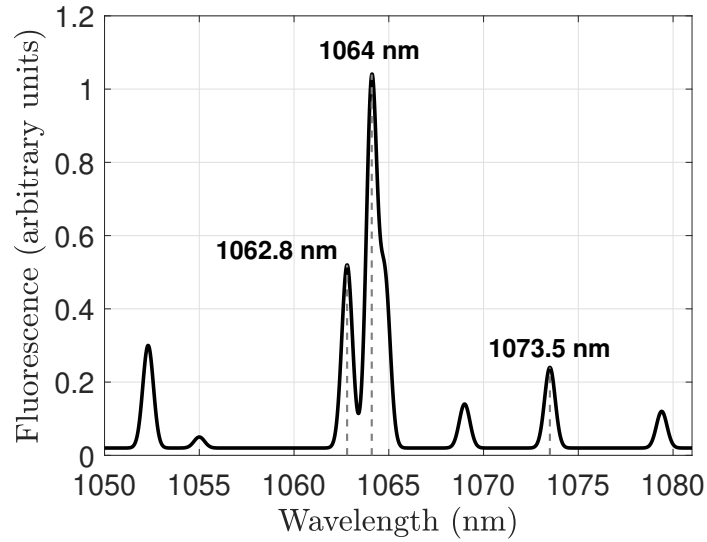


Figure 3.12: Schematic fluorescence spectrum of Nd:YAG near the 1064 nm laser transition. The labeled peaks indicate fluorescence lines near 1062.8, 1064.0, and 1073.5 nm, with the spectrum plotted in arbitrary units based on reported Nd:YAG emission line locations.

harmonic output was limited to approximately 600 μJ per pulse, corresponding to an overall conversion efficiency of approximately 6%, significantly below the nominal 70% efficiency observed at 1064 nm.

Because SHG efficiency scales nonlinearly with input intensity, the reduced fundamental energy at 1073 nm likely limited conversion efficiency. Future efforts will focus on increasing fundamental pulse energy prior to the SHG stage to improve harmonic output.

Chapter 4

Acetone Tagging

This chapter presents the first molecular tagging velocimetry campaign conducted in the Stevens Shock Tunnel using acetone tagging velocimetry (ATV). The goal of these experiments is to obtain high-resolution measurements of streamwise velocity and turbulence statistics in a hypersonic turbulent boundary layer over a hollow-cylinder flare (HCF). These measurements serve as an initial demonstration of tagging velocimetry in the facility and provide experimental data for comparison with computational and experimental studies discussed in Chapter 1.

4.1 Experimental Setup

4.1.1 Ground-Test Facility and Run Conditions

Experiments were conducted in the Stevens Shock Tunnel described in Chapter 2. For the present ATV campaign, the ≈ 0.5 MJ/kg condition at $Re_\infty^u = 16 \times 10^6 \text{ m}^{-1}$ with an N_2 driver was chosen because of the relatively high Reynolds number and modest reservoir temperature ($T_R = 500 \text{ K}$) which precludes the breakdown of acetone in the reservoir (details of this assertion follow). In the current configuration, the facility has a double-diaphragm separating the driver section (5.0 m long), and the driven section (11.07 m long), both of which have a 194 mm inner diameter. The nozzle exit diameter is 406 mm with a core of approximately 330 mm. The hollow-cylinder flare (HCF) test article fits within the inviscid test rhombus, as shown in Fig. 4.4. The HCF is 1 m-long with a 0.102 m outer diameter hollow cylinder and a 34° flare of maximum diameter 0.203 m. A 0.15 m (6 in) long trip composed of 50 grit sanding

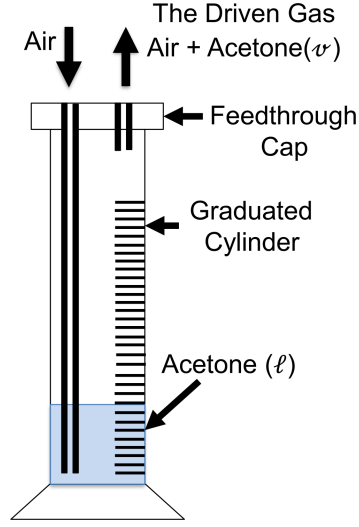


Figure 4.1: Schematic of Acetone Bubbler.

belt was placed 0.26 m (10 in) from the sharp leading edge to induce turbulence.

Acetone vapor was seeded into air as the driven-section test gas via a bubbler, which consisted of a capped graduated cylinder as shown in Fig. 4.1. The graduated cylinder was used to measure the volume of liquid acetone before and after the filling process to determine the mass of acetone vapor in the driven gas. Given the acetone mass and the composition of air, a fixed-point iteration method using Cantera [113] was used to calculate the mole fraction of acetone in the test gas. Run conditions are shown in Table 4.1. Reservoir conditions were calculated with both Cantera and the Shock and Detonation Toolbox [114]. The appropriate thermodynamic data are found in the literature [115, 116].

Freestream conditions were calculated via the method described in Mustafa et al. [60] and [117], utilizing reservoir and Pitot pressure traces and assuming an isentropic expansion to iteratively find freestream Mach number M_∞ . In Fig. 4.2, we show representative reservoir and Pitot-pressure traces as well as the freestream velocity as measured at the boundary-layer edge with ATV. The freestream velocity measurements compare well with the calculated velocity denoted by a dashed green

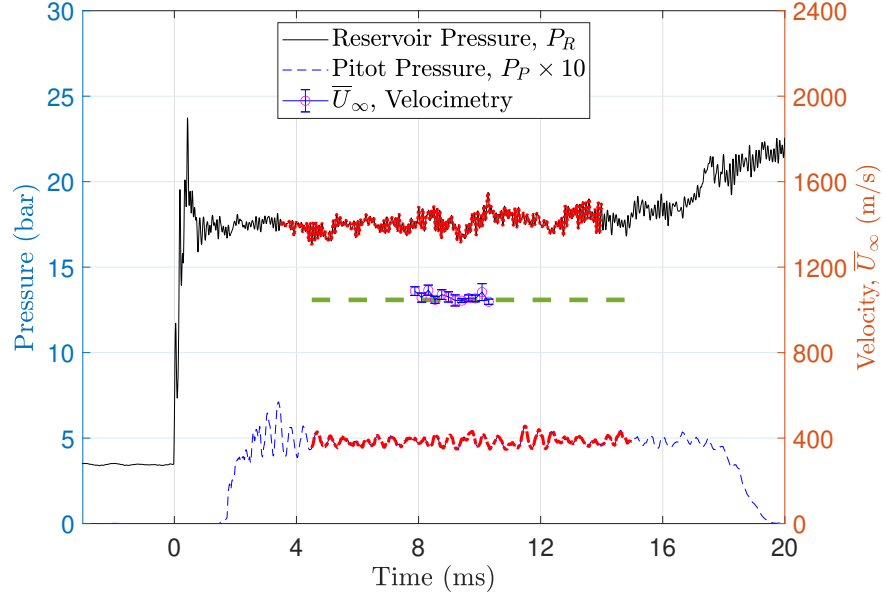
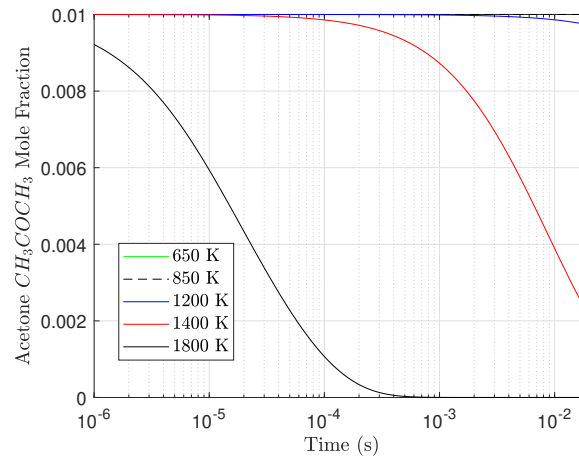


Figure 4.2: Pressure Traces and Freestream Velocimetry Data for Shot 270. Test times are indicated by a thick red linewidth. Run conditions are listed in Table 4.1.

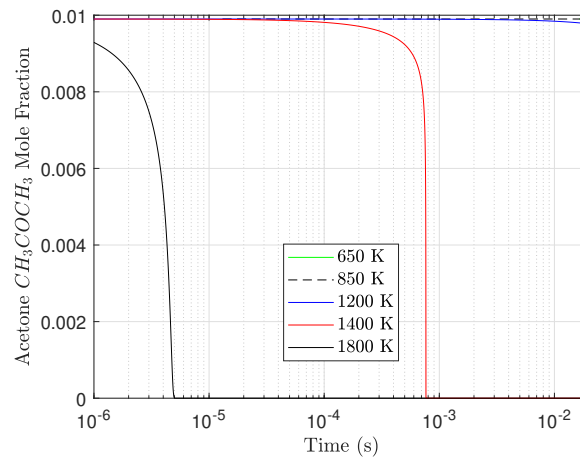
line, bringing confidence to the run-condition calculations. To assess model alignment, static pressure was measured on the surface of the cylinder at azimuthal angles of 60 and 300 degrees at $x = 0.100$ mm and 0, 120, and 240 degrees at $x = 0.600$ mm. These measurements were all within 3% providing confidence to model alignment in the test section and flow uniformity.

Of concern during this campaign was the degree to which the acetone would pyrolyze and react with the test gas. Using nitrogen as the driver gas and a pressure ratio of 10, the reservoir temperature is typically $T_R = 550$ K. Sato and Hidaka [88] found that at a steady state temperature exceeding 573 K, acetone will pyrolyze, which in the present experiments would complicate reservoir-condition calculation and potentially reduce tracer signal. However, the risk of acetone pyrolysis is reduced considerably by the short run times of the tunnel (≈ 20 ms). To investigate this, we perform a kinetics study assuming an initial acetone mole fraction of 0.01 with a

mechanism from [118]. Results are presented for acetone-seeded air and nitrogen at various temperatures in Fig. 4.3. For temperatures below 1200 K for acetone-seeded nitrogen and air at pressure 3.5 MPa, there is little predicted acetone decomposition during the useful test time of the Stevens Shock Tunnel. Therefore, acetone-seeded air can be used for shock tunnel experiments in which the reservoir temperature does not exceed 1200 K, given a reservoir pressure of 3.5 MPa (500 psi). From this, we conclude that there is no appreciable acetone degradation in the reservoir for the current test campaign.



(a)



(b)

Figure 4.3: Acetone Mole Fraction vs. Time in seeded Nitrogen (a) and Air (b).

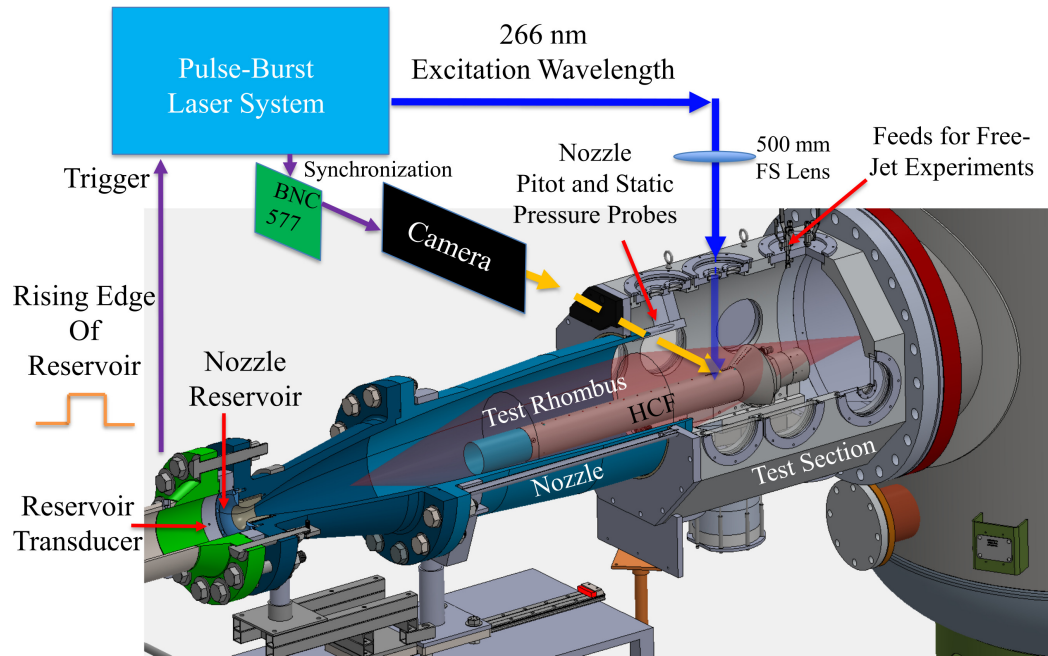


Figure 4.4: Diagram of the experimental setup for tagging velocimetry in the Stevens Shock Tunnel.

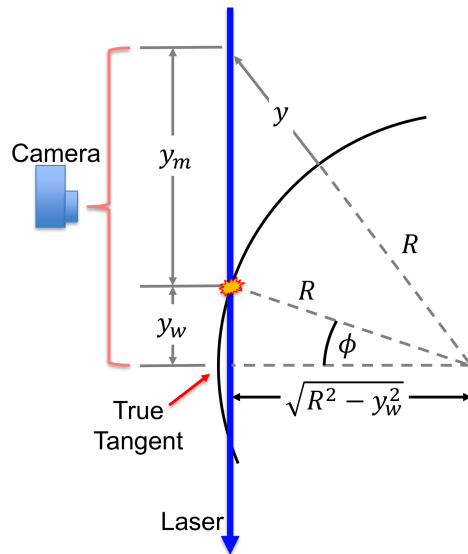


Figure 4.5: Schematic geometrically relating the image coordinate y_m to the wall-normal coordinate y .

Table 4.1: Run Conditions, where P_4 is driver-section nitrogen gas pressure; P_4/P_1 is the driver-to-driven pressure ratio that determines the driven air gas pressure P_1 ; u_s is the shock speed; M_s is the shock Mach number; T_R is the reservoir temperature; P_R is the reservoir pressure; P_∞ is the freestream pressure, T_∞ is the freestream temperature, M_∞ is the freestream Mach number, U_∞ is the time-averaged freestream velocity, and Re_∞^u is the unit Reynolds number.

Shot	P_4 (MPa)	$\frac{P_4}{P_1}$	$X_{C_3H_6O}$	u_s (m/s)	M_s	P_R (MPa)	T_R (K)	P_∞ (kPa)	T_∞ (K)	M_∞	U_∞ (m/s)	Re_∞^u (10^6 1/m)
270	3.44	10	0.0088	504	1.46	1.76	485	1.01	58.7	6.07	935	14.7
273	3.44	10	0.012	511	1.48	1.79	488	1.01	59.0	6.08	940	14.6
274	3.46	10	0.023	505	1.48	2.48	505	1.25	62.9	6.02	960	16.0
276	3.46	10	0.0079	528	1.52	2.11	516	1.25	62.9	6.04	964	16.1
280	3.45	10	0.015	500	1.46	2.17	510	1.21	61.9	6.08	962	16.1
282	3.45	10	0.015	510	1.48	2.12	509	1.26	62.9	6.02	960	16.2
Average	3.45	10	0.013	510	1.48	2.07	502	1.17	61.4	6.05	954	15.6

4.1.2 Experimental Tagging Velocimetry Setup

The motivation of this work is to make high-fidelity observations of high-speed turbulence for comparison to computations. To design a relevant experiment, we must make estimates of the prominent scales in the flow, noting an approximate boundary-layer thickness of $\delta \approx 10$ mm, which was determined from the mean-flow profile, and freestream velocity of $U_\infty \approx 1000$ m/s. As such, we estimate some large structures in the flow to have a timescale of $\tau_L \approx \delta/U_\infty \approx 10$ μ s. We aim to resolve structures of this timescale and smaller. This is why we choose the method of successive imaging of fluorescence of [65, 66] and choose a write/read-delay time of 450 ns with an exposure time of 95 ns. Due to Nyquist, the smallest turbulent time scale we can resolve is twice the 450 ns write/read-delay time, resulting in a Nyquist frequency for each snapshot of approximately 1 MHz.

In contrast to other tagging velocimetry techniques, the laser and imaging setup is relatively simple. The imaging system consisted of a Phantom TMX 7510 with a 135 mm Carl Zeiss lens and a 20 mm lens tube to increase magnification. The image length scale is 15.0 pixels/mm. The camera timing was controlled by a

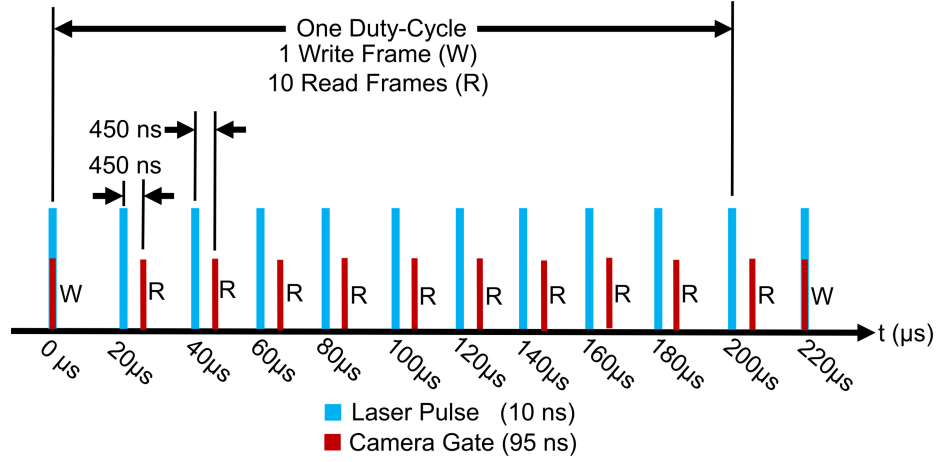


Figure 4.6: Duty-cycle Timing Diagram.

Berkeley Nucleonics Model 577 pulse/delay generator (PDG), which was triggered at 50 kHz by the QuasiModo Quantum Composer (a signal generator inside the laser), resulting in a frame rate of 50 kHz (20 μ s between frames). The PDG synchronized the camera strobe (gate) using its duty-cycle feature, timing the capture of the write images at user-defined locations within a given pulse train. As shown in Fig. 4.6, one write image is captured at zero delay from a single 266 nm laser pulse. Then, ten consecutive read images are captured by gating the camera at a delay of 450 ns from ten respective laser pulses (the write/read-delay time); one laser pulse per read image. Then, the duty-cycle resets and one write image is captured at zero delay from a single 266 nm laser pulse and ten read images are captured at a delay of 450 ns from ten respective laser pulses, and so on.

As shown in Fig. 4.4, ATV experiments utilized an externally triggered Spectral Energies QuasiModo 1200 Burst-Mode Laser System, where we prescribe a 4 ms burst train duration, a repetition rate of 50 kHz, and a pulse duration of 10 ns. The average pulse train energy was 7.75 J for 200 pulses at 266 nm, resulting in an average pulse energy of approximately 38 mJ. Frequency conversion from the fundamental 1064 nm was performed with two successive frequency-doubling crystals built into the system.

Of the 4 ms burst width, the laser generated sufficient intensity for ≈ 3 ms, which corresponds to 150 frames. During a test run, the start of the pulse burst train was triggered by the rising edge of a reservoir pressure transducer. At a prescribed delay (7.5 ms) from this rising edge, the burst-mode laser system pulses at the excitation wavelength into the test section (see the timing in Fig. 4.2). The pulse is directed towards the test section and is focused by an $f = 500$ mm plano-convex lens. The beam enters the test section through a 0.0254 m (1 in) diameter, 3 mm (0.12 in) thick fused-silica glass window.

The approach and coordinate transformation of Mustafa et al. [54] were implemented to interrogate the boundary layer. As shown in Fig. 4.5, the 266 nm excitation laser beam was oriented approximately tangential to the hollow cylinder portion of the test article, in a location upstream of the shock-wave/turbulent boundary-layer interaction to increase wall resolution. The beam was focused to approximately 68 μm near the true tangent of the cylinder. The spot size was determined to be diffraction limited, assuming $M^2 = 4$ (per manufacturer specification) for the 266 nm laser. By grazing the surface of the cylinder, the size of the laser-induced ablation plume at the wall was minimized. The probing location was 0.9 m (35.5 in) from the cylinder leading edge. Using Fig. 4.5 and the Pythagorean Theorem, a coordinate transformation between the measured image coordinate y_m to the wall-normal coordinate y was derived,

$$y = \sqrt{R_c^2 + y_m^2 + 2y_my_w} - R_c, \quad (4.1.1)$$

where R_c is the radius of the hollow cylinder (0.0508 m) and y_w is the location of the no-slip condition from the true tangent of the cylinder, as seen by the camera.

4.2 Results

Acetone tagging results are presented for the shock tunnel experiments whose run conditions are listed in Table 4.1. Acetone mole fractions were $0.9 - 2.3\%$. A write/read-delay time of 450 ns was used to allow sufficient displacement to capture the signal, observe structure, and minimize velocity-measurement error. Sample processed images are shown in Fig. 4.7. Write and read images were processed as follows: (1) cropped to the region of interest; (2) smoothed via a Wiener Filter; (3) zeroed to the minimum value and normalized by the resulting image maximum value; (4) smoothed by 10 pixels across each row, and (5) processed via the peak-finding algorithm of [119], which fits a Gaussian to the peak of each row.

The write lines were fitted to the write images which were recorded every 11

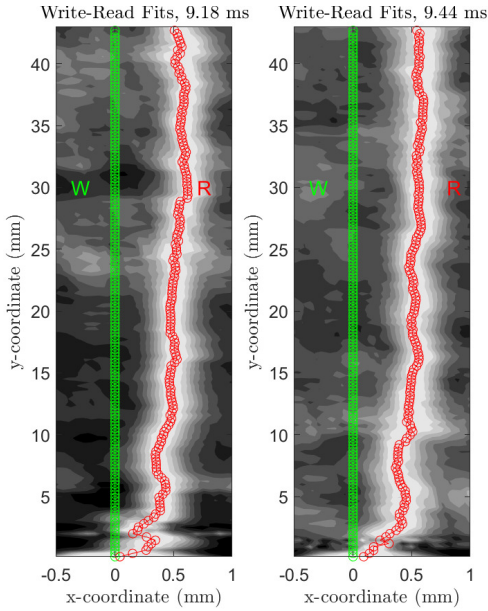


Figure 4.7: Representative write- and read-frame fitting from Shot 280. The letters W and R indicate a write fit and a read fit, respectively. The y-coordinate is y_m in Fig. 4.5.

frames to monitor any displacement between the laser and the HCF during a run. These write lines are marked as W in Fig. 4.7. The read lines, marked as R, were extracted from the processed read images which were recorded for 10 out of 11 frames. Also note that the y-coordinate in Fig. 4.7 is y_m in Fig. 4.5 prior to the scaling of Eq. 4.1.1. Profiles with velocity less than -1000 m/s and greater than $+2000$ m/s were removed from the ensuing analysis. These erroneous samples comprised less than 5% of the data.

Uncertainties for the ATV measurements are calculated in the same fashion as Zahradka et al. [57]

$$\tilde{U} = \sqrt{\left(\tilde{\Delta x} \frac{\partial U}{\partial \Delta x}\right)^2 + \left(\tilde{\Delta t} \frac{\partial U}{\partial \Delta t}\right)^2 + \left(v'_{RMS} \frac{\partial U}{\partial y} \Delta t\right)^2}, \quad (4.2.1)$$

where uncertainty estimates of a variable are indicated with a tilde. The uncertainty in the measured displacement distance, $\tilde{\Delta x}$, of the tracer is estimated as $25 \mu\text{m}$ from the 95% confidence bound on the write and read locations from the Gaussian fits due to a bootstrap error analysis [119]. The uncertainty in time, $\tilde{\Delta t}$, is estimated to be half the camera gate width, 47.5 ns , which causes fluorescence blurring as considered in Bathel et al. [120]. The third term in Eq. (4.2.1) is uncertainty in streamwise velocity due to wall-normal flow in the xy -plane. This formulation is taken from Hill and Klewicki [121] and Bathel et al. [120]. The wall-normal fluctuations used in Eq. (4.2.1) (v'_{RMS}) are conservatively estimated to be 10% of the edge velocity.

For a given experiment, 150 usable frames were recorded ($\approx 3 \text{ ms}$ at 50 kHz). This limitation is imposed by concerns of burning optics in the laser or elsewhere if the pulse train were to be made longer. This is not sufficient to generate statistically significant mean and fluctuating boundary-layer profiles. As such, we repeat the same condition six times to generate sufficient realizations. In Fig. 4.8, profiles are plotted

for the six experiments listed in Table 4.1. Of the six experiments performed, two experiments (Shot 270 and Shot 273) were unable to capture near-wall resolution due to laser-induced breakdown on the surface of the hollow cylinder. The average near-wall velocity-profile data for the remaining four shots are plotted in red and includes 377 realizations while the blue line depicts the average profile for all six experiments and includes 555 realizations.

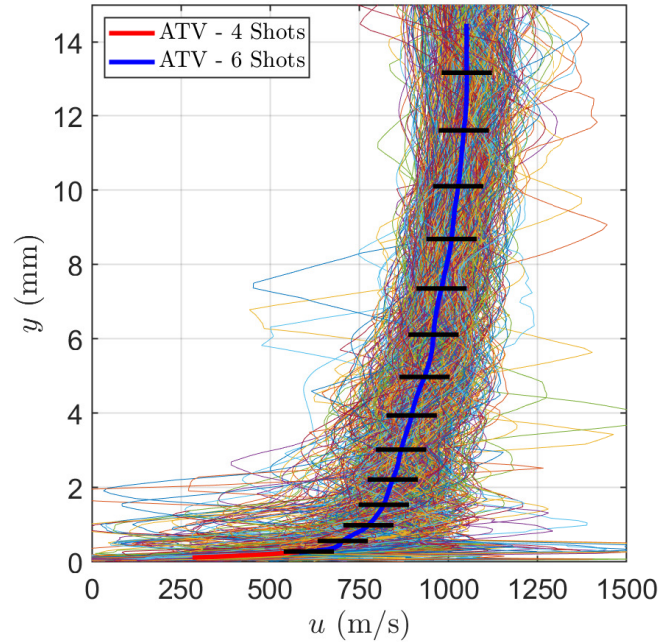


Figure 4.8: All velocity profiles for all six experiments are plotted as thin multicolored lines. The mean velocity profile of all six experiments is blue and the red profile mean of four shots where we found sufficient near-wall SNR. Black error bars due to Eq. 4.2.1.

To average the data for non-dimensionalization, there is a concern that the run conditions are disparate. In Fig. 4.9, a distribution of the velocity at the boundary-layer edge, U_e , is presented along with the computed error of mean velocity. The measured mean velocity is 975 ± 70 m/s (7%), which is well within the computed freestream velocity reported for any experiment in Table 4.1. This brings confidence to

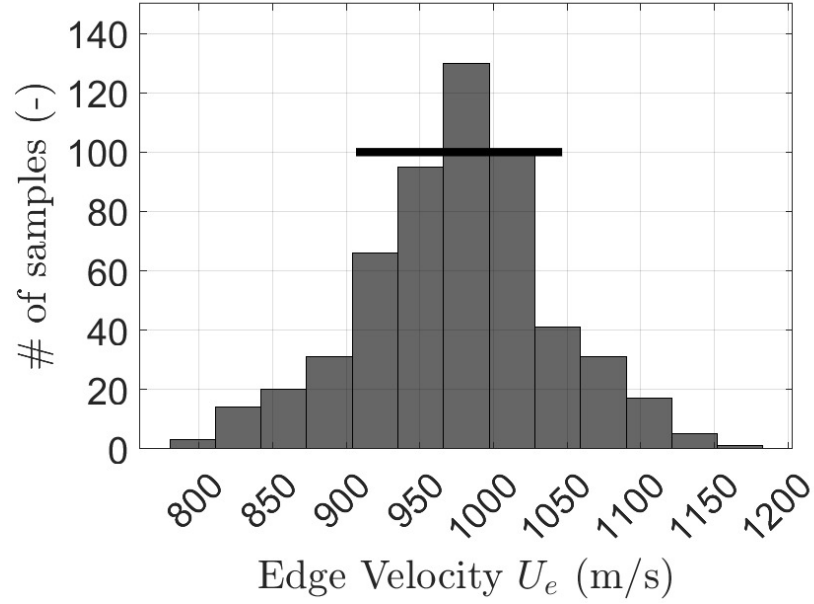


Figure 4.9: Distribution of velocity at the boundary layer edge. Error bar is in black and corresponds to the boundary-layer edge in Fig. 4.8.

the run-condition calculation, as well as the presentation of an averaged run condition at the bottom of Table 4.1.

4.3 Discussion

In this section, the data are non-dimensionalized and compared to data from the literature. To compare data from separate shots where the run conditions were very similar but not identical, the data from each experiment is first non-dimensionalized and then the resulting non-dimensional data are averaged. Note that the alternative method of first averaging all the profiles and then performing the non-dimensionalization results in essentially the same results. The averaged conditions are listed in the bottom row of Table 4.1 and, in an additional form, here in Table 4.2.

Following the formulation of Huang et al. [23], the experimentally-obtained

Table 4.2: Turbulence Quantities

M_e	T_w/T_r	Re_θ	Re_τ	$C_f \times 10^3$	$\frac{u_\tau}{(\text{m/s})}$
6.05	0.65	8925	408	1.07	52.4

velocity can be non-dimensionalized with the van Driest transformation as

$$U_{VD}^+ = \frac{1}{R} \left[\sin^{-1} \frac{R(U^+ + H)}{D} - \sin^{-1} \left(\frac{RH}{D} \right) \right], \quad (4.3.1)$$

where $U^+ = u/u_\tau$, shear velocity $u_\tau = \sqrt{\tau_w/\rho_w}$, $H = B_q/[(\gamma - 1)M_\tau^2]$, $R = M_\tau \sqrt{(\gamma - 1)Pr_t/2}$, and $D = \sqrt{1 + R^2 H^2}$ with parameters defined as follows: $M_\tau = u_\tau/\sqrt{(\gamma - 1)C_p T_w}$, heat transfer coefficient $B_q = q_w/(\rho_w C_p u_\tau T_w)$, and turbulent Prandtl number $Pr_t = 0.9$ [122]. Here, u is the measured velocity, u' is the measured fluctuating velocity, q_w is the heat flux into the flow from the wall, ρ_w is the density of the gas at the wall, C_p is the constant-pressure, mass-specific heat capacity, γ is the specific heat ratio, T_w is the wall temperature, and τ_w is the wall shear stress. The temperature profile is obtained via Walz's use [123] of the Crocco-Busemann relation [20] as

$$\bar{T} = T_w \left[1 + \left(\frac{T_r}{T_w} - 1 \right) \frac{\bar{u}}{U_e} - \frac{\gamma - 1}{2} M_e^2 \left(\frac{T_e}{T_w} \right) \left(\frac{\bar{u}}{U_e} \right)^2 \right] \quad (4.3.2)$$

where U_e is the edge velocity, T_e is the edge temperature, M_e is the edge Mach number, T_R is the recovery temperature and r is the recovery factor estimated to be $0.9 \approx Pr^{1/3}$. The mean density profile is determined as $\bar{\rho} = \rho_w T_w / \bar{T}$, assuming negligible wall-normal pressure gradient. Using the resulting mean density profile, $\bar{\rho}$, the momentum thickness Reynolds number is calculated to estimate a friction factor coefficient, C_f , from a compressibility-corrected friction-factor correlation, which is then used to estimate wall shear stress $\tau_w = (1/2)\rho_e u_e^2 C_f$. Recent work by Huang et

al. [34] states that “[a]mong the various combinations of compressible transformations and incompressible correlations for the skin friction, the combination of van Driest II transformation with the power-law relation of [124] correlates best with the DNS data” at Mach numbers and wall-temperature ratios similar to the present work. The Smits et al. [124] compressibility-corrected friction factor is calculated as

$$C_f = \frac{1}{F_c} 0.024 (F_\theta Re_\theta)^{-1/4}, \quad (4.3.3)$$

where the compressibility corrections are $F_\theta = \mu_e/\mu_w$ and

$$F_c = \frac{(\gamma - 1) r M_e^2}{2(\sin^{-1}(\alpha) + \sin^{-1}(\beta))^2}, \quad (4.3.4)$$

with parameters defined as $\alpha = (2A^2 - B)/\sqrt{4A^2 + B^2}$, $\beta = B/\sqrt{4A^2 + B^2}$, $A = \sqrt{(0.5(\gamma - 1) r M_e^2 / F)}$, $B = (1 + 0.5(\gamma - 1) r M_e^2 - F)/F$, and $F = T_w/T_e$.

Using the Reynolds analogy, the heat flux, q_w , into the flow from the wall can be calculated as

$$q_w = \rho_e U_e C_p (T_w - T_r) C_{he}, \quad (4.3.5)$$

where the Stanton number is $C_{he} = C_f/(2Pr^{2/3})$ [20].

The non-dimensionalized ATV data are plotted for all six shots (555 realizations) for $y^+ > 10$ and for four shots (377 realizations) for $y^+ > 3$ in Fig. 4.10. Two shots were excluded because, in those experiments, the laser-induced breakdown and reflections precluded near-wall measurement. Also included are data from the literature, DNS calculations from [1] and PIV measurements from [5] at conditions similar to those reported in Table 4.2. The error bars of the current ATV data overlap with existing data in the literature. Discrepancies among the data sets are greatest in the wake region which could be Reynolds number dependent.

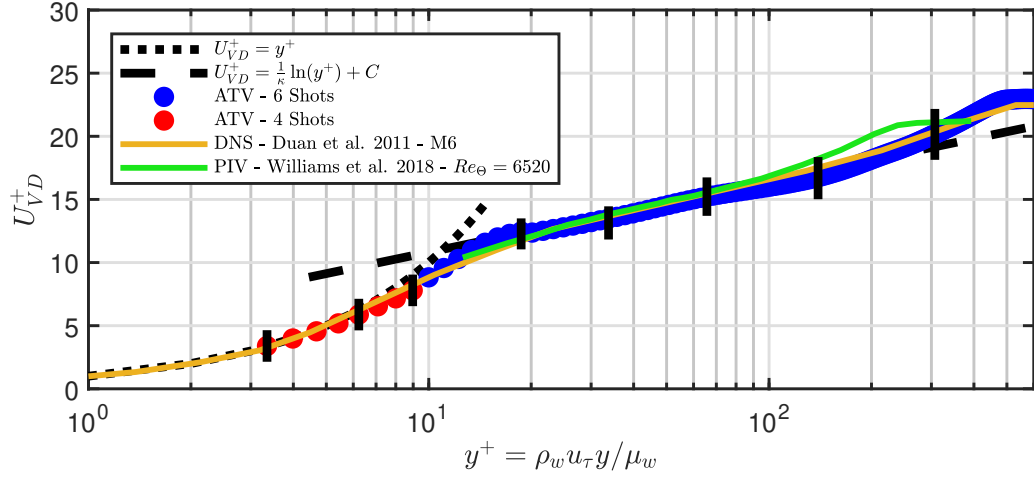


Figure 4.10: Van Driest non-dimensionalization of ATV data and comparison to DNS in [1] and PIV in [5]. Vertical black lines are uncertainty estimates for ATV data.

The Morkovin [3] scaling of the streamwise velocity fluctuations is presented in Fig. 4.11. This scaling, which was intended to account for the density variation in the boundary layer, is also compared against data from the literature, again from Duan et al. [1] and Williams et al. [5]. In Fig. 4.11, the streamwise fluctuations between the present experiments and the literature match well for the subset of 4 experiments above $y^+ > 10$ ($y/\delta > 0.03$). However, the fluctuations reported for $3 < y^+ < 10$ ($y/\delta < 0.03$) do not appear to be physical or follow the trends in the literature. We assert that, in the current experiments, there is noise arising from the measurement technique which artificially increases the reported fluctuation level but does not appear to affect the mean. Additionally, we note that, near the wall, u'_{rms} is the difference between a smaller instantaneous velocity and its mean, and therefore the associated error with tagging velocimetry is larger (see Eq. 4.2.1). The noise from the measurement technique is the result of laser reflections and small laser-induced breakdown at the model surface. For $y/\delta > 1$, the fluctuations do not trend to zero due to (a) camera jitter, which is ≈ 20 ns (less than half of the $\widetilde{\Delta t} = 47.5$ ns uncertainty in time used to determine the error bars in Eq. 4.2.1) and (b) decreased

SNR at and above the boundary-layer edge because the laser was not as focused here.

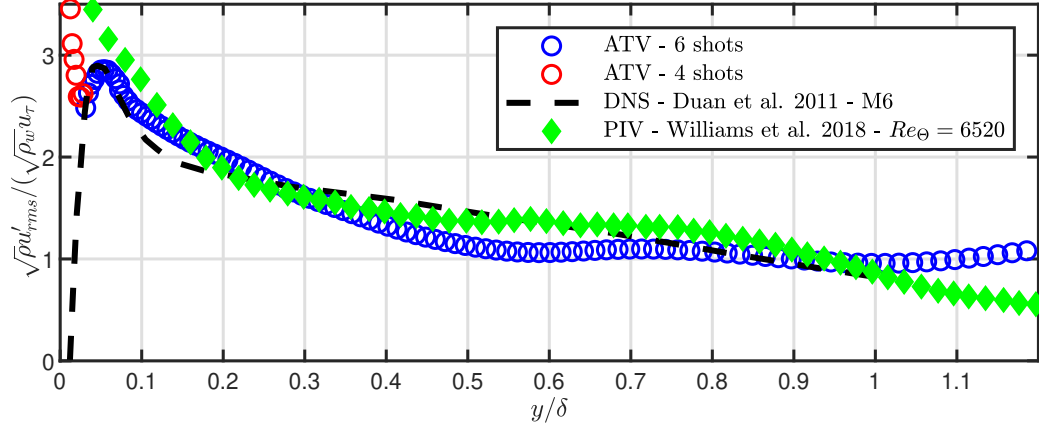


Figure 4.11: Morkovin scaling of ATV data of streamwise fluctuations, u'_{rms} , and comparison to the DNS data in [1] and PIV data in [5].

In Table 4.3, we present some of the relevant length scales in the flow in the context of some of the length scales of the measurement technique to bring confidence in the presentation of results for the mean-velocity profile into the viscous sublayer of a hypersonic turbulent boundary layer. Specifically, the scales of the lowest measurement point and the overlap region are compared to the camera resolution and the calculated beam-waist diameter. The relation between y_m and y is found in Eq. 4.1.1 and illustrated in Fig. 4.5. The mapping of y_m to y appreciably stretches the boundary layer out, especially at the wall, as [125] Section 3.1 showed via a Taylor-series expansion that $y \approx y_m^2 / (2R_c)$ for $y_m / R < 1$ and $y_w \approx 0$. This stretching alleviates concerns that beam waist and camera resolution are on the same order of the measurand. The uncertainty in determining y from recorded images as in Fig. 4.5 and processed with Eq. 4.1.1 is estimated as

$$\tilde{y} = \sqrt{\left(\widetilde{\Delta R_c} \frac{\partial y}{\partial R_c}\right)^2 + \left(\widetilde{\Delta y_w} \frac{\partial y}{\partial y_w}\right)^2 + \left(\widetilde{\Delta y_m} \frac{\partial y}{\partial y_m}\right)^2}. \quad (4.3.6)$$

Table 4.3: Flow and measurement technique length scales.

	y μm	y_m μm	y/δ	y^+
Lowest $\bar{u}$ Measurement	100	1400	0.01	3
Overlap Region	300	3300	0.03	10
Camera Resolution at $y_m = 0$	< 10	67/pix	< 0.001	< 0.3
Beam Waist at $y_m = 0$	< 10	68	< 0.001	< 0.3

Here, we estimate that the uncertainty in the HCF radius is $\widetilde{\Delta R_c} = 0.001$ in ($25.4 \mu\text{m}$) from doubling the machining tolerance. The uncertainty in the recorded location of where the laser strikes the HCF is estimated to be $\widetilde{\Delta y_w} = 134 \mu\text{m}$ (within 2 pixels). The uncertainty in the measured image coordinate is estimated to be $\widetilde{\Delta y_m} = 134 \mu\text{m}$ (within 2 pixels). The spanwise location of where the laser strikes the HCF is $\sqrt{R^2 - y_w^2}$ (Fig. 4.5). Assuming this location is known to within $500 \mu\text{m}$ (that is, the spanwise location of where the laser strikes the HCF) the uncertainty in the transformed wall-normal coordinate is $\tilde{y} = 20 \mu\text{m}$.

4.4 Conclusion

Acetone tagging velocimetry (ATV) with nominally 1 % mole fraction acetone seeding was conducted in a shock-tunnel-generated hypersonic turbulent boundary layer over a hollow-cylinder flare at $M_e = 6.05$, $T_w/T_r = 0.65$, $Re_\theta = 8925$, and $Re_\tau = 408$. This single-laser tagging velocimetry experiment, which recorded mean and fluctuating streamwise velocity ($\bar{u}$ and u'_{RMS}), had a repetition rate of 50 kHz and did not require the use of an image intensifier. Results are presented for six experiments at nominally the same run condition. The measured velocity at the boundary-layer edge is in good

agreement with the freestream velocity calculated with the Stevens Shock Tunnel run-condition calculation procedure.

The ATV data are non-dimensionalized via the van Driest transformation and compared to literature with DNS data from Duan et al. [1] and PIV measurements from Williams et al. [5]. For all six shots at $y^+ > 10$ (555 realizations), the velocity compares well with the literature, and for four shots (377 realizations), velocity was reported down to $y^+ \approx 3$; this is into the viscous sublayer. Two shots were excluded because, in those experiments, the laser-induced breakdown and reflections precluded near-wall measurement. The Morkovin-scaled fluctuations matched the literature for $y^+ > 10$; however, below $y^+ \approx 10$, it appears that the noise introduced by the measurement technique, in addition to the velocity differences being small near the wall, rendered the ATV-measured Morkovin-scaled fluctuations below $y^+ = 10$ not useful.

Future work could involve manipulating the photophysics of acetone with an additional CW laser, as is done with certain applications of other tagging techniques, could be performed to increase the tracer lifetime, increase SNR, and decrease required laser-pulse energy. Alfonso et al. [126] published preliminary results of using a 470 nm laser diode to re-excite acetone photolysis byproducts. The 470 nm diode, with a frequency of 10 Hz and power of 20 mJ/pulse, was introduced to the experiment with a delay of 400 ns from the initial acetone excitation. Fluorescence of the gas was observed and recorded but an understanding of the composition and fluorescence lifetime of the re-excited gas remain areas of interest.

Of technical interest will be efforts to increase the frequency and pulse count from the burst-mode laser to get statistically significant results out of a single experiment. In the current configuration, at least four experiments are needed to obtain 300 samples, a benchmark for statistically significant data. Increasing the pulse-burst du-

ration while maintaining enough pulse power and intensity will increase test efficiency and decrease effects from test-to-test variability.

Chapter 5

Streamwise Tagging Using KTV

This chapter continues the experimental in the SST using KTV and was featured in a publication, Segall et al. [127].

5.1 Experimental Setup

Experiments were performed in the Stevens Shock Tunnel. The Mach 6 enthalpy-matched condition has a test time of 5 ms at unit Reynolds number up to 8×10^6 1/m (a 3×10^6 1/m condition is used for this work) and is executed with a He driver. For the cold-flow condition, a N₂ driver was used, providing a longer 15 ms test time and an increased Reynolds number of up to 32×10^6 1/m (a 14×10^6 1/m condition is used for this work). This reduces the reservoir enthalpy but provides sufficient heating such that the flow in the freestream does not condense. For this campaign, the test gas was 95% N₂ / 5% Kr. Doping the flow with 5% Kr slightly alters the transport properties and freestream conditions which are calculated with Cantera [128] and the Shock and Detonation Toolbox [129] following the methods in Korte et al. [117] and Segall et al. [130]. For reference, seeding the N₂ test gas with 5% Kr changes the ratio of specific heats from 1.400 to 1.407 or $\approx 0.5\%$. The conditions from this campaign and those for some in the relevant literature are presented in Table 5.1. Five shock tunnel runs were executed for both the streamwise and wall-normal tests for the cold-flow condition. Each run yields ≈ 70 realizations for a total of ≈ 350 realizations per data set. Three shock tunnel runs were executed for the streamwise and five for the wall-normal tests for the enthalpy-matched condition. All tunnel run data are available in Table B.1 in Appendix B.

Table 5.1: Run conditions. P_R and T_R are the reservoir pressure and temperature, respectively. P_∞ , T_∞ , U_∞ , M_∞ , and Re_∞^U are the freestream pressure, temperature, velocity, Mach number, and unit Reynolds number, respectively. $u_\tau = \sqrt{\tau_w/\rho_w}$, δ , $Re_\tau = u_\tau\delta/\nu_w$, $Re_\theta = \rho_\infty u_\infty\theta/\mu_\infty$, and T_w/T_r are the friction velocity, boundary-layer thickness, friction Reynolds number, momentum thickness Reynolds number, and wall-to-recovery-temperature ratio, respectively.

Method	P_R	T_R	P_∞	T_∞	U_∞	M_∞	Re_∞^U	u_τ	δ	Re_τ	Re_θ	T_w/T_r
	(MPa)	(K)	(kPa)	(K)	(m/s)	(-)	(1/m)	(m/s)	(mm)	(-)	(-)	(-)
PIV - Williams et al.	9.60	542	1.37	48.3	1046	7.45	4.42×10^6	53.8	10.2	280	9340	0.79
LDA - DeGraaff and Eaton	-	-	-	-	18.95	0.056	1.30×10^6	0.727	35.87	1700	5200	1.00
DNS - Zhang et al.	0.93	433	0.68	55.0	870	5.86	9.95×10^6	45.1	23.8	453	9455	0.76
DNS - Zhang et al.	3.40	693	0.39	51.8	1155	7.87	8.16×10^6	54.3	35.2	480	9714	0.48
KTV - (N ₂ Driver)	2.18	533	1.37	67.5	932	5.82	15.0×10^6	44.1	10.5	450	9880	0.54
KTV - (He Driver)	2.43	1698	1.41	223.8	1740	5.97	2.89×10^6	49.5	15.0	780	2047	0.17

The test article is a hollow-cylinder flare (HCF) that fits within the inviscid test rhombus. The HCF is 1 m-long with a 0.102 m (4 in) outer diameter hollow cylinder and a 34° flare of maximum diameter 0.203 m (8 in). Three 40-grit sanding belt trips were placed on the test article to induce turbulence with distributed roughness: (1) 0.076 m (3 in) long at 0.026 m (1 in) from the sharp leading edge; (2) 0.076 m (3 in) long at 0.137 m (5.375 in); and (3) 0.222 m (8.75 in) long at 0.267 m (10.5 inches). The tagging location is 0.84 m (32.9 in) downstream of the HCF leading edge. This is $Re_x = 11.5 \times 10^6$ and $Re_x = 2.4 \times 10^6$ for the cold-flow and enthalpy-matched conditions, respectively, and at least 45δ downstream of the end of the trip. The boundary layer is assumed to be in an equilibrium, turbulent state as supported by focused laser differential interferometry (FLDI) measurements reported in Segall et al. [127] and Chen et al. [131]. The FLDI measurements were made upstream of the KTV measurements at 0.7 m (27.6 in) or $Re_x = 9.8 \times 10^6$ and $Re_x = 2.0 \times 10^6$ from the HCF leading edge for the cold-flow and enthalpy-matched conditions, respectively. Those results match those of DNS, bringing confidence to the understanding of the state of the boundary layer.

Following Mustafa et al. [54] and Shekhtman et al. [107], KTV is performed by issuing a laser beam into the flow to excite the Kr along a line. “Write” and

Table 5.2: Relevant NIST Atomic Spectra Database Lines Data, labels match Fig. 5.1. Racah nl[K]_J notation. Transition J is not listed because it represents the recombination process and is not an atomic-level transition. G/H/I represents ranges and order of magnitude estimates since they represent numerous transitions in the 5p-5s band.

Transition	λ_{air} (nm)	Nature	Lower Level	Upper Level	A_{ij} (1/s)	E_j (cm ⁻¹)	E_i (cm ⁻¹)
A	212.556	Two-Photon	$4s^2 4p^6, ^1S_0$	$5p[1/2]_0$	(-)	0	94092.8626
B	212.556	Single-Photon	$5p[1/2]_0$	Kr+	(-)	94092.8626	112917.62
C	758.950	Single-Photon	$5s[3/2]_1$	$5p[1/2]_0$	4.31×10^7	80916.7680	94092.8626
D/E	769.454	Single-Photon	$5s[3/2]_2$	$5p[3/2]_1$	4.3×10^6	79971.7417	92964.3943
F	829.811	Single-Photon	$5s[3/2]_1$	$5p[3/2]_1$	2.9×10^7	80916.7680	92964.3943
G/H/I	750-830	Single-Photon	5s	5p	$10^6 - 10^7$	80000.0000	90000.0000

“read” exposures of laser-induced fluorescence (LIF) from the excited and recombining Kr/Kr⁺ are recorded at a prescribed delay. The relevant transitions are marked in blue and red in the energy level diagram in Fig. 5.1, along with the relevant transition data in Table 5.2. The “write step” is performed by exciting Kr atoms to form the tagged tracer through a (2+1) resonance-enhanced, multiphoton ionization (REMPI) process. First, two-photon excitation occurs via $4p^6(^1S_0) \rightarrow 5p[1/2]_0$ (two 212.6 nm photons, transition A) with subsequent one-photon ionization (one 212.6 nm

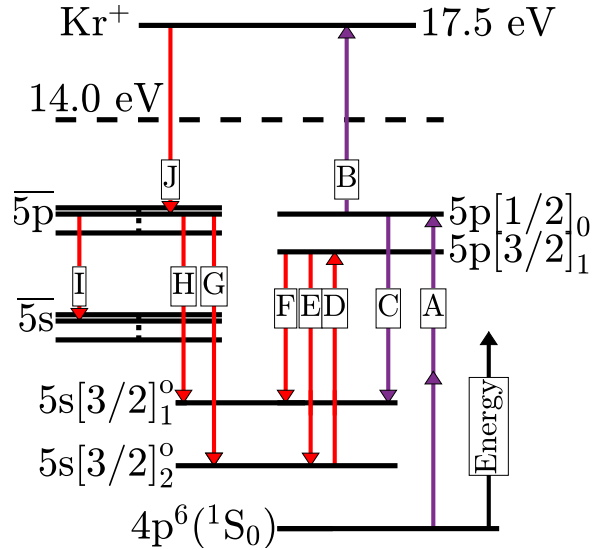


Figure 5.1: Energy diagram for Kr, Racah nl[K]_J notation. Write and read steps in blue and purple, respectively.

photon, transition B). Fluorescence for the write step is recorded primarily from the decay to the resonance state $5p[1/2]_0 \rightarrow 5s[3/2]_1^o$ (758.9 nm, transition C). After a prescribed delay, the “read step” is performed. The displacement of the tagged Kr is imaged from transitions E-I resulting from the recombination process J. The streamwise measurements made use of a CW 769.5 nm tunable laser diode to boost signal with re-excitation of the metastable state discussed in [107, 132] (transition D in Fig. 5.1). This boosts the signal by up to $\approx 50\%$ and was not used for the wall-normal experiments.

The 100 kHz, 212.6 nm laser pulses for (2+1) REMPI were generated using an externally triggered Spectral Energies QuasiModo 1200 Burst-Mode Laser (Fig. 5.2-top) previously discussed in Chapter 3. The fundamental Nd:YAG output near 1064.5 nm was frequency tripled to produce 354.8 nm pulses. A beam splitter was used to direct 66% of the power to pump an optical parametric oscillator (OPO) described in Grib et al. [101]. The OPO output 530.1 nm pulses, which were then sum-frequency mixed (SFM) with the residual (33%) 354.8 nm pulses to form the 212.6 nm beam. The OPO process is shown in Fig. 5.3.

For the streamwise measurements, the 212.6 nm beam was directed into the test section such that it traversed the spanwise direction and was focused at a point tangent to the HCF just upstream of the flare with a 300 mm lens, (Fig. 5.2-top). This was done following Segall et al. [130] to reduce the laser-induced ablation plume and increase wall resolution using a coordinate transformation. The imaging system for the streamwise measurements was a Carl Zeiss 85 mm lens, a 28 mm lens tube, a Video Scope VS4-1845HS image intensifier (50 ns gate for the cold-flow condition and 20 ns gate for the enthalpy-matched condition), and a Phantom 2512 high-speed camera (10 pixel/mm resolution). The inclusion of the image intensifier for these experiments and its benefits are discussed in Section 3.3. The prescribed write/read

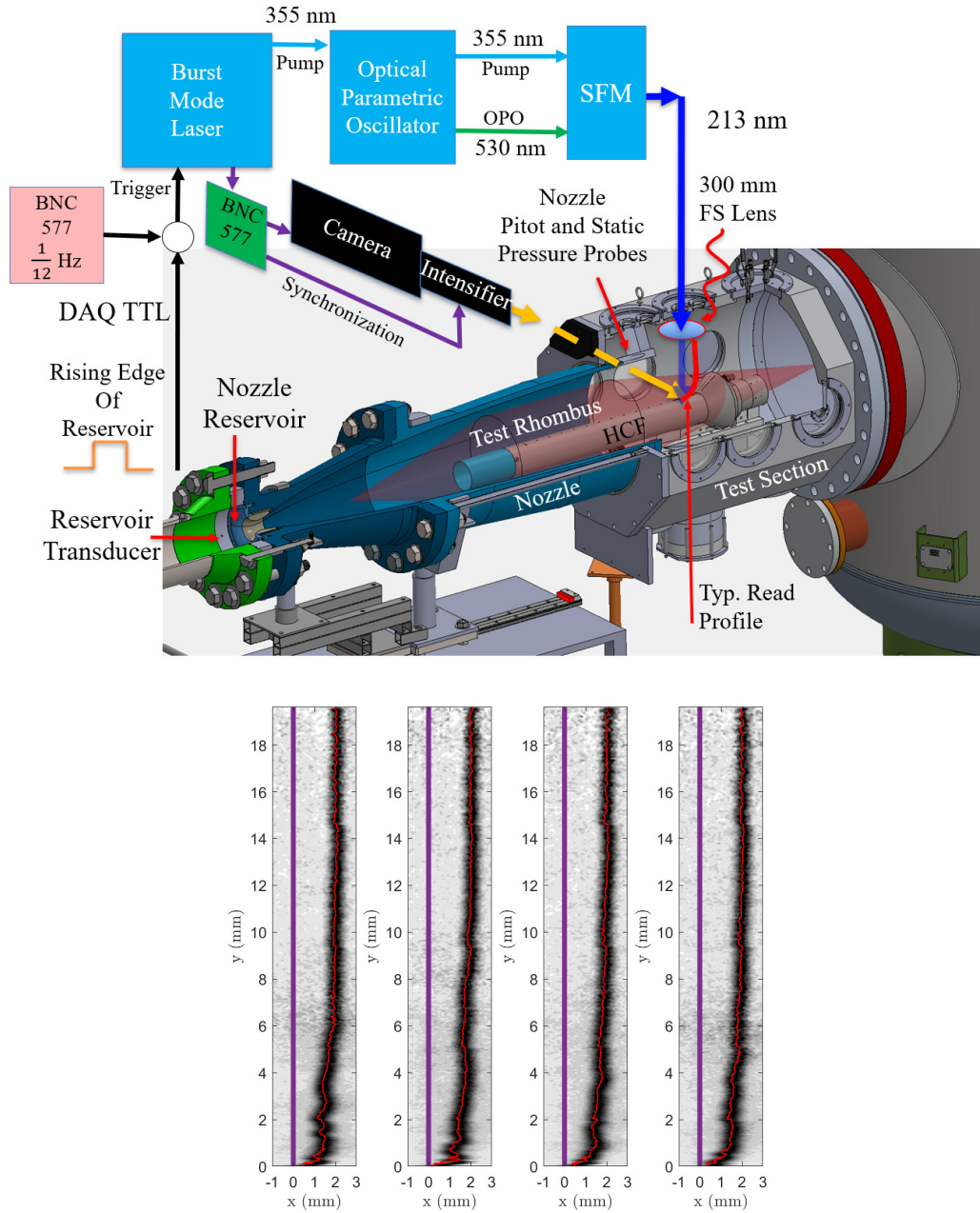


Figure 5.2: Top: Setup for streamwise fluctuation measurement. Bottom: Representative write (purple) read (red) lines. $\delta \approx 10.5$ mm. The ordinate is y , the wall-normal coordinate obtained from the coordinate transformation in Fig 4.5.

delay was $2 \mu\text{s}$ for the cold-flow condition and $1 \mu\text{s}$ for the enthalpy-matched condition.

Sample write/read exposures are shown in Fig. 5.2 (bottom).

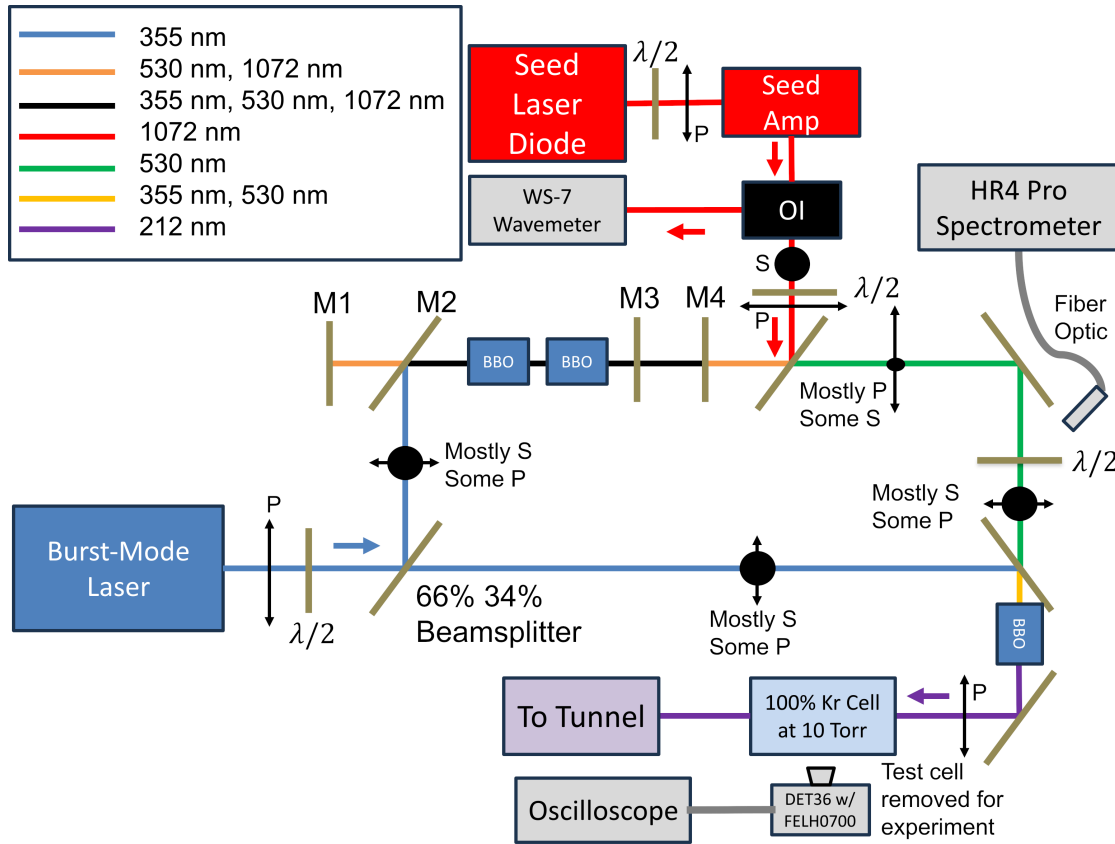


Figure 5.3: Schematic of the OPO system. Polarization of the beam is noted throughout the process. Diagnostic components include a wavemeter, spectrometer, and test cell with photodetector for seed wavelength and beam intensity monitoring.

5.2 Results and Discussion

The write/read image pairs are processed in the same way as Segall et al. [130] and Chapter 4 to yield dimensional velocity profiles shown in Fig. 5.4. The translated fluorescing lines are captured with the intensified camera and Gaussian fits are made to find their location. The streamwise data processing procedure is as follows:

- (1) Binning by six pixels vertically (collapsing in the wall-normal direction);
- (2) Cropping to the desired region;
- (3) Smoothing using a Wiener filter with a 2x8 pixel stencil;

- (4) Normalizing the image intensity;
- (5) Using the peak-finding algorithm of O'Haver [119] to fit a Gaussian to the peak intensity value of each row. The displacement from the write line to the read line, Δx , over the prescribed delay, Δt , resulted in the velocity profile for each frame, $u = \Delta x / \Delta t$. For reference, see the sample exposures in Fig. 5.2 (bottom).

Following Segall et al. [133], uncertainty in the streamwise KTV measurements was estimated as

$$\tilde{U} = \sqrt{\left(\tilde{\Delta x} \frac{\partial U}{\partial \Delta x}\right)^2 + \left(\tilde{\Delta t} \frac{\partial U}{\partial \Delta t}\right)^2 + \left(\widetilde{v'_{RMS}} \frac{\partial U}{\partial y} \Delta t\right)^2}, \quad (5.2.1)$$

where uncertainty estimates of a variable are indicated with a tilde. The uncertainty in the measured displacement distance, $\tilde{\Delta x}$, of the tracer is estimated as 10 μm , which is the 95% confidence interval of peak location for the Gaussian fit. The uncertainty in time, $\tilde{\Delta t}$, is estimated to be half the intensifier gate width (25 ns), which causes fluorescence blurring as considered in Bathel et al. [120]. The third term in Eq. (5.2.1) is uncertainty in streamwise velocity due to wall-normal flow in the xy -plane [121]. The wall-normal fluctuation uncertainty used in Eq. (5.2.1) is conservatively estimated to be $\widetilde{v'_{RMS}} = u_\tau \sqrt{\rho_w / \rho}$, which is $\approx 2\text{-}5\%$ of the velocity at the boundary-layer edge. The error bars on the streamwise KTV data for both the cold-flow and enthalpy-matched conditions, seen in Fig. 5.4 and Fig. 5.7, are a result of these uncertainties. Additionally, the distribution of velocity at the boundary layer edge is plotted in Fig. 5.5 for both conditions to confirm that the experiments are not disparate. The measured mean velocities is $943 \pm 65 \text{ m/s}$ (7%) and $1780 \pm 124 \text{ m/s}$ (7%), are well within the computed freestream velocity reported for the experiments in Table 5.1.

To understand how close KTV can get to the wall, Fig. 5.6 maps each point

in the image coordinate (y_m) to the wall-normal coordinate (y) using the geometry transformation in Fig. 4.5. The close-up of the mapping shows that with the geometric transformation, the technique can get very high resolution, on the order of microns.

The mean streamwise velocity is then nondimensionalized via the van Driest

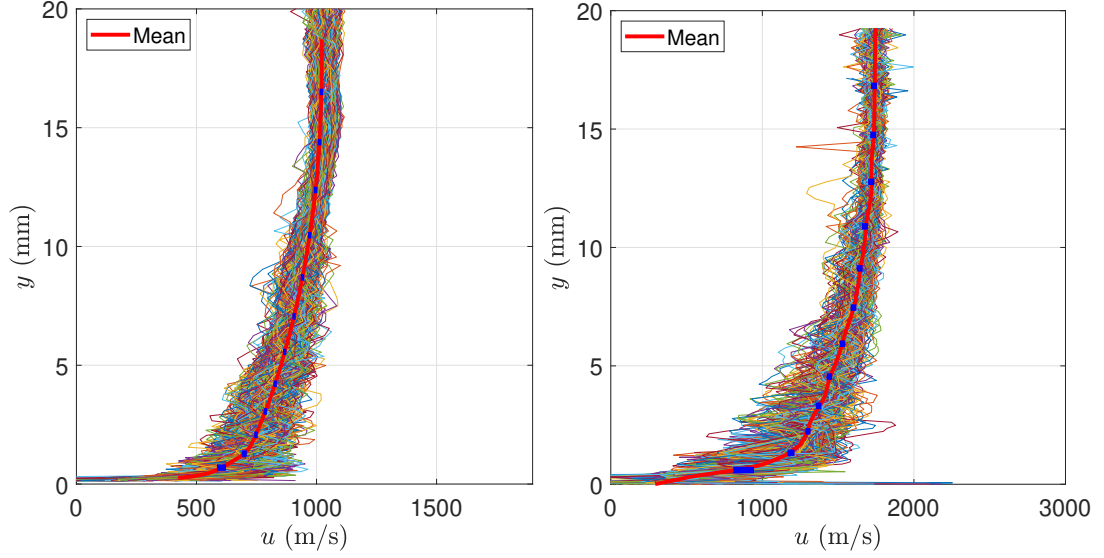


Figure 5.4: All velocity profiles for all cold-flow (left) and enthalpy-matched (right) conditions, plotted as thin multicolored lines. The mean velocity profile of all experiments is red. Blue error bars due to Eq. 5.2.1.

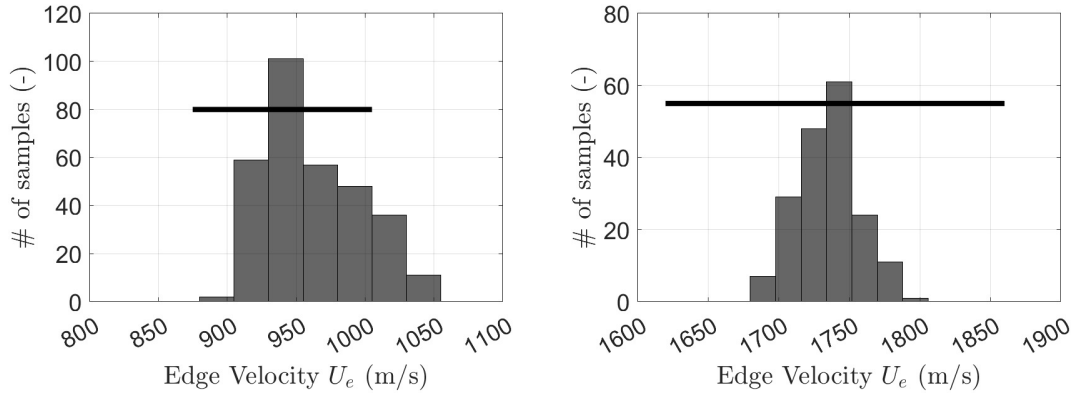


Figure 5.5: Distribution of velocity at the boundary layer edge for cold-flow (left) and enthalpy-matched (right) conditions. Error bar is in black and corresponds to the boundary-layer edge in Fig. 5.4.

[134] transformation following [23, 34] with the local mean density profile coming from the generalized Reynolds analogy from Zhang et al. [25], as opposed to the simpler Walz [123] relation. The skin-friction coefficient is evaluated following Zhao and Fu [36], replacing the previous technique of using the Kármán–Schoenherr correlation and the van Driest II [134] transformation. The Zhao and Fu [36] method is chosen, as the van Driest II method struggles to map the compressible scaling of skin-friction to the incompressible relation for a highly cooled wall, as confirmed by Huang [34] and Zhao and Fu [36]. Nondimensional streamwise velocity from computations (DNS), comparable experiments (PIV), and the present work (KTV) data are presented in Fig. 5.7. Agreement among the methods is within the estimated uncertainty of the current measurement technique, bringing confidence to the results.

The streamwise (u'_{RMS}) fluctuation data are shown in Fig. 5.8 as normalized by the Morkovin scaling for computations (DNS), comparable experiments (PIV), and the present work (KTV). Additionally, the HWA data from Klebanoff [37] is shown as a point of comparison for incompressible flow. The data are presented in this way to demonstrate the validity of the Morkovin hypothesis. That is, the fluctuation profiles may be transformed from the compressible form to a similar incompressible

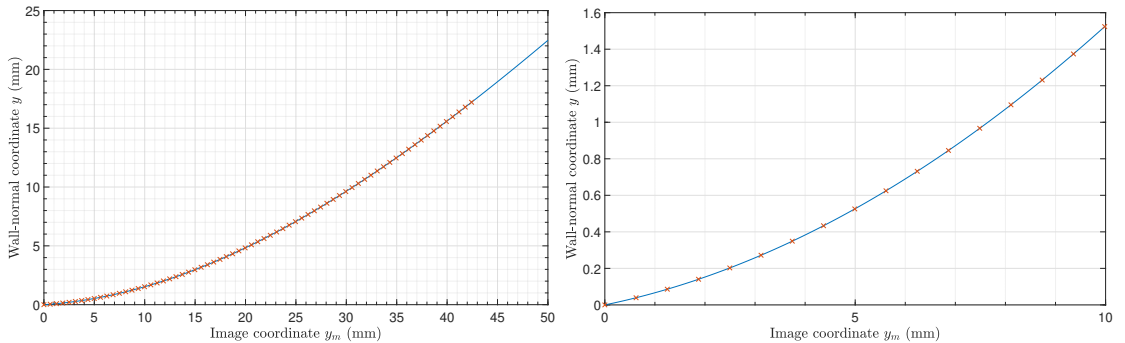


Figure 5.6: Mapping of the image coordinate (y_m) to the wall-normal coordinate (y) using the geometry transformation in Fig. 4.5. The right image is a zoomed in version of the left.

form by simply accounting for changes to the mean flow variables in the boundary layer and at the wall [4]. The DNS, PIV, and KTV streamwise fluctuation data are in reasonable agreement with the incompressible data when appropriately transformed. The streamwise KTV fluctuation data show good agreement with DNS from Zhang et al. [4]. KTV data indicate a higher level of noise outside the boundary layer than in the DNS. We speculate that this is the result of the noisy tunnel freestream.

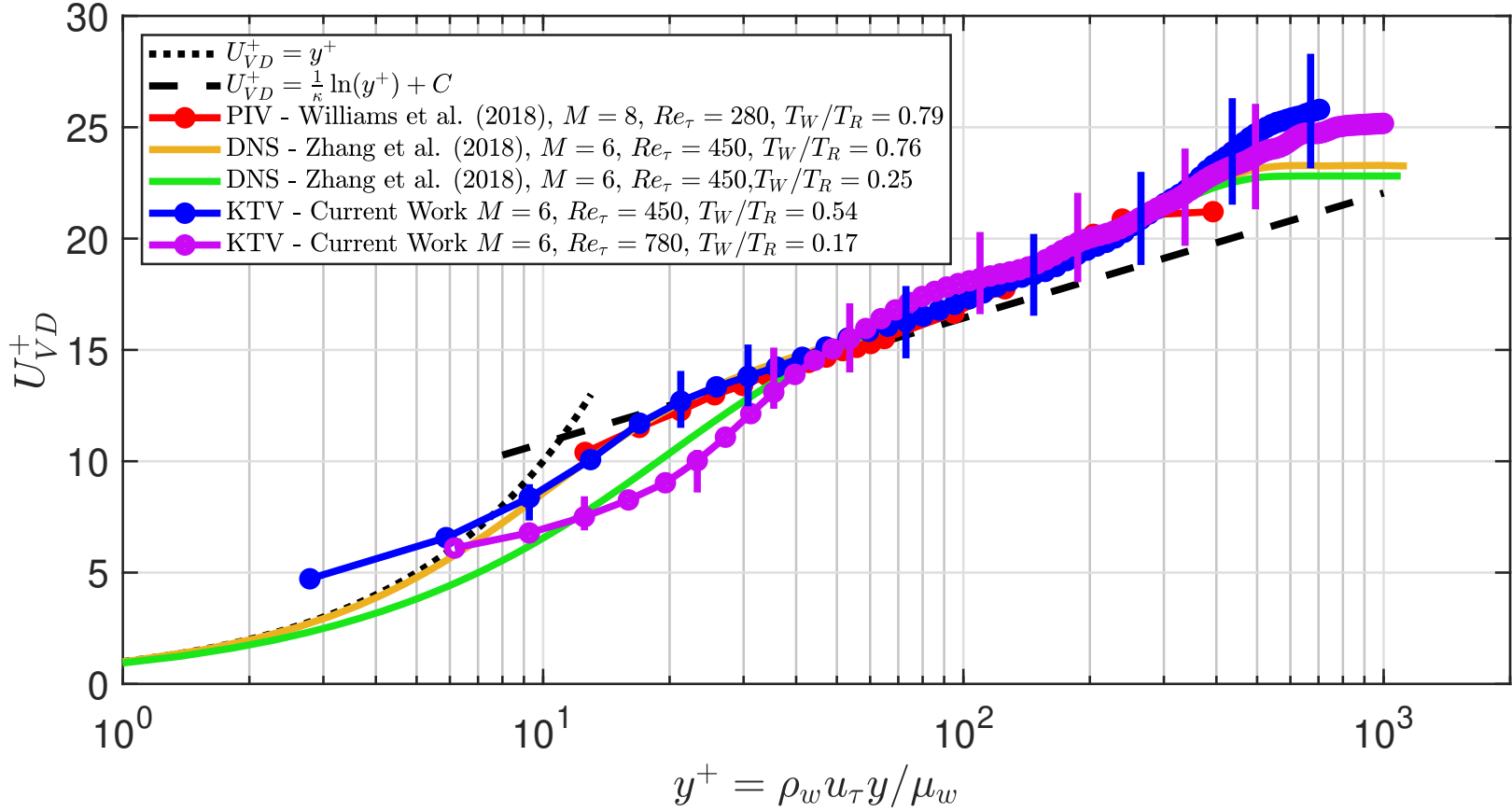


Figure 5.7: Van Driest non-dimensionalization of KTV data (current work) and comparison to DNS in Zhang et al. [4] and PIV in Williams et al. [5]. Vertical lines are error for KTV data.

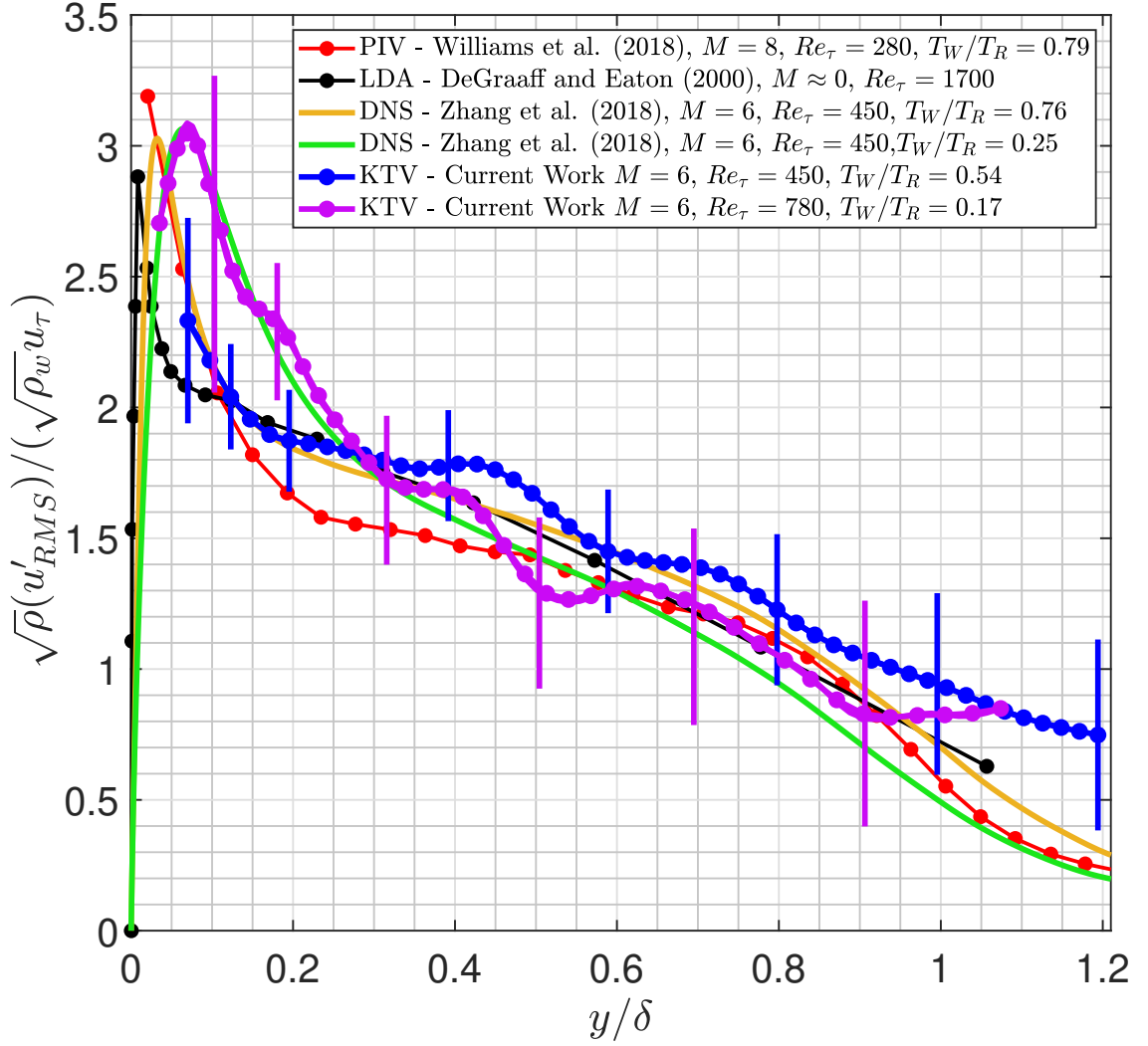


Figure 5.8: Morkovin scaled u'_{RMS} for KTV experimental data (present work) compared to DNS, LDA, and PIV.

5.3 Conclusions and Future Work

Streamwise KTV was performed in the Stevens Shock Tunnel, yielding Van-Driest transformed mean data and Morkovin scaled fluctuation data. The case above shows the nitrogen low-enthalpy, cold-flow shot and the helium driver condition, matching Mach 6 free-flight enthalpy. More on this will be discussed in Chapter 7.

Chapter 6

Wall-Normal Fluctuations Using KTV

This chapter continues experimentation and introduces a novel single-line method of KTV to measure wall-normal fluctuations. This is part of the first published KTV campaign from the SST facility and was published in Nature Communications in 2025.

6.1 Experimental Setup

Five shock tunnel runs were executed for the wall-normal configuration with both the cold-flow and enthalpy-matched condition. Each run yields ≈ 70 realizations for a total of ≈ 350 realizations per data set. All tunnel-run data are available in Appendix B.

To measure wall-normal fluctuations, an approach was implemented as shown in Fig. 6.1 (top). Juxtaposed to the streamwise setup, the 212.6 nm beam entered a window on the side of the facility, traversing the spanwise direction to a mirror just behind the flare on the HCF. Then, the beam is turned at nearly a right angle and directed upstream at a shallow angle ($\theta_{\text{RF}} = 14^\circ$ in Fig. 6.1) to the freestream to strike the top of the HCF. This tagged line captures wall-normal fluctuation information. The angle had to be shallow to reduce streamwise fluctuation contamination of the wall-normal data and to capture the entire boundary layer and a significant portion of the freestream ($y/\delta \approx 1.6$). This relatively long traversal length required the use of a 750 mm lens to provide a long Rayleigh length. The wall-normal fluctuation imaging system was a Carl Zeiss 85 mm Lens, an 8 mm lens tube, a Video Scope VS4-1845HS image intensifier (200 ns gate), and a TMX 7510 high-speed camera (12 pixel/mm

resolution). The prescribed write/read delay was 2 μ s. Sample write/read exposures are shown in Fig. 6.1 (bottom).

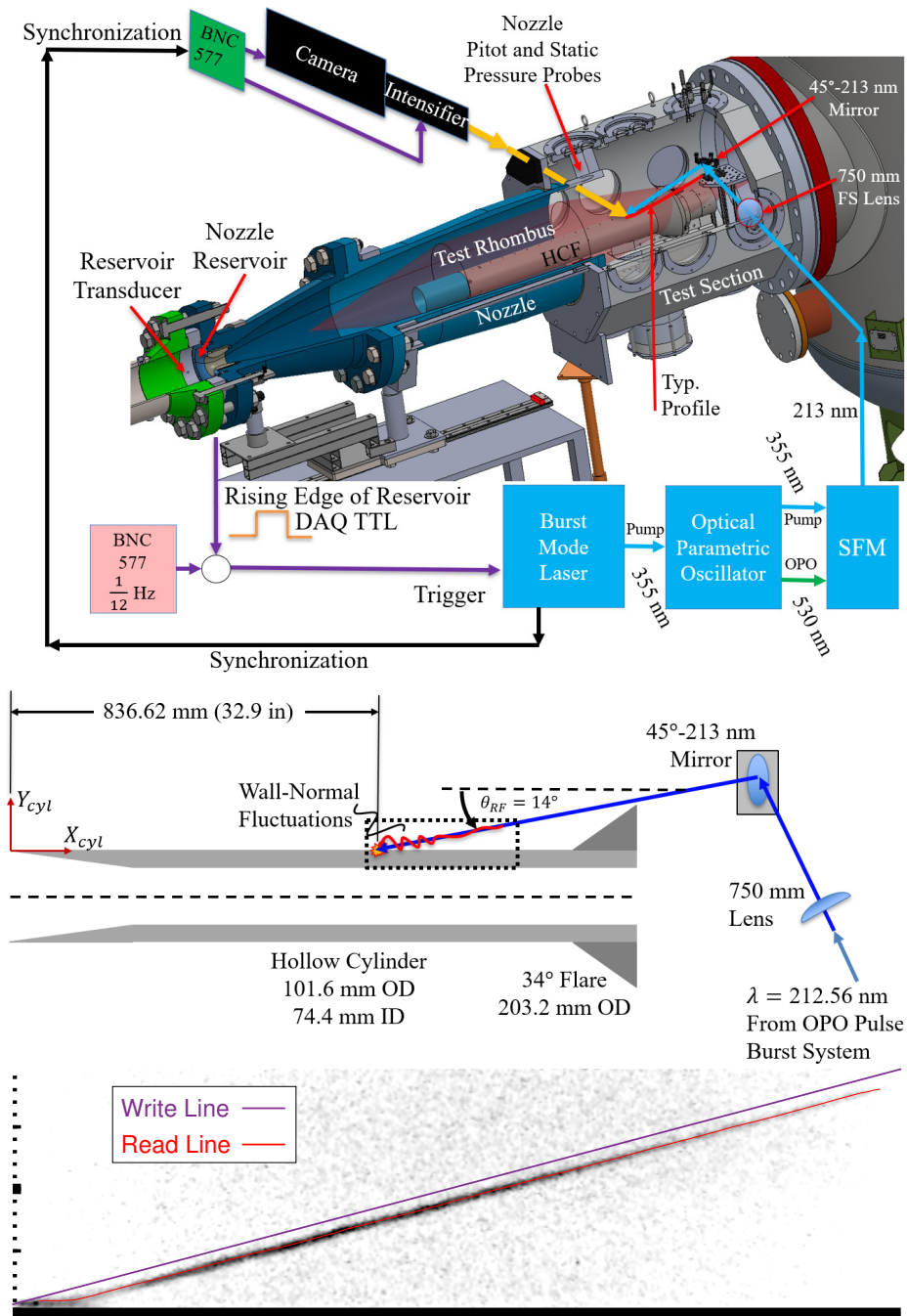


Figure 6.1: Top: Experimental setup for wall-normal fluctuations Middle: Side view of HCF and wall-normal fluctuation setup. Bottom: Sample wall-normal exposure with fitted line.

6.2 Results and Discussion

Because the gas is moving in nearly the opposite direction of the 212.6 nm beam, the absorption wavelength of the Kr will change substantively. Following Demtröder [135], this is calculated as

$$\lambda_s = \frac{\lambda_0}{1 - (U/c) \cos(\theta_{\text{RF}})} \quad (6.2.1)$$

where λ_s , λ_0 and c are the Doppler-shifted wavelength, the un-Doppler-shifted wavelength, and the speed of light, respectively.

An experiment was conducted to validate the Doppler shift calculation described in Eq. 6.2.1. Reference images were first acquired in the facility test section at 10 torr with the laser beam oriented at 14° to replicate the intended beam path. Images were captured using a TMX 7510 camera with an image intensifier at 100 kHz, with a 100 ns delay between the laser pulse and camera gate to prevent saturation. The same camera frame was used for each case to ensure direct comparison, corresponding to approximately the 15th pulse in the 30-pulse burst train of the pulse-burst laser. Following each acquisition, the OPO idler seed wavelength was incrementally tuned to observe the resulting change in fluorescence intensity due to Doppler effects. The results are shown in Fig. 6.3. The peak response under reference conditions occurs at a seed wavelength of 1072.934 nm, corresponding to excitation of the 212.556 nm krypton transition.

To visualize the Doppler shift under a known velocity, an underexpanded jet was used to generate a controlled flow field. A schematic of the experiment setup can be seen in Fig. 6.2. The jet conditions are detailed in Mustafa and Parziale [55], with a calculated Mach number of approximately 5 and a static temperature of approxi-

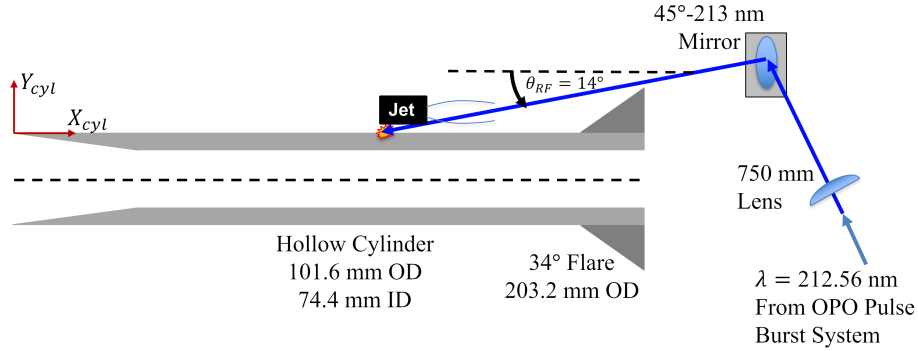


Figure 6.2: Side view of HCF and wall-normal Doppler shift setup.

mately 50 K in a 5% krypton, N_2 balance gas mixture. These conditions correspond to a freestream velocity of approximately 700 m/s. Using Eq. 6.2.1, the expected Doppler shift is a 12.3 pm redshift of the KTV line, corresponding to a 12.3 pm blueshift of the idler seed wavelength for maximum krypton emission intensity.

Images of the underexpanded jet were acquired over seed wavelengths spanning 1072.914–1072.943 nm to identify the peak response and determine the spectral limits of detectable fluorescence. As shown in Fig. 6.4, the highest intensity occurs between 1072.921 nm and 1072.924 nm, consistent with the predicted Doppler shift from Eq. 6.2.1. Detectable fluorescence with acceptable signal-to-noise ratio was observed up to approximately 15 pm from the peak seed wavelength.

Because the velocity varies throughout the boundary layer, selection of a single seed wavelength must account for the spatially varying Doppler shift. For the cold-flow condition, the mean streamwise velocity ranges from approximately 0 m/s at the wall to nearly 1 km/s at the boundary-layer edge, increasing to approximately 2 km/s in the “enthalpy-matched” condition. A representative Doppler shift profile for the cold-flow case is shown in Fig. 6.5, computed using the local mean velocities from Shot 385 (see Chapter 5).

For the “cold-flow” condition, the predicted Doppler shift spans up to approx-

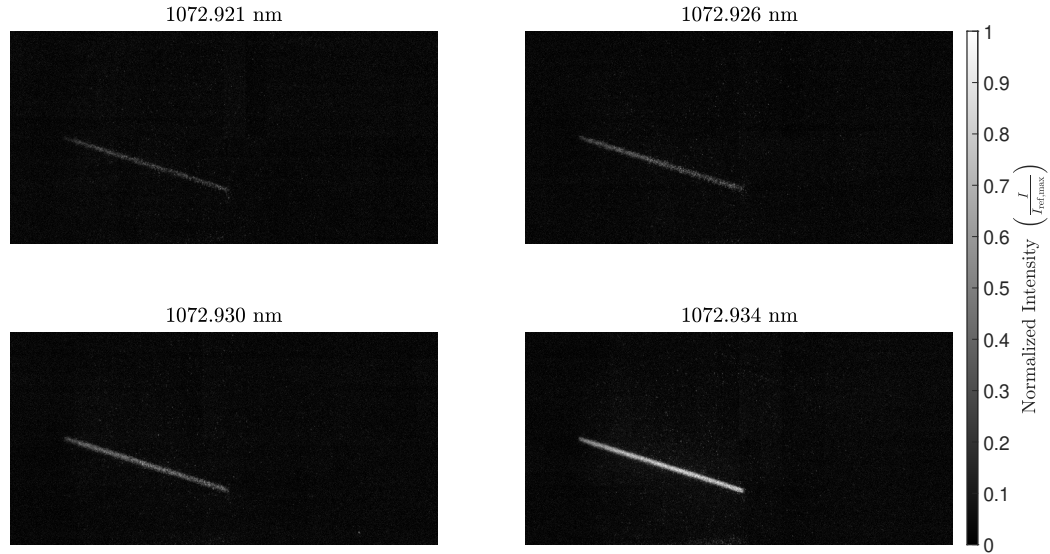


Figure 6.3: Fluorescence images acquired at seed wavelengths of 1072.921–1072.934 nm. All images are normalized to the maximum pixel intensity in the 1072.934 nm case to enable direct comparison of signal strength across wavelengths. The decrease in intensity with wavelength confirms the peak seed wavelength with no opposing flow is 1072.934 nm.

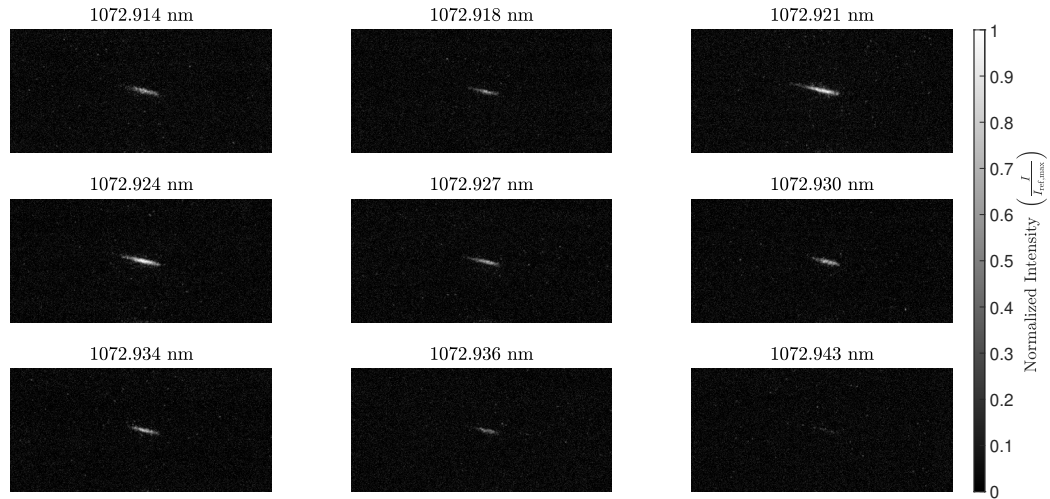


Figure 6.4: Fluorescence images acquired at seed wavelengths ranging from 1072.914 to 1072.943 nm. All images are normalized by the maximum pixel intensity from the 1072.921 nm condition to enable direct comparison of signal strength across wavelengths. The highest intensities are observed for the 1072.921 and 1072.924 nm conditions, consistent with the shift predicted by Eq. 6.2.1.

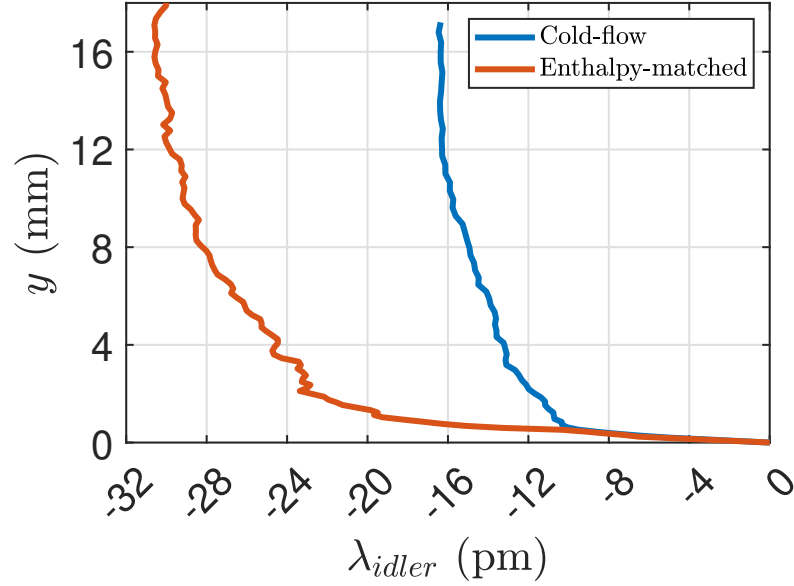


Figure 6.5: Doppler shift throughout the boundary layer for both the cold-flow and enthalpy-matched run condition. The mean streamwise velocity profiles for each condition were substituted into Eq. 6.2.1 to determine the corresponding shift in OPO idler wavelength.

imately 16 pm across the boundary layer. Accordingly, the idler seed wavelength was shifted by 16 pm to maximize fluorescence intensity throughout the Rayleigh length of the beam. Tests conducted with smaller wavelength shifts resulted in substantially reduced signal levels compared to the 16–18 pm shifted cases.

For the “enthalpy-matched” condition, the freestream velocity corresponds to a predicted Doppler shift of approximately 30 pm. However, to maintain strong fluorescence throughout the boundary layer while preserving signal uniformity, a shift of 20 pm was selected as a compromise between freestream resonance and near-wall signal strength.

A novel processing algorithm was used to quantify the wall-normal fluctuations, v'_{RMS} . These fluctuations may be found through the Reynolds decomposition, $v = \bar{v} + v'$; that is, the mean will be used as a reference to find the fluctuating quantity. The process is as follows:

- (1) Referring to Fig. 6.6-top, in physical space (X_{cyl}, Y_{cyl}) , fit a straight line (X_F, Y_F) to the read-line data (X_R, Y_R) outside the boundary layer at $y/\delta = 1.0 - 1.6$ and extend that line to the model surface (green line in Fig. 6.6-top);
- 2) Find the difference in the wall-normal direction between the read data and the freestream fitted line to form $\Delta Y_1 = Y_R - Y_F$;
- 3) Recast the data in $(Y_{cyl}, \Delta Y_1)$ space as shown in Fig. 6.6-bottom;
- 4) Find the mean of ΔY_1 to form $\overline{\Delta Y_1}$ shown in Fig. 6.7;
- 5) Subtract $\overline{\Delta Y_1}$ from ΔY_1 to form $\Delta Y_2 = \Delta Y_1 - \overline{\Delta Y_1}$, as shown in Fig. 6.8;
- 6) Find the RMS of ΔY_2 to form $\Delta Y_{2\text{RMS}}$, shown in black in Fig 6.8 and divide by the interframe delay to form $v'_{\text{RMS}} = \Delta Y_{2\text{RMS}}/\Delta t$.

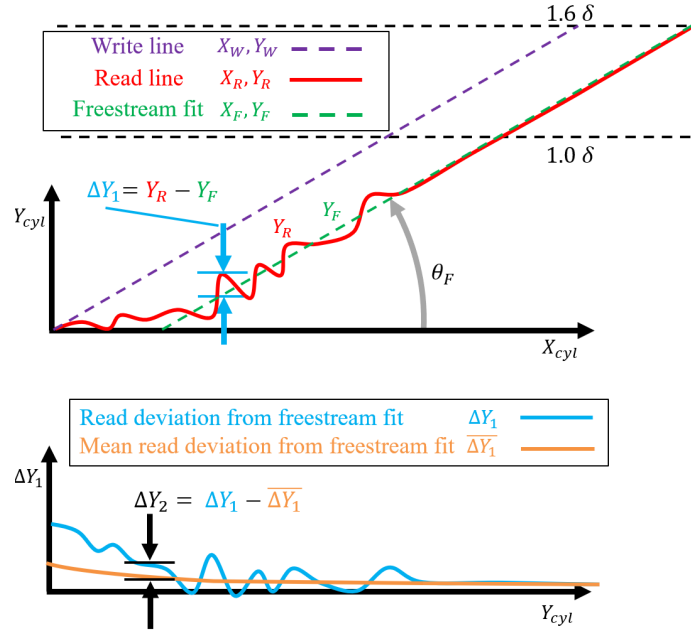


Figure 6.6: Approach for wall-normal velocity fluctuation measurement. Top: Sketch of write line in purple, translated read line in red, and fit to read line above boundary layer in blue, $y/\delta = 1.0 - 1.6$. ΔY_1 Bottom: Sketch of read deviation from freestream fit in blue, and mean read deviation from freestream fit in gold.

Exposures for the wall-normal fluctuations were processed nearly-identically to the streamwise data. The images were rotated such that the line was forced to

be nearly vertical to help the Gaussian peak finding algorithm and the images were unwarped. Because the fluorescing line extends the entire length of the camera lens, from bottom left to top right, image warping became appreciable. To rectify this, images were unwarped before analysis using the `imwarp` MATLAB command and prescribing a displacement. The displacement was quantified by:

- 1) Capturing a write line to be used as a reference (with no flow);
- 2) Rotating the reference line such that it is vertical;
- 3) Performing the binning, truncation and peak fitting to the reference line as done in the streamwise data;
- 4) Performing a linear fit to the reference line;
- 5) Record the displacements in x between the reference line and linearly-fitted line for every y .

A sanity check on this unwarping procedure was to plot the reference write line, which can be assumed to be straight, before and after applying `imwarp` in MATLAB. Because of the differential nature of the technique, this unwarping procedure was probably not necessary but for the presentation of ΔY_1 and $\overline{\Delta Y_1}$ in Fig. 6.7. That is, the results in ΔY_2 in Fig. 6.8 do not change appreciably relative to the error in the measurement with or without the unwarping procedure.

Uncertainty in the wall-normal KTV measurements is estimated in the usual way as

$$\widetilde{v'_{\text{RMS}}} = \sqrt{\left(\widetilde{\Delta t} \frac{\partial v'_{\text{RMS}}}{\partial \Delta t}\right)^2 + \left(\widetilde{\Delta Y} \frac{\partial v'_{\text{RMS}}}{\partial \Delta Y}\right)^2} + v'_{\text{RMS-Bias}}, \quad (6.2.2)$$

where the first two terms consider precision error and the final term considers bias error. The uncertainty in the measured displacement distance, $\widetilde{\Delta Y}$, of the tracer is estimated as 10 μm (from Gaussian fitting). The uncertainty in time, $\widetilde{\Delta t}$, is estimated to be 100 ns (half the intensifier gate width).

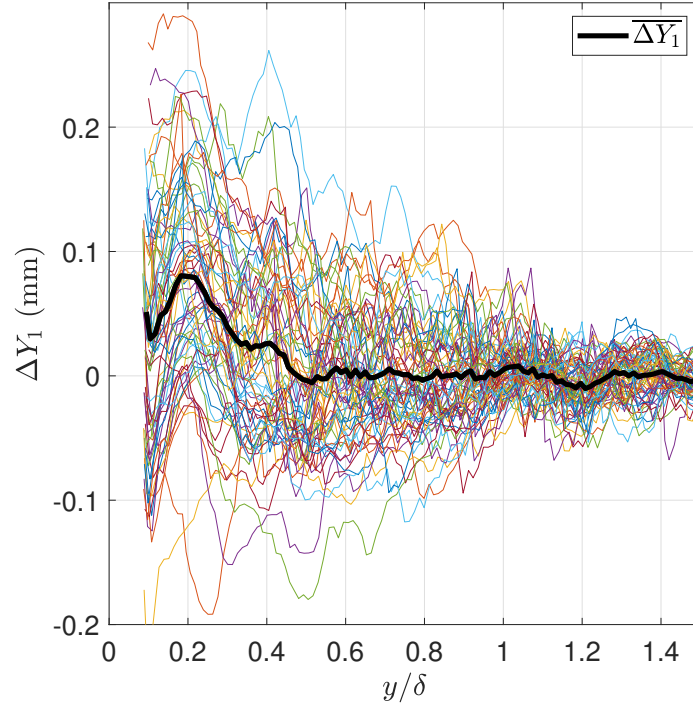


Figure 6.7: Data visualization of Fig 6.6-bottom for all realizations in one shock tunnel run.

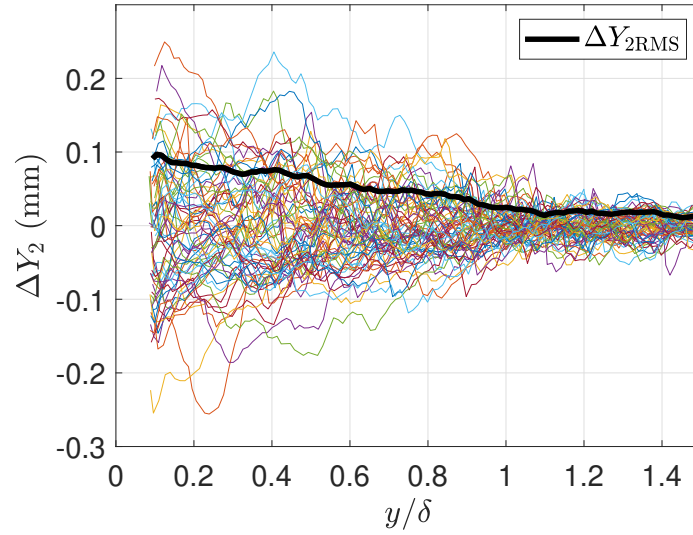


Figure 6.8: Data visualization of $\Delta Y_{2\text{RMS}}$ for all realizations in one shock tunnel run.

The last term, $v'_{\text{RMS-Bias}}$, accounts for the uncertainty due to the potential in-plane contamination from $\bar{u}$, u' , and $\bar{v}$. Two strategies were used to assess the in-

plane contamination. The first approach is to analyze the magnitude of the measured wall-normal velocity fluctuation as

$$v'_{\text{RMS-Measured}} = \frac{\Delta Y_{2\text{-RMS}}}{\Delta t} = \sqrt{(\delta v'_{u'})^2 + (v'_{\text{DNS}})^2}, \quad (6.2.3)$$

where the wall-normal velocity fluctuation from Zhang et al. [4] DNS is v'_{DNS} and the uncertainty in the wall-normal fluctuation due to the streamwise fluctuation, by geometry, is $\delta v'_{u'} = u'_{\text{DNS}} \tan \theta_F$. The second approach is to create and analyze pseudo data via a Monte Carlo (MC) simulation of the kinematics of the write/read line as

$$\begin{aligned} X_R &= X_W + \underline{u}_{\text{DNS}} \Delta t = (\bar{u}_{\text{DNS}} + u'_{\text{DNS}}) \Delta t \\ Y_R &= Y_W + \underline{v}_{\text{DNS}} \Delta t = (\bar{v}_{\text{DNS}} + v'_{\text{DNS}}) \Delta t. \end{aligned} \quad (6.2.4)$$

The velocity vectors in Eqs. 6.2.4 are created from Zhang et al. [4], assuming 300 normally distributed random samples for the fluctuations. The pseudo data were processed identically to the experimental data.

Processing the Monte Carlo (MC) pseudo data in the same way as the experimental data results in a nearly identical v'_{RMS} as that from the analytical prediction in Eq. 6.2.3, bringing confidence to both (Fig. 6.9 (left)). Additionally, in Fig. 6.9 (left), we evaluate the response of the measurement technique to turning on/off different components of velocity. Including all components, $(\bar{u}, u', \bar{v}, v')$, most closely mimics the present experiment and was used to evaluate the bias error in Eq. 6.2.2, as there is a slight positive bias (plotted in red). Because of the differencing step in the measurement technique, there is no error in v'_{RMS} from $(\bar{u}, \bar{v})$ (plotted in blue). The primary form of bias error is then from the (u') (purple). Finally, a sanity check including only $(\bar{v}, v')$ matches the prescribed DNS data nearly identically (plotted in green).

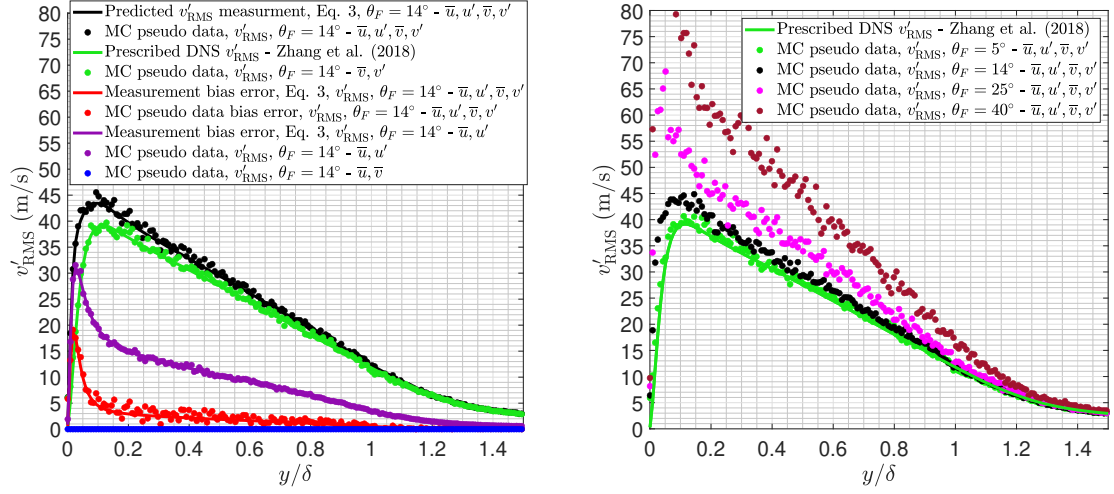


Figure 6.9: Bias error assessment. Left: Response of the measurement technique to turning on/off different velocity components to assess their effect on $v'_{\text{RMS-Bias}}$. Data labeled with a component include only those components; e.g., data labeled with $\bar{u}, u'$ includes only those two components. Right: Response of the measurement technique to variable incident angle θ_F

Fig. 6.9 (right) shows that the bias error of the technique is highly dependent on the incident angle, θ_F . The bias error is low outside the boundary layer and increases closer to the wall due to the onset of streamwise velocity fluctuation contamination in wall-normal measurements. This is apparent in the bias error increasing in Fig. 6.9 (left) nearer to the wall. The data appear to nearly collapse at $\theta_F = 5^\circ$. Ultimately, for the current experimental setup with $\theta_F = 14^\circ$, the maximum bias error in v'_{RMS} is at most $\approx 8\%$ high at $y/\delta \approx 0.1$. We assert this bias error is acceptable, given the dearth of data for this flowfield, and report error bars in the final results per Eqs. 6.2.2 and 6.2.3.

The wall-normal (v'_{RMS}) fluctuation data are shown in Fig. 6.10 as normalized by the Morkovin scaling for computations (DNS), comparable experiments (PIV), and the present work (KTV). Additionally, the LDA data from DeGraaff and Eaton [6] is shown as a point of comparison for incompressible flow. That is, the fluctuation

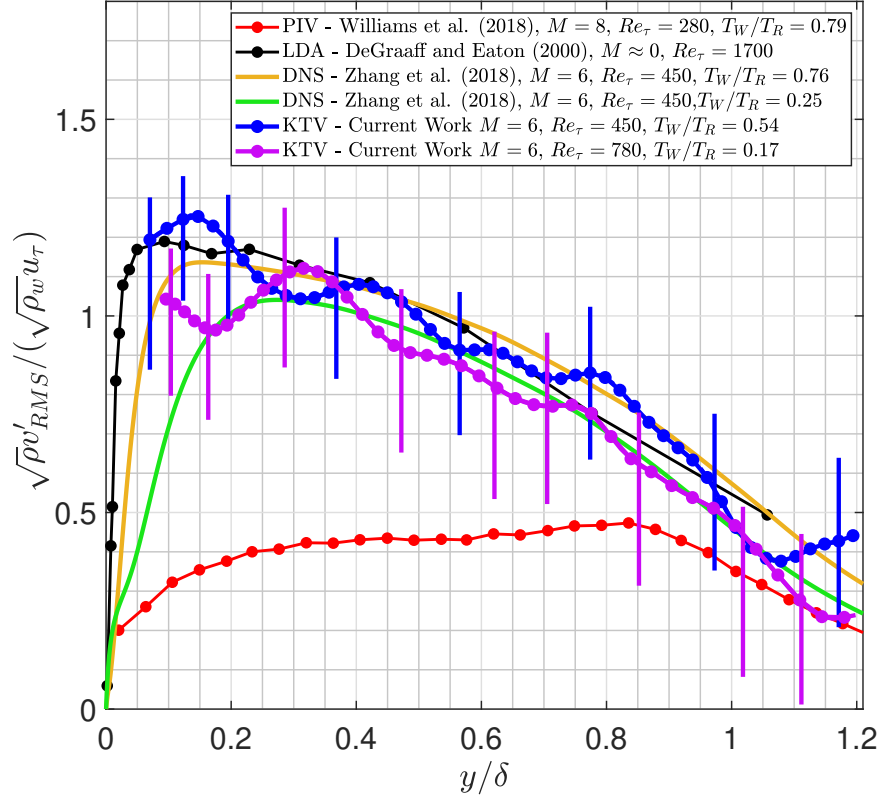


Figure 6.10: Fluctuation comparison. Morkovin scaled v'_{RMS} for KTV experimental data (present work plotted every fifth point) compared to DNS, LDA, and PIV (plotted every other data point).

profiles may be transformed from the compressible form to a similar incompressible form by simply accounting for mean density variations across the boundary layer [4]. The nondimensional mean and fluctuation profiles were not found to change substantially ($< 2\%$) when the mean density variation was evaluated with the Walz [123] relation and compared to the generalized Reynolds analogy in Zhang et al. [25].

The data are presented in the outer units (against y/δ) where sensitivity to Reynolds number differences between data sets would be most modest; this is because the fluctuations could be expected to collapse with Reynolds number as in incompressible flow. Closer to the wall, the KTV bias error increases, per Chapter 5. The DNS, PIV, and KTV streamwise fluctuation data are in reasonable agreement

with the incompressible data when appropriately transformed. However, it is apparent that PIV, which is the current state-of-the-art (SOA) measurement technique, cannot capture the wall-normal fluctuations at hypersonic conditions. Williams et al. [35] suggests that this discrepancy is because PIV is fundamentally limited by insufficient particle response times due to rarefaction effects, especially near the wall.

In contrast, because of the use of a tagging technique which does not suffer from reduced particle response, the wall-normal KTV fluctuation data in Fig. 6.10 show good agreement with DNS from Zhang et al. [4] performed at two wall-temperature ratios. DNS predicts that the wall-temperature ratio will shift the peak magnitude and location slightly at wall-normal locations, a feature apparently captured in the $T_W/T_R = 0.17$ case. The transformed DNS and KTV data agree with the incompressible data from DeGraaff and Eaton [6]. These KTV data are the first such measurements of hypersonic wall-normal fluctuations in support of Morkovin’s hypothesis [3]. There is a slight undulation in the KTV data with a wavelength of $y/\delta \approx 0.2$. This feature is not believed to be physical and falls well within the error bars of the measurement. KTV also indicates a higher level of noise outside the boundary layer than in the DNS, a topic of future investigation.

6.3 Conclusions and Future Work

Wall-normal KTV was performed in the Stevens Shock Tunnel, yielding Morkovin scaled fluctuation data, for both the cold-flow and enthalpy-matched conditions, that shows good agreement to the DNS. These are the first successful experimental measurements of v'_{RMS} at hypersonic Mach number. Future work will include expanding the run conditions for varying T_W/T_R and Re_τ .

Chapter 7

Turbulence Results and Discussion

This chapter presents the analysis and interpretation of turbulence measurements obtained from optical diagnostics in a Mach 6 turbulent boundary layer. Building on the experimental methodologies developed in previous chapters, the results are examined in the context of established theoretical frameworks and high-fidelity numerical simulations to assess the accuracy of the measurements and the underlying flow physics.

The focus of this chapter is on characterizing the mean and fluctuating components of the velocity field, evaluating compressibility effects through appropriate transformations, and examining the structure of turbulence in hypersonic conditions. Comparisons to DNS and prior experimental data are used to validate the measurements and to assess the applicability of Morkovin's hypothesis. In addition, turbulence structure is explored through the identification of coherent features such as uniform momentum zones (UMZs), and through the estimation of characteristic turbulence length and time scales.

7.1 Turbulence Statistics from MTV Campaigns

An extensive analysis of the data measured in previous chapters is needed to characterize the flow. First, the Van Driest transformed mean streamwise velocities are compared to experiments with similar conditions as well as the DNS from Zhang et al. [4]. The mean streamwise velocity of the enthalpy-matched condition, detailed in Chapter 5 was compared to the Zhang et al. data for similar conditions, Mach 6 flow with $Re_\tau = 450$ and a recovery ratio $T_w/T_r = 0.25$. This data set was the closest DNS

match in terms of a cold-wall boundary-layer flow. As mentioned earlier, the mean velocities are nondimensionalized via the van Driest [134] transformation following [23, 34] with the local mean density profile coming from the generalized Reynolds analogy from Zhang et al. [25], as opposed to the simpler Walz [123] relation. The skin-friction coefficient is evaluated following Zhao and Fu [36], replacing the previous technique of using the Kármán–Schoenherr correlation and the van Driest II [134] transformation. The Zhao and Fu [36] method is chosen, as the van Driest II method struggles to map the compressible scaling of skin-friction to the incompressible relation for a highly cooled wall, as confirmed by Huang [34] and Zhao and Fu [36]. As seen in Figure 7.1, the data agrees with the DNS within the uncertainty of the experiment including similar capturing of the buffer region and viscous subregion in the law of the wall. The ATV experiments, detailed in Chapter 4 were conducted at the cold-flow condition, matching the KTV experiments for the same condition. These are compared to the Zhang et al. data for similar conditions, Mach 6 flow with $\text{Re}_\tau = 450$ and a recovery ratio $T_w/T_r = 0.76$.

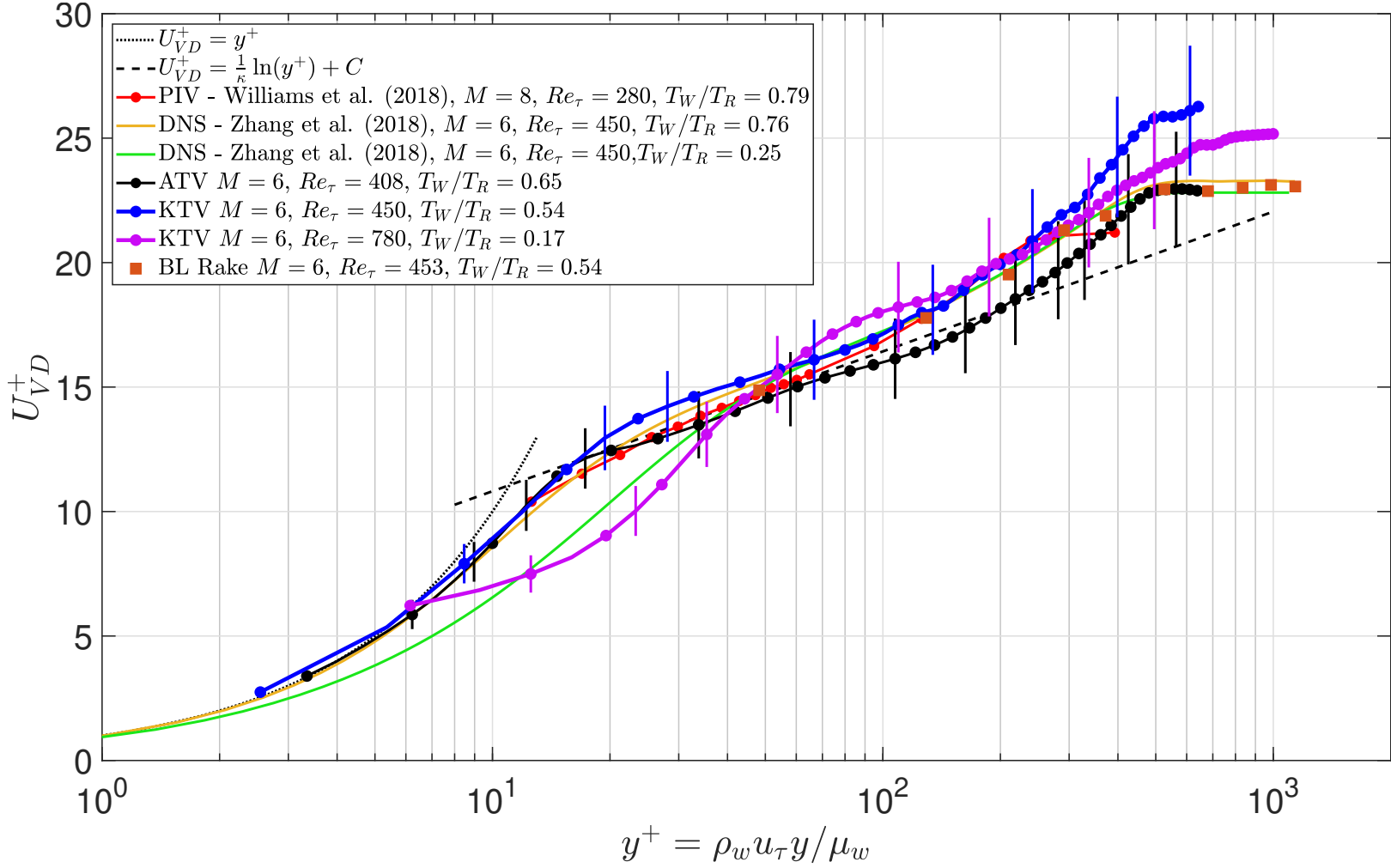


Figure 7.1: Mean velocity comparison. van Driest non-dimensionalization of KTV data (present work) and comparison to DNS in Zhang et al. [4] and PIV in Williams et al. [5]. Vertical lines are error bars for KTV data.

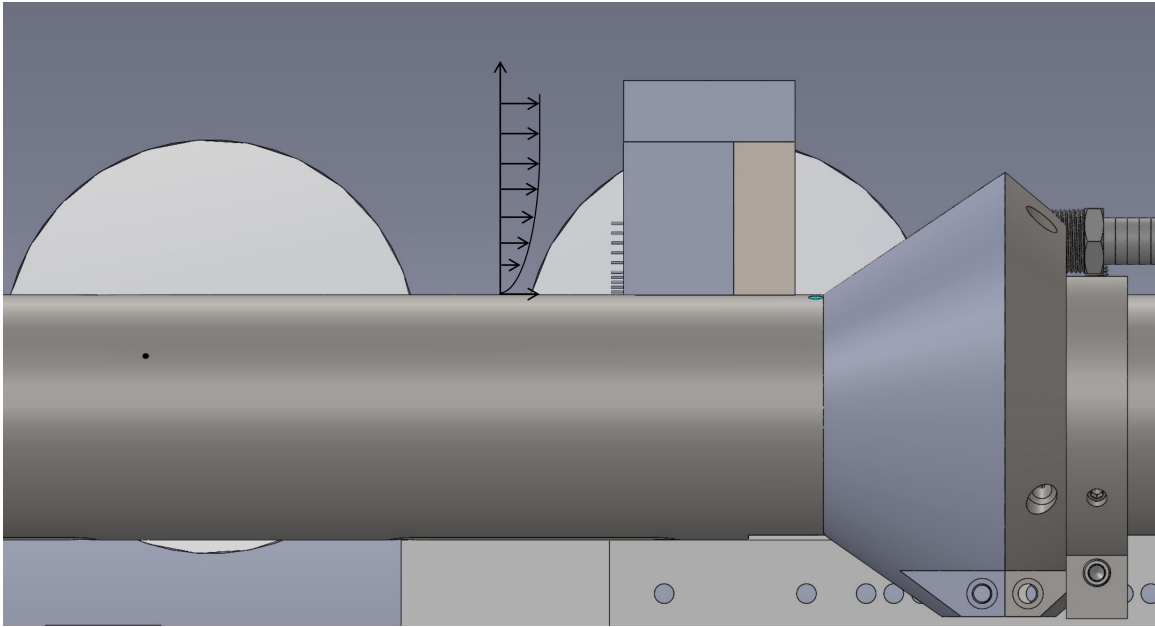


Figure 7.2: Probed boundary-layer rake schematic in SST. Flow goes from left to right.

Additionally, a probed-boundary layer rake was used to make turbulence measurements on the surface of the hollow-cylinder test article. The design is adapted from a design from Bui et al. [136] utilizing thin-walled tubes and pressure transducers (PCB Piezotronics Model 113B28) to record pitot pressures throughout the boundary layer. The rake was mounted at the same axial location downstream of the hollow-cylinder leading edge and included ten measurements at wall-normal values of 1.27 mm, 3.40 mm, 5.53 mm, 7.67 mm, 9.80 mm, 13.86 mm, 17.93 mm, 21.99 mm, 26 mm, 30 mm. An important design consideration for the rake is ensuring the filling time of the cavity is fast enough for the transducer to record measurements during the test time. The equation for the time constant is defined from “Response of Pressure Tap” by H.G. Hornung as $\tau = 5V_{tube}/(A_{tube} * a)$, where V is the volume of the tube and A is the cross sectional area of the tube. The mounted rake on the test article can be seen in Fig. 7.2.

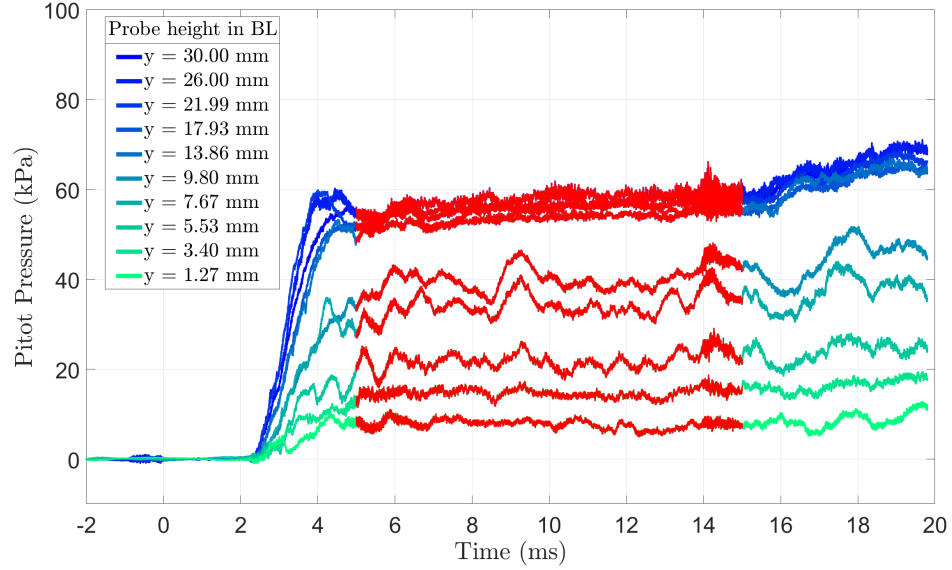


Figure 7.3: Boundary-layer rake pitot pressure for every probe height. Steady test time is highlighted in red.

To compute turbulence quantities from the probe, the pitot pressures at each rake location were first time-averaged over the useful test window to obtain a mean pressure distribution across the boundary layer, shown in Fig. 7.3. The traces also give a range of the boundary layer thickness, as the pitot pressure increases to the freestream pitot pressure between the 9.80 mm and 13.86 mm probes. The mean pitot pressure during the test time for each probe was referenced against the static pressure which was then used to infer the local Mach number using the Rayleigh pitot relation. The local Mach number profile was then used to compute the corresponding temperature distribution, assuming a known reservoir temperature. The local speed of sound was then evaluated from the temperatures, and the velocity profile was obtained as $u(y) = M(y) a(y)$. The outer-edge velocity, U_E , was estimated from the uppermost rake locations, where the flow approaches freestream conditions. Comparisons between the DNS from Zhang et al. [4] and the boundary-layer rake in both local temperature and velocity, non-dimensionalized by the wall temperature

and edge velocity respectively, are shown in Fig. 7.4. The profiles from the rake show good agreement in value and overall shape compared to the DNS.

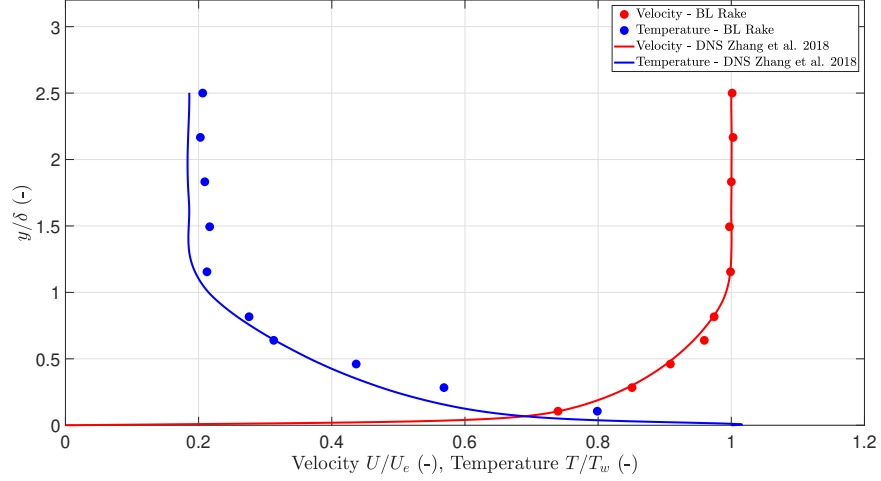


Figure 7.4: Boundary-layer rake non-dimensionalized velocity (U/U_E) temperature (T/T_w) and compared to DNS data from Zhang et al. [4].

With these values calculated, the momentum thickness, friction velocity, skin friction and Van Driest transformed velocities are all computed in the same manner as the tagging experiments. The transformed mean velocities for the cold-flow condition are plotted with the tagging experiments on Fig 7.1 and show good agreement with both the experiments and the DNS. Tests were conducted for the enthalpy-matched condition but the time constant for the cavity filling were too slow to capture the steady test time for the facility.

The streamwise (u'_{RMS}) and wall-normal (v'_{RMS}) fluctuation data are shown in Fig. 7.5 as normalized by the Morkovin scaling for computations (DNS), comparable experiments (PIV), and the present work (KTV). Additionally, the LDA data from DeGraaff and Eaton [6] is shown as a point of comparison for incompressible flow. That is, the fluctuation profiles may be transformed from the compressible form to a similar incompressible form by simply accounting for mean density variations across the boundary layer [4]. The nondimensional mean and fluctuation profiles were not

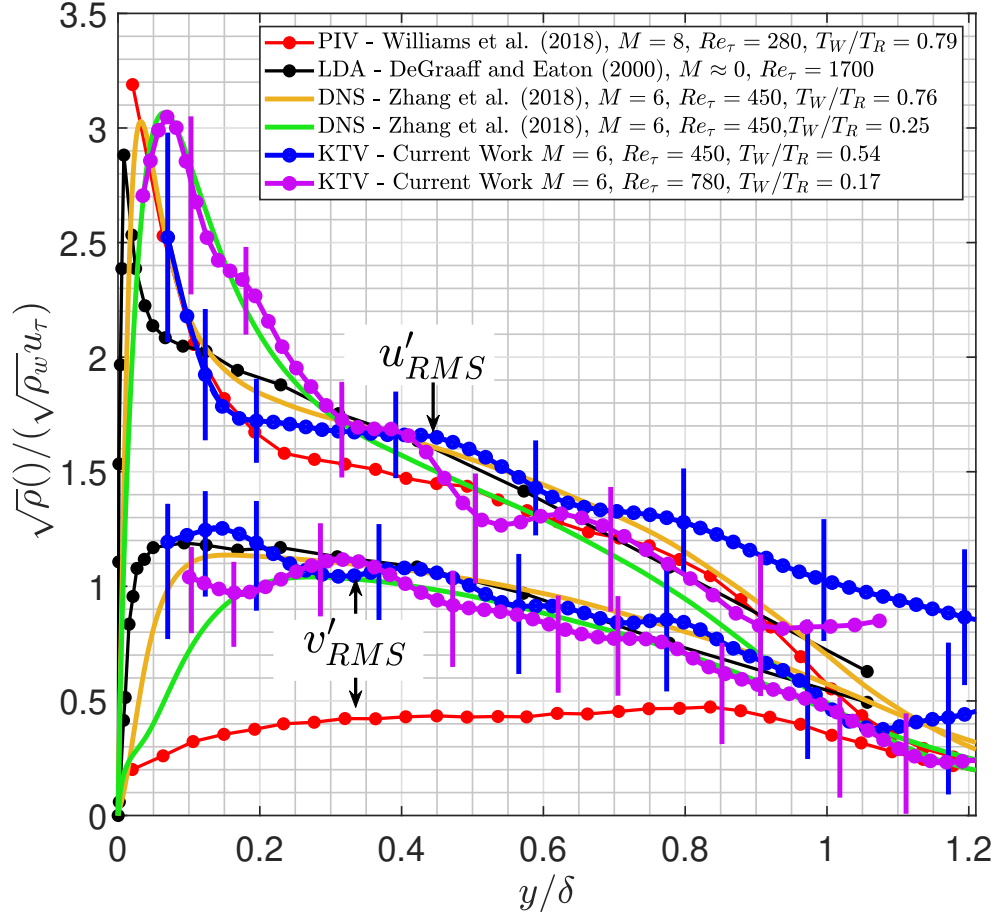


Figure 7.5: Fluctuation comparison. Morkovin scaled u'_{RMS} and v'_{RMS} for KTV experimental data (present work plotted every fifth point) compared to DNS, LDA, and PIV (plotted every other data point).

found to change substantially ($< 2\%$) when the mean density variation was evaluated with the Walz [123] relation and compared to the generalized Reynolds analogy in Zhang et al. [25].

The data are presented in the outer units (against y/δ) where sensitivity to Reynolds number differences between data sets would be most modest; this is because the fluctuations could be expected to collapse with Reynolds number as in incompressible flow. Closer to the wall, the KTV bias error increases, per Chapter 5. The DNS, PIV, and KTV streamwise fluctuation data are in reasonable agreement

with the incompressible data when appropriately transformed. However, it is apparent that PIV, which is the current SOA measurement technique, cannot capture the wall-normal fluctuations at hypersonic conditions. Williams et al. [35] postulate that this discrepancy is because PIV is fundamentally limited by insufficient particle response times due to rarefaction effects, especially near the wall.

In contrast, because of the use of a tagging technique which does not suffer from reduced particle response, both the streamwise and wall-normal KTV fluctuation data in Fig. 7.5 show good agreement with DNS from Zhang et al. [4] performed at two wall-temperature ratios. DNS predicts that the wall-temperature ratio will shift the peak magnitude and location slightly at wall-normal locations below $y/\delta \approx 0.1$, a feature apparently captured in the u' data for the $T_W/T_R = 0.17$ case. The fluctuation data plotted in outer units nominally collapse for $y/\delta \gtrsim 0.1$, as in Duan et al. [27]. The transformed DNS and KTV data agree with the incompressible data from DeGraaff and Eaton [6]. These KTV data are the first such measurements of hypersonic wall-normal fluctuations in support of Morkovin's hypothesis [3]. There is a slight undulation in the KTV data with a wavelength of $y/\delta \approx 0.2$. This feature is not believed to be physical and falls well within the error bars of the measurement. KTV also indicates a higher level of noise outside the boundary layer than in the DNS, a topic of future investigation.

Figure 7.6 shows the streamwise fluctuation u'_{RMS} data plotted against inner units and compared to the DNS data. The experimental data were plotted with inner-scale units. Unlike the outer-scaled streamwise fluctuation profiles, which collapsed well between the KTV measurements and the DNS, the inner scaling collapse deteriorates for the higher- Re_τ discrepancy for the enthalpy-matched condition. This behavior is expected because increasing Re_τ maps the same physical measurement locations to larger values of y^+ . This behavior is similar to that reported by Zhang

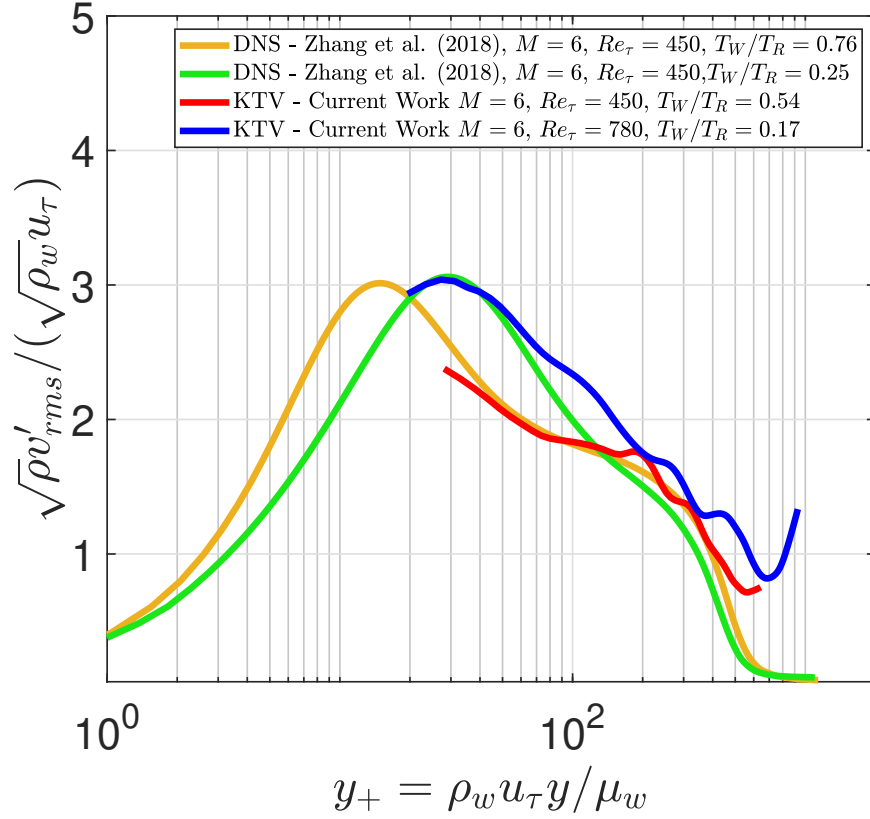


Figure 7.6: Comparison between DNS and KTV experiment Morkovin scaled u'_{RMS} with inner-scaling.

et al. [137]. In their DNS database, the turbulence intensities showed reasonable collapse when plotted using outer scaling. However, when plotted in classical inner coordinates, the peak locations shifted away from the wall as wall-cooling and Reynolds-number effects became more significant.

7.2 Turbulent Structures and Characterization of Uniform Momentum Zones (UMZ)

Uniform momentum zones (UMZs) provide a useful framework for examining the organization of coherent structures in turbulent flows. These zones, detailed by Adrian et al. [138], are defined as regions of approximately uniform streamwise velocity,

are closely associated with the presence of vortical structures such as hairpin and streamwise vortices, which play a central role in turbulence production and dissipation. Previous experimental studies have demonstrated the presence of UMZs using two-dimensional particle image velocimetry (PIV) [9, 139] for a large number of velocity fields in the boundary layer. Using the streamwise profile measured by KTV, turbulent structures and UMZs are assessed.

The KTV image capturing technique provides approximately 70 successive instantaneous velocity profiles per each shock-tunnel run. Each realization is aggregated to create the mean velocity profile through the boundary layer but the instantaneous realizations can help identify and interpret the structures that are advecting through the flow in time. Figure 7.7 shows an individual streamwise realization compared to the mean profile. Throughout the logarithmic region of the boundary layer, undershoots and overshoots from the mean flow are experienced. Although definitive conclusions cannot be made from these fluctuations, the clockwise pattern of the deviations could indicate a structure, such as a hairpin vortex. As detailed by Adrian et al. [138], these events can eject fluid from the viscous boundary layer at the wall in streaks.

Although KTV experimentation can only capture flow in one-dimension at a time, the two-dimensional spatial field is reconstructed from time-resolved streamwise measurements using Taylor's frozen turbulence hypothesis, wherein temporal variations at a fixed location are converted to spatial variations via the convection velocity, $x = U_E t$. This approach assumes that turbulent structures are advected downstream with negligible temporal evolution over the sampling interval. The time between realizations is the time between laser pulses, 10 μ s for a sampling rate 100 kHz. This sampling period is comparable to the characteristic convective time scale of the boundary layer, $\delta/U_e \approx 10.5 \mu$ s. Although the 10 μ s sampling period does

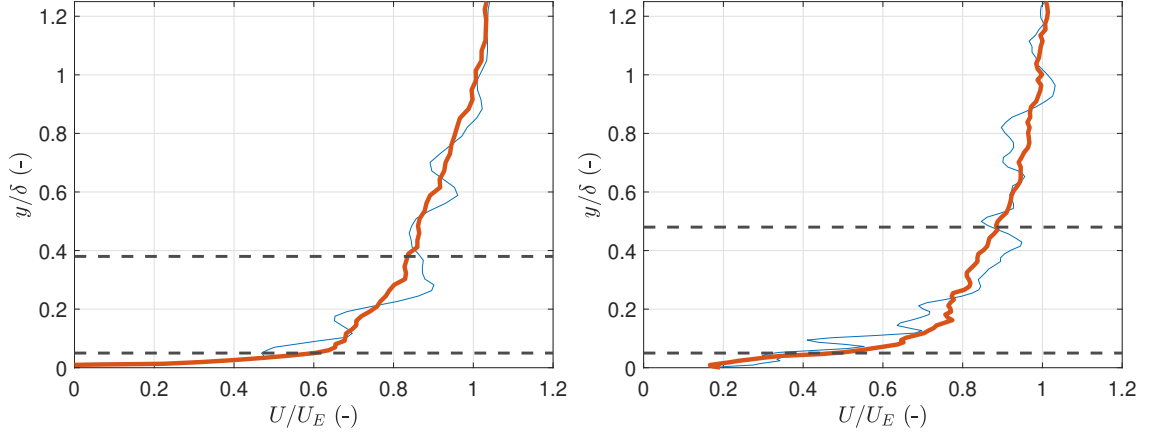


Figure 7.7: Instantaneous non-dimensionalized streamwise realization (U/U_E) compared to mean profile for both cold-flow and enthalpy-matched conditions. Black dashed lines represent the heights of UMZ regions determined in Fig. 7.8.

not strictly satisfy the Nyquist criterion for fully resolving motions at this time scale, it provides sufficient temporal resolution to capture the primary convective evolution of the large-scale flow structures. Therefore, the present measurements should be interpreted as resolving the dominant convective behavior rather than the full temporal spectrum of turbulent fluctuations. Taylor's frozen-flow hypothesis is applied to map the temporal data onto a two-dimensional spatial frame. However, because the sampling frequency is near the characteristic convective frequency, the resolution of small-scale turbulent fluctuations is limited, and the comparison is best suited for analyzing large-scale features and mean flow properties. Figure 7.8 shows the 2-D reconstruction of the KTV data with (a-d) showing contours of normalized streamwise velocity U/U_E and turbulence intensity $\log_{10}(2k/U_E^2)$ respectively for both the cold-flow and enthalpy-matched run conditions. In (a-b), for both run conditions, the flow exhibits clear stratification into regions of relatively uniform streamwise velocity, separated by sharp gradients that define the UMZ interfaces. These interfaces, indicated by the overlaid contours, display significant spatial variability and intermittent step-like behavior.

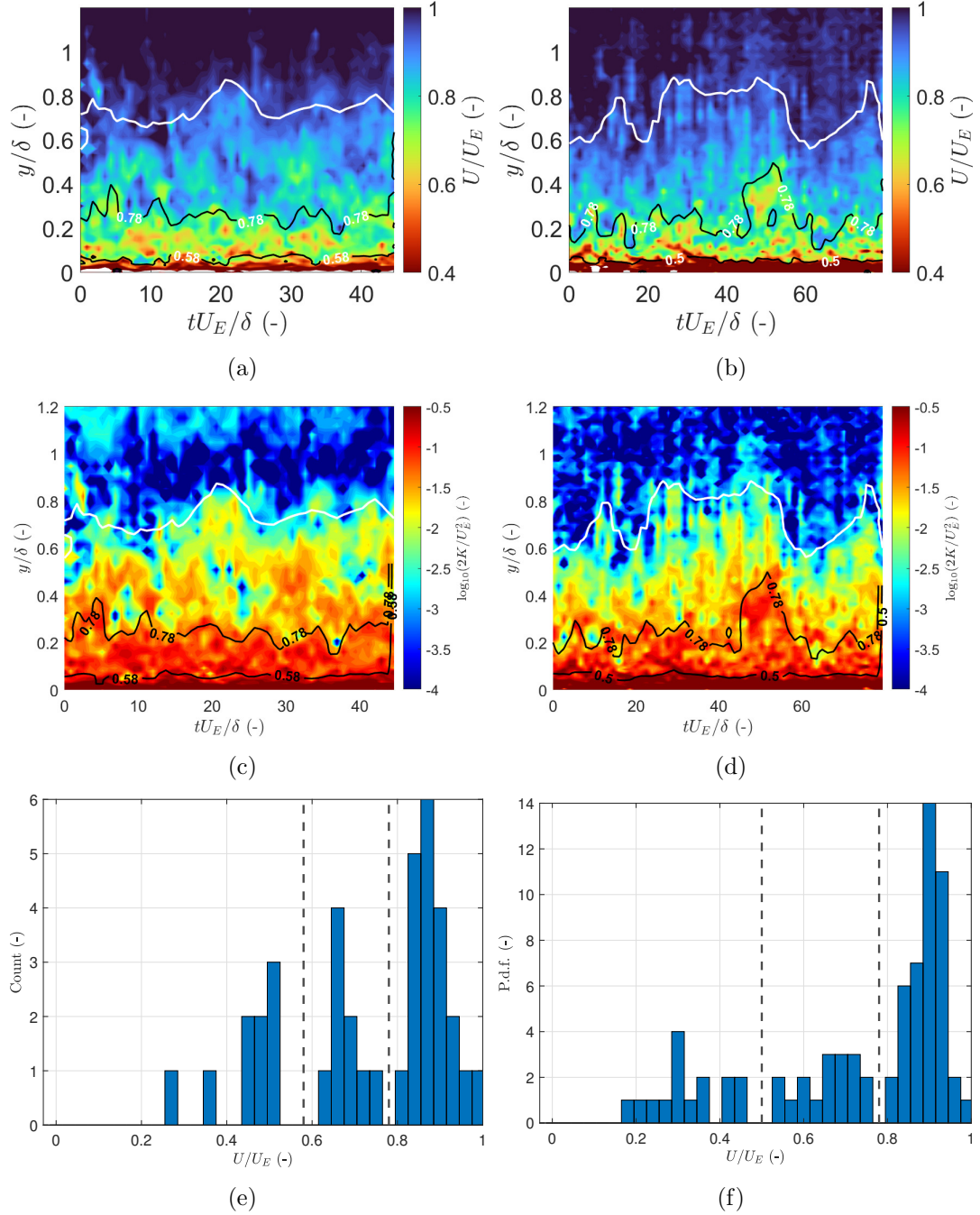


Figure 7.8: Comparison of UMZ statistics for a shot representative of the cold-flow and enthalpy-matched run conditions. Panels (a), (c), and (e) correspond to cold-flow shot 385, while panels (b), (d), and (f) correspond to enthalpy-matched shot 500. The TNTI line show in white and UMZ regions in black. Top row: normalized streamwise velocity field. Middle row: turbulent kinetic energy distribution. Bottom row: histogram of U/U_E , with UMZ regions shown with black dashed lines.

The corresponding turbulent kinetic energy fields, shown in (c) and (d), highlight regions of elevated turbulence intensity concentrated near the wall, with intermittent penetration into the outer flow. The turbulent-nonturbulent interface is denoted in white and is the boundary separating the turbulent boundary layer from the irrotational freestream. The boundary is defined, in a similar manner to de Silva et al. [139], as $K = c_0((1/2)U_E^2)$ with $c_0 = 5 \times 10^{-3}$ and $c_0 = 3 \times 10^{-3}$ for the cold-flow and enthalpy-matched conditions respectively. Various coefficients were tried and the chosen c_0 best visualizes the boundary. The difference in c_0 between the conditions is due to the SNR towards the edge of the boundary layer. The higher SNR enthalpy-matched case has a lower c_0 and reduces smoothing in that region.

The UMZ boundaries exhibit significant spatial variability, with pronounced undulations and step-like features. These interface deformations are closely correlated with regions of elevated turbulence intensity, indicating a strong coupling between UMZ organization and the underlying turbulent motions. In particular, inclined regions of high turbulence intensity extending away from the wall are observed to coincide with upward displacements of the UMZ interfaces. Such features are consistent with ejection events, in which low-momentum fluid is transported away from the wall [9, 138]. For example, in (d), the UMZ interface is raised at $tU_E/\delta = 40 - 50$, which is indicative of a strong coherent structure that disrupts the local momentum distribution and reorganizes the UMZ structure. Similar correlations between UMZ interface topology and turbulent activity have been reported in previous experimental studies, in incompressible and low-speed flow ($U_\infty \approx 10 - 20$ m/s) [9, 139], where UMZ boundaries were shown to align with shear layers separating regions of distinct momentum. The observed inclined structures and associated interface displacements are consistent with the dynamics of coherent vortical motions, such as hairpin vortex packets or other wall-attached eddies, which are known to induce ejection and sweep

events. However, it should be noted that the present analysis is based on a reconstructed two-dimensional field and does not directly resolve the three-dimensional vortical topology. As such, the results should be interpreted as evidence of coherent structure activity influencing UMZ organization, rather than a direct identification of specific vortex geometries.

The histograms shown in (e) and (f) provide a statistical representation of the UMZ structure. Using a single realization, non-dimensionalized streamwise velocities are counted to visualize the UMZs throughout the turbulent region. Distinct peaks in the distributions correspond to preferred velocity levels, which are interpreted as the modal values associated with individual UMZs. The areas of no counts between the modes account for the interfaces and are overlaid as black contour lines shown in Figure 7.8. For the present conditions, both the cold-flow and enthalpy-matched

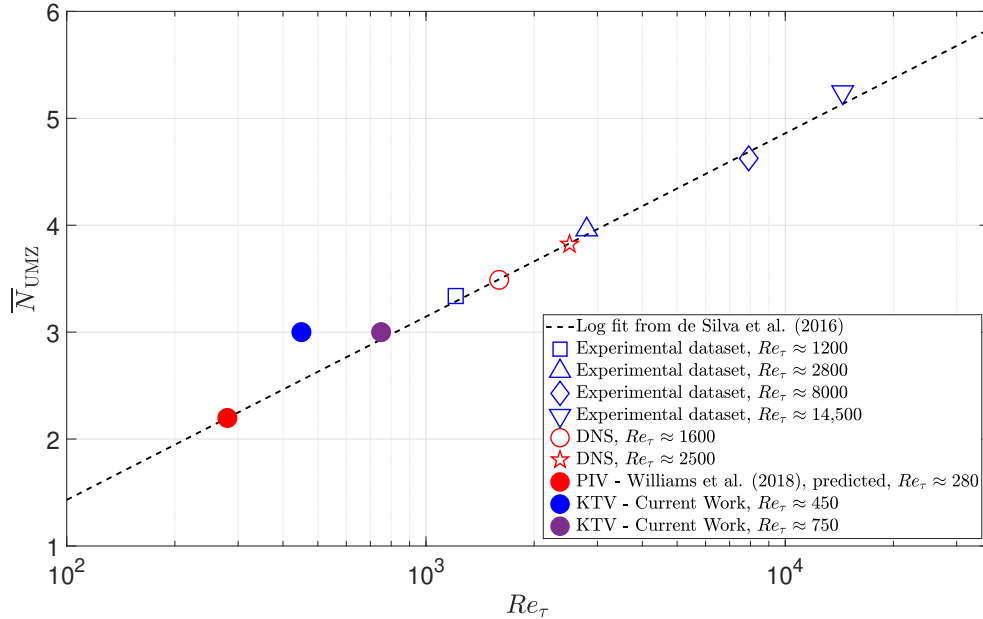


Figure 7.9: Comparison of experimental and DNS data with logarithmic fit for number of UMZ regions ($\bar{N}_{UMZ}$) for varying Re_τ from de Silva et al. [9]. Additionally plotted are the predicted number of UMZs for both the cold-flow and enthalpy-matched run conditions and PIV experiment from Williams et al. [5].

conditions experience 3 UMZ peaks. de Silva et al. [9] fitted experimental and DNS data to a plot of number of UMZs as a function of Re_τ . Fig. 7.9 plots the data from de Silva et al. [9] in addition to predicted values for the KTV experiments and PIV experiments from Williams et al. [5]. For both the cold-flow ($Re_\tau = 450$) and enthalpy-matched ($Re_\tau = 750$), the fit expects approximately 3 UMZs which is consistent with the experiment.

7.3 Turbulence Length and Time Scales

The spatial variation of turbulence length and time scales within the boundary layer provides insight into the distribution of dissipative and inertial processes. The Kolmogorov scales define the smallest dissipative scales in the flow [140]. At these scales, viscous dissipation dominates and turbulent kinetic energy is dissipated into heat [141]. This scale determines the smallest eddies that should be resolved in high-fidelity simulations such as DNS.

Pope [140] defines the Kolmogorov length and time scales as

$$\eta = \left(\frac{\nu^3}{\epsilon} \right)^{1/4} \quad (7.3.1)$$

and

$$\tau_\eta = \left(\frac{\nu}{\epsilon} \right)^{1/2} \quad (7.3.2)$$

where the dissipation rate is defined as

$$\epsilon = \frac{\nu}{\tau_\eta^2}. \quad (7.3.3)$$

Dussauge et al. [32], following Townsend [142] and Perry et al. [143], assume that in the turbulent wall region dissipation is equal to production, so

$$\epsilon = -\overline{u'v'} \frac{\partial U}{\partial y}. \quad (7.3.4)$$

Continuing, Dussauge et al. [32] assumes that the velocity profile follows the log-law and $-\overline{u'v'} = u_\tau^2$. These assumptions can be used with Equations 7.3.1 and 7.3.2 to obtain the Kolmogorov length and time scales as

$$\eta = \left(\frac{\nu^3 \kappa y}{u_\tau^3} \right)^{1/4} \quad (7.3.5)$$

and

$$\tau_\eta = \left(\frac{\nu \kappa y}{u_\tau^3} \right)^{1/2}, \quad (7.3.6)$$

where κ is the von Kármán constant and ϵ is

$$\epsilon = \frac{u_\tau^3}{\kappa y}. \quad (7.3.7)$$

The Taylor microscale, λ , characterizes the intermediate range of turbulent motion, bridging the gap between the large energy-containing eddies and the dissipative scales [141]. Tennekes and Lumley [141] define the Taylor microscale as

$$\lambda = \sqrt{15 \nu u_{\text{RMS}}'^2 / \epsilon}, \quad (7.3.8)$$

and can be interpreted as a measure of the scale at which viscous effects begin to influence the dynamics, but do not yet dominate. Using the relation between η and ϵ in Equation 7.3.4, the Taylor microscale can be rewritten as

$$\lambda = \sqrt{15} \frac{u_{\text{RMS}}' \eta^2}{\nu}. \quad (7.3.9)$$

Tennekes and Lumley define the time scale for the Taylor microscale as

$$\tau_\lambda = \sqrt{15} \tau_\eta. \quad (7.3.10)$$

The Taylor microscale is larger than the Kolmogorov scale. Figure 7.10 shows the length scales throughout the boundary layer for both run conditions. The turbulence length scales are observed to be significantly larger in the enthalpy-matched condition compared to the cold-flow case. This behavior is primarily attributed to the strong dependence of viscosity on temperature in high-speed flows. In the enthalpy-matched condition, the higher gas temperature results in an increased dynamic viscosity, and consequently a larger kinematic viscosity. By using local values for density and viscosity throughout the boundary layer, the estimated microscales are influenced by the strong variations of kinematic viscosity. Although the dissipation rate from

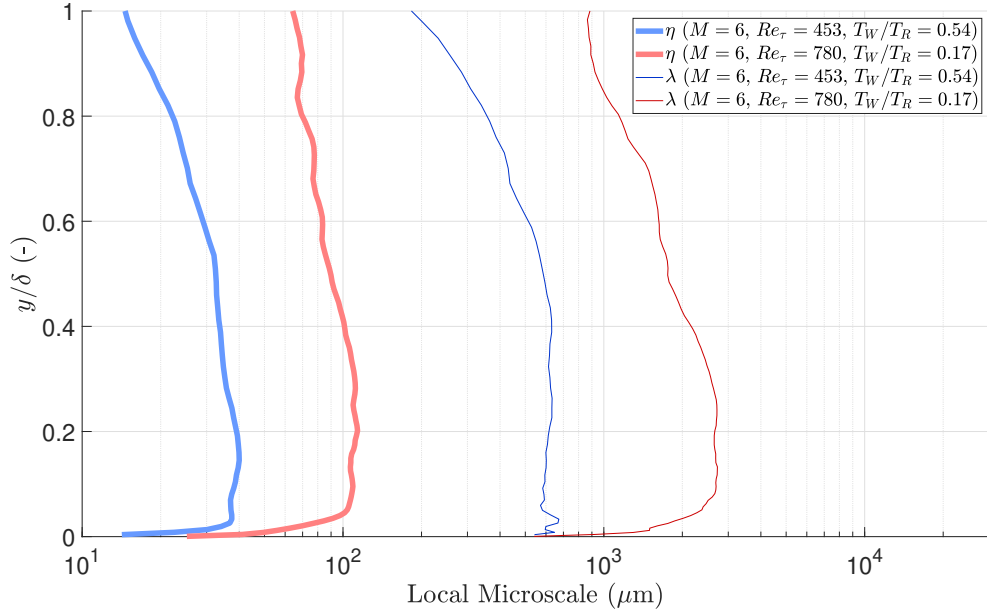


Figure 7.10: Dimensional Kolmogorov length scale (η) and Taylor microscale (λ_g) throughout the boundary layer for both cold-flow ($M = 6$, $Re_\tau = 453$, $T_W/T_R = 0.54$) and enthalpy-matched ($M = 6$, $Re_\tau = 780$, $T_W/T_R = 0.17$) run conditions.

Equation 7.3.7 would cause η to increase with wall-normal distance for constant ν , the local kinematic viscosity decreases from the hot, low-density near-wall region toward the boundary-layer edge. The Taylor microscale shows a similar behavior, with additional sensitivity to the measured fluctuation intensity.

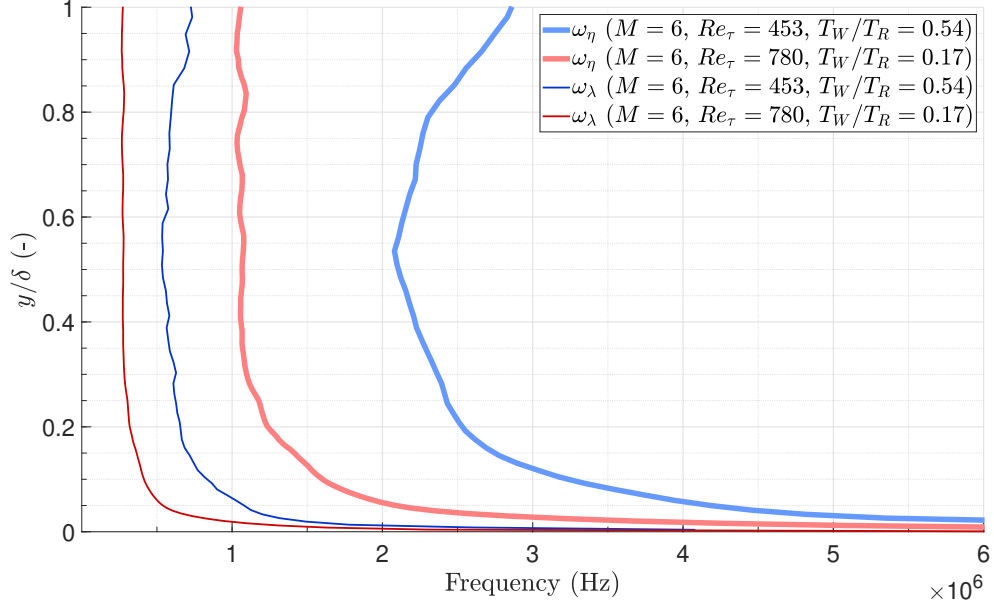


Figure 7.11: Comparison of Kolmogorov and Taylor characteristic frequencies throughout the boundary layer cold-flow ($M = 6$, $Re_\tau = 453$, $T_W/T_R = 0.54$) and enthalpy-matched ($M = 6$, $Re_\tau = 780$, $T_W/T_R = 0.17$) run conditions.

The characteristic frequencies shown in Fig. 7.11 are calculated as $\omega_\eta = 2\pi/\tau_\eta$ and $\omega_\lambda = 2\pi/\tau_\lambda$ for the Kolmogorov and Taylor scales respectively. The enthalpy-matched case exhibits larger Kolmogorov and Taylor microscales, therefore its corresponding characteristic frequencies are lower throughout most of the boundary layer. Conversely, the smaller scales observed in the cold-flow condition produce higher characteristic frequencies. This indicates that turbulent motions in the cold-flow case evolve on shorter time scales, while the enthalpy-matched cold-wall condition is associated with comparatively slower dissipative and intermediate-scale dynamics.

7.4 Conclusion and Future Work

The combined KTV data show that both mean velocity and Morkovin-scaled fluctuation profiles are consistent with DNS and established incompressible reference data under the tested Mach 6 conditions. The observation of UMZs in the present hypersonic measurements, consistent with the incompressible low-speed experiments of de Silva et al. [9], further supports Morkovin's hypothesis [3] by indicating that the organized shear-layer structure of compressible turbulent boundary layers retains key similarities to incompressible wall-bounded turbulence. Future work includes scaling ultra-high-speed FLDI spectra with estimates in Figure 7.11.

Chapter 8

Conclusions and Future Work

8.1 Conclusions

A detailed experimental investigation of turbulence in hypersonic boundary layers was performed using non-intrusive optical diagnostics in the Stevens Shock Tunnel. The primary objective of this work was to improve the understanding of compressible turbulence behavior and to provide experimental data for evaluating scaling approaches, particularly Morkovin’s hypothesis, in Mach 6 flows over a hollow-cylinder-flare configuration.

Streamwise velocity and fluctuation measurements were obtained using both acetone tagging velocimetry (ATV) and krypton tagging velocimetry (KTV). These measurements showed favorable agreement with existing experimental data and direct numerical simulation (DNS) results when scaled using the Morkovin transformations, supporting the applicability of incompressible-based scaling to compressible turbulent boundary layers under the conditions tested.

A novel measurement technique was developed to obtain wall-normal velocity fluctuations using a single inclined laser line. This method enabled non-intrusive access to wall-normal turbulence quantities that are difficult to measure using conventional diagnostics. The resulting wall-normal fluctuation profiles exhibited trends consistent with DNS predictions and addressed known discrepancies observed in particle-based measurement techniques, which are subject to attenuation due to particle lag and spatial filtering effects.

The results of this work provide new experimental evidence supporting Morkovin’s

hypothesis in hypersonic turbulent boundary layers, particularly for both streamwise and wall-normal fluctuation components. The ability to accurately measure these quantities improves the understanding of turbulence structure in high-speed flows and provides valuable data for the development and validation of turbulence models used in computational fluid dynamics. These findings contribute to reducing uncertainty in the prediction of aerodynamic drag and heat transfer in hypersonic vehicle design.

From an engineering perspective, the measurement techniques developed in this work are not only tools for studying turbulence, but also tools for reducing uncertainty in hypersonic vehicle design. In practical applications, quantities such as skin friction, heat transfer, and turbulent mixing directly affect thermal protection system sizing, structural loading, propulsion-system performance, and vehicle range. Because these quantities are typically predicted using CFD and turbulence models, the accuracy of the design process depends on experimental data that can validate the modeled Reynolds stresses and boundary-layer structure under realistic high-speed conditions. Non-intrusive diagnostics such as ATV and KTV provide this information without placing probes into the flow, which is especially important in hypersonic boundary layers where intrusive sensors can disturb the very turbulence being measured. The ability to resolve streamwise and wall-normal velocity fluctuations, identify turbulent structure, and estimate relevant length and time scales therefore provides a physical basis for improving turbulence closures, assessing whether ground-test measurements capture the dynamically important flow scales, and ultimately making more reliable predictions of aerodynamic drag and heating on hypersonic flight vehicles.

8.2 Future Work

Future work will focus on expanding the experimental capabilities and measurement techniques developed in this study. A primary effort will involve the implementation of a Mach 3.5 nozzle in the Stevens Shock Tunnel. This addition will expand the test envelope and allow for broader parameterization of run conditions. Additionally, a heated blanket system was recently commissioned and instrumented for the Stevens Shock Tunnel. This system will enable operation at higher reservoir temperatures, allowing access to higher-enthalpy flow conditions. This capability will extend the parameter space of the facility and enable investigation of turbulence behavior under conditions more representative of true hypersonic flight, including stronger thermodynamic coupling and increased compressibility effects.

Efforts have been started to develop and implement surface instrumentation to quantify heat transfer and wall shear stress on the HCF test article. While non-intrusive optical diagnostics such as KTV provide detailed information on velocity fluctuations and turbulence structure, surface measurements are required to directly assess aerodynamic heating and skin friction, which are critical to high-speed vehicle design and validation.

To address this need, the HCF test article was instrumented with fast-response unsteady temperature gauges and wall shear stress sensors capable of resolving the short test times inherent to reflected-shock tunnel facilities. Seven unsteady surface temperature gauges and two wall shear sensors, all supplied by Ahmic Aerospace, were instrumented to the HCF. The locations of the upgraded instrumentation can be seen in Fig. 8.1. The locations on the cylinder and the Reynolds number (Re_x) for each of the test conditions are presented in Table 8.1. Seven unsteady temperature gauges were installed on the cylinder; however, only four are highlighted in this work

because only four signal-conditioning amplifiers were available. Each temperature gauge provides a fast time response greater than 100 kHz and a high output sensitivity of 0.5 mV/°F, with a thermal response of 0.0045 °F/μs. The sensing element is housed in a cartridge with an outer diameter of 1.59 mm (0.0625 in), which is encased in a 316 stainless steel sleeve with a 6.36 mm (0.25 in) outer diameter to match the material and wall thickness of the HCF model. The exposed surfaces of the gauges are contoured to align with the outer diameter of the hollow cylinder.

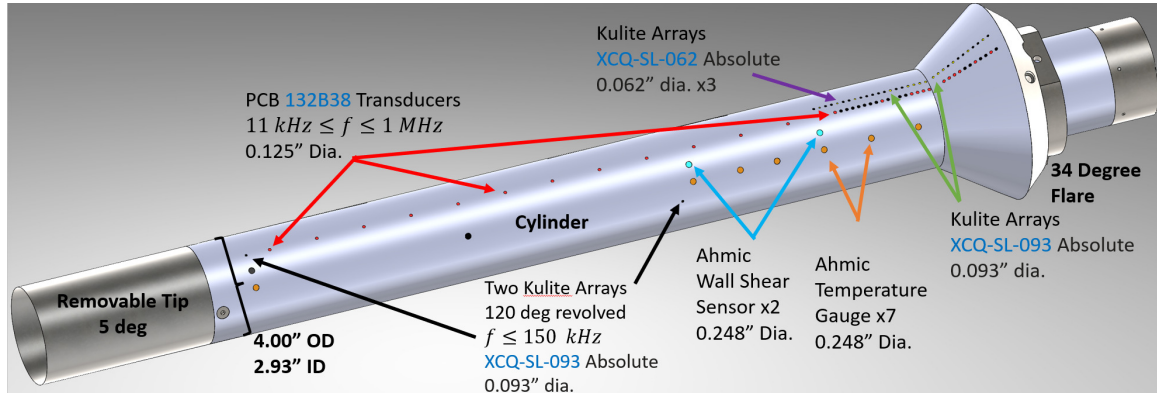


Figure 8.1: Schematic of re-instrumented hollow-cylinder flare test article. Unsteady temperature gauges are denoted in orange and wall shear stress gauges are denoted in cyan.

Table 8.1: Axial locations and distance from the leading edge for the unsteady temperature gauges. Local Reynolds number for both run conditions, defined as $\rho_e U_e x / \mu_e$.

Gauge	Distance from leading edge	$Re_x \times 10^6$ cold-flow	$Re_x \times 10^6$ enthalpy match
1	9.0	2.65	1.24
3 (FLDI Location)	27.6	8.10	3.79
5	31.1	9.15	4.28
6 (KTV Location)	33.1	9.73	4.56

Each cartridge provides two outputs: an instrumentation cable for the unsteady temperature signal and a Type-T thermocouple embedded within the housing. The instrumentation lead is routed through a facility feed-through plate and terminated at an external completion box. The completion box output is then connected to a

Stanford Research SR560 constant-voltage preamplifier, configured with a 100 kHz low-pass filter and a gain of 1. The amplified signal is subsequently digitized using a PXIe-6378 data acquisition card operating at 3 MHz.

The wall shear sensors provide a measurement bandwidth from 250 Hz upwards, with the ability to resolve shear stresses as low as 25 Pa. Each sensor is housed in a 6.36 mm (0.25 in) diameter package, making it fully interchangeable with the temperature gauges on the HCF cylinder. Like the temperature sensors, the exposed surfaces of the shear sensors are contoured to match the outer diameter of the hollow cylinder. Proper orientation of each sensor is critical, as the device produces a voltage drop when the applied shear stress is directed upstream along the surface of the cylinder. The sensor output is routed to a built-in amplifier that completes the Wheatstone bridge and applies a 5 kHz low-pass filter and a gain of 1000 prior to external data acquisition. The amplified signal is digitized using a Cleverscope C328A digital oscilloscope operating at 100 MHz.

In addition to facility modifications, future investigations will explore alternative non-intrusive diagnostic techniques that reduce experimental complexity. One promising approach is using the fifth harmonic of the Nd:YAG to excite krypton, which enables high-speed, spatially resolved flow measurements without the need for a wavelength-tuned optical parametric oscillator (OPO).

Further work may also include extending the current measurement methodology to additional flow configurations, such as shock-wave/boundary-layer interactions and varying wall thermal conditions, to assess the generality of the observed turbulence behavior. Continued development of both experimental diagnostics and facility capabilities will contribute to improving the measurement and modeling of hypersonic turbulent boundary layers.

Appendix A

Previous Turbulence Experimental Data

Table A.1: Run conditions for various turbulence experiments. P_R and T_R are the reservoir pressure and temperature, respectively. P_∞ , T_∞ , U_∞ , M_∞ , and Re_∞^U are the freestream pressure, temperature, velocity, Mach number, and unit Reynolds number, respectively. $u_\tau = \sqrt{\tau_w/\rho_w}$, δ , $Re_\tau = u_\tau \delta / \nu_w$, $Re_\theta = \rho_\infty u_\infty \theta / \mu_\infty$, T_w/T_r , and St are the friction velocity, boundary-layer thickness, friction Reynolds number, momentum thickness Reynolds number, wall to recovery temperature ratio, and Stokes number respectively.

Method	P_R (MPa)	T_R (K)	P_∞ (kPa)	T_∞ (K)	U_∞ (m/s)	M_∞ (-)	Re_∞^U (1/m)	u_τ (m/s)	δ (mm)	Re_τ (-)	Re_θ (-)	T_w/T_r (-)	St (-)
PIV - Williams et al.	9.60	542	1.37	48.3	1046	7.45	4.42×10^6	59.5	10.2	280	9340	0.79	~ 1
PIV - Brooks et al. Exp 1A	2.95×10^{-2}	306	1.09	122.4	621	2.80	2.30×10^6	59.9	8.2	202	1167	1	1.0
PIV - Brooks et al. Exp 3C	1.01×10^{-1}	307	4.35	123.3	601	2.70	8.90×10^6	47.5	11.4	916	6945	1	0.3
PIV - Tichenor et al.	2.35	380	4.98	65.5	795	4.90	4.61×10^7	39.0	7.8	755	43000	1	~ 1
LDA - DeGraaff and Eaton	-	-	-	-	18.95	0.056	1.30×10^6	0.727	35.9	1700	5200	1.00	(-)
DNS - Zhang et al.	0.93	433	0.68	55.0	870	5.86	9.95×10^6	45.1	23.8	453	9455	0.76	(-)
DNS - Zhang et al.	3.40	693	0.39	51.8	1155	7.87	8.16×10^6	54.3	35.2	480	9714	0.48	(-)

Appendix B

Streamwise and Wall-Normal KTV Experiment Data

The streamwise experiments ($\bar{u}, u'_{\text{RMS}}$) were averaged over shots 385–390, and the wall-normal experiments (v'_{RMS}) were averaged over shots 415, 418, 420–421 and 423. Conditions are given in Tables B.1 and B.2. Run conditions are sufficiently consistent such that the results may be averaged between runs, as the Re_∞^u standard deviation is 0.87×10^6 1/m or 6%, which is slightly higher than the estimated tunnel noise.

Table B.1: Run conditions for all N₂ driver experiments.

Shot	Experiment	P_R (MPa)	h_R (MJ/kg)	T_R (K)	P_∞ (kPa)	T_∞ (K)	M_∞ (-)	U_∞ (m/s)	Re_∞^u (10^6 1/m)	T_w/T_r (-)
385	u,u'	2.18	0.51	532.6	1.11	65.3	6.31	942	13.9	0.55
386	u,u'	2.19	0.51	533.4	1.08	64.9	6.35	943	13.6	0.55
387	u,u'	2.19	0.51	534.0	1.07	64.9	6.35	944	13.6	0.55
388	u,u'	2.56	0.54	564.1	1.29	68.0	6.34	970	15.2	0.52
389	u,u'	2.94	0.57	597.5	1.49	71.3	6.33	998	16.4	0.49
390	u,u'	2.17	0.51	532.3	1.06	64.6	6.36	943	13.4	0.55
415	v'	2.16	0.51	531.8	1.05	64.6	6.36	943	13.4	0.55
418	v'	2.20	0.51	534.5	1.07	64.8	6.36	945	13.5	0.55
420	v'	2.21	0.51	535.9	1.09	65.1	6.35	946	13.7	0.55
421	v'	2.20	0.51	534.7	1.07	64.8	6.36	945	13.5	0.55
423	v'	2.18	0.51	531.5	1.04	64.5	6.36	942	13.4	0.55
Average	u,u',v'	2.28	0.52	542.0	1.13	65.7	6.35	951	14.0	0.54

Table B.2: Run conditions for all He driver experiments.

Shot	Experiment	P_R (MPa)	h_R (MJ/kg)	T_R (K)	P_∞ (kPa)	T_∞ (K)	M_∞ (-)	U_∞ (m/s)	Re_∞^u (10^6 1/m)	T_w/T_r (-)
500	u,u'	2.53	2.01	1780	1.47	236.7	5.96	1786	2.78	0.16
502	u,u'	2.46	2.00	1770	1.46	236.7	5.94	1780	2.77	0.17
503	u,u'	2.52	2.00	1774	1.48	236.4	5.95	1783	2.81	0.17
481	v'	2.40	1.82	1627	1.37	211.5	6.00	1702	3.02	0.18
482	v'	2.43	1.83	1632	1.39	212.4	6.00	1704	3.04	0.18
485	v'	2.26	1.77	1586	1.28	204.7	6.02	1678	2.96	0.18
486	v'	2.39	1.86	1658	1.40	217.8	5.97	1718	2.95	0.18
488	v'	2.41	1.88	1676	1.37	219.0	5.99	1729	2.88	0.17
Average	u,u',v'	2.43	1.90	1698	1.41	223.8	5.97	1740	2.89	0.17

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