

Effect of density ratio on velocity dynamics in the blast-driven instabilitySamuel J. Petter[✉], Benjamin C. Musci[✉], Gokul Pathikonda^{✉,*},Prasoon Suchandra[✉], and Devesh Ranjan^{†,‡}*George W. Woodruff School of Mechanical Engineering, Georgia Institute of Technology,
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Results are presented for the explosive blast-driven instability of a stratified interface in a diverging cylindrical geometry using high-speed optical diagnostics. Specifically, the perturbation growth was studied using the qualitative Mie scattering technique, and the velocity, mixing transition, and vorticity characteristics are studied to investigate the behavior of the interface using planar particle image velocimetry (PIV). The role of density contrast is studied using CO₂-air (low Atwood number, A) and SF₆-air (high- A) interfaces. The perturbation growth $h \sim t^\theta$ was studied, where the decay parameter θ was found for $A \sim 0.22$ and $A \sim 0.68$ to be 0.34 and 0.50, respectively. Using velocity information from PIV, it was found that the initial vorticity deposition from the pressure impulse can accurately be estimated using known models previously developed for Richtmyer-Meshkov instability. There is a significant addition to the interface circulation by the long pressure decay phase of the blast wave, which gives rise to a variable acceleration Rayleigh-Taylor instability contribution. The high- A case showed a prolonged increase in circulation, plausibly due to higher inertia. This trend was also supported by variations in peak turbulence intensity, where a sharp increase at early times quickly decays with increased mixing between the two fluids. Finally, the interface Reynolds number of the higher A was found to exceed all criteria for the mixing transition (also indicated by the qualitative observations of the mixing in spike-structures in the Mie scattering images).

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Blast-driven instability (BDI) has generated particular interest due to the fluid instabilities present, which often lead to a turbulent state. The instabilities present have an important effect in violent supernova (SN) explosions, as well as in inertial confinement fusion (ICF). These instabilities are difficult to account for in variable density, compressible flow models (like those used for SN and ICF predictions) [1–7], particularly because numerical techniques used to capture shock discontinuities impair resolution of turbulence [8,9]. In SN they are also large contributors to events where material is ejected from the bulk of the mixing region, and in ICF they inhibit successful fusion [10]. The two primary fluid instabilities present are the buoyancy-driven Rayleigh-Taylor

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instability (RTI) and the shock-driven Richtmyer-Meshkov instability (RMI). RTI is a variable density fluid instability that occurs when the density and pressure gradients of a fluid system point in opposite directions (i.e., there is a misalignment of density and pressure) [11–14]. This can be conveniently imagined as a heavy fluid resting over a lighter fluid under the influence of gravity. Here, even the smallest imperfections along the fluid interface (i.e., perturbations) will grow, stretching the interface between the two fluids. At the impulsive limit a special form of this instability occurs, namely the Richtmyer-Meshkov instability [15,16]. RMI occurs when the interface between two fluids of different density is impulsively accelerated. This is conventionally studied in shock tubes, where the impulsive force is provided by a shock wave. A fluid interface becomes unstable when impulsively accelerated (shocked) regardless of the direction of acceleration [17]. A detailed review of these hydrodynamic instabilities and previous experimental and computational works are presented in [18–20].

RTI and RMI have been extensively studied by numerous authors. The theoretical aspects of these instabilities in their combined form have been successfully captured using the group theory approach, which enabled the studies of both RTI and RMI dynamics in a broad range of conditions, including accelerations varying in time and space [1,21–24]. Traditional experiments typically focused on either RTI or RMI individually, though recent efforts have begun to study these instabilities in their combined form, particularly high energy density physics (HEDP) experiments like those performed at the National Ignition Facility (NIF). These experiments involve nonfluid materials that are ablated with high-power lasers to create plasmas, which then mix like fluids [25,26]. Samples for these experiments are designed with initial perturbations machined into the plastic to amplify and explore RMI and RTI effects. The early growth of these experiments is captured and compared with the commonly used buoyancy drag model [27]. Two features to note about these studies are that (i) very large density ratios can be achieved using solid targets, and (ii) heat transfer effects are present, making them closely representative of ICF and SN conditions. Nevertheless, the extreme nature of these experiments is challenging for diagnostics, where x-ray diagnostics have only faintly resolved snapshots in time of the development of this combined instability. These efforts have revealed the need for later growth times, finer temporal and spatial resolution, and a wider mixing plane for more comprehensive statistical studies. A limited number of studies have explored time-varying acceleration as the driving mechanism for RTI/RMI, contributing valuable insights into this complex phenomenon. Notable among these are the works that investigated the interplay of acceleration profiles and instability growth [28,29]. While these contributions have advanced our understanding of the field, further research to build upon these foundational studies and fully characterize the effects of dynamic acceleration on RTI/RMI behavior is still necessary. The particular case considered here is a combined form of RTI/RMI, i.e., the BDI.

In the classical sense, the term “blast” is used when describing strong radially expanding shocks whose pressure and energy release can be characterized by a similarity of the first and second type [21,30–33]. Here, the term “blast” describes a radially expanding wave that is generated by a high energy density “point source,” and is characterized by overpressure and the shape of the pressure waveform [32,34–38]; see Fig. 2(a). Note that here, the words “blast wave” are used for convenience and are traditional for this specific facility, though it is emphasized that “explosively driven” is also suitable terminology. Blast waves are directly coupled with the expansion fan, and when this interacts with a variable density fluid interface, there is an initial RMI-phase, immediately followed by an RTI-phase, where the coupled rarefaction causes a variable deceleration driving mechanism that propels the instability to develop. Thus, when a blast wave interacts with a fluid interface, there is a combined RMI/RTI effect, due to the nature of the blast. Other authors have also used the term “blast” to describe weaker pressure waves due to the shape of the shock and the nature of the pressure wave [9,39,40].

Although the BDI has no shortage of attention, capturing experimental data is challenged by the difficulties present in both laser and blast-driven environments. Additionally, little experimental work has been done to resolve the velocity fields of the BDI. The current study experimentally investigates this combined RMI/RTI instability, and specifically the following unexplored areas:

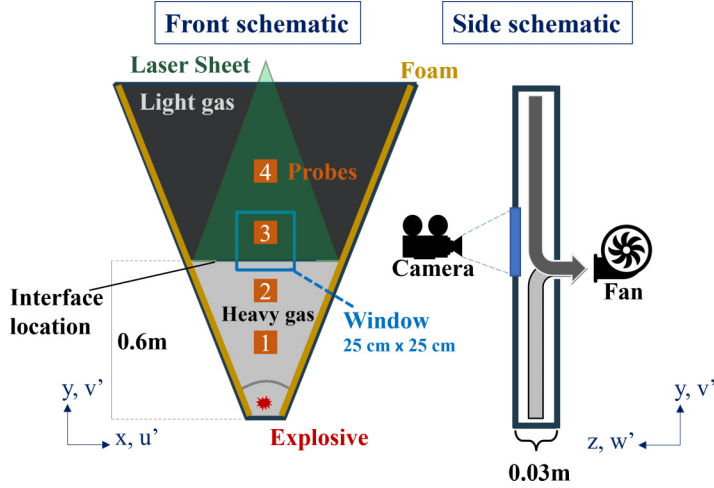


FIG. 1. A front schematic (left) and side schematic (right) of the Georgia Tech blast wave facility.

(i) isolating the physics of the variable acceleration generated from the interface and blast wave interaction, and (ii) a diverging geometry. Creating a radially diverging blast wave produces a pressure front that decays immediately and rapidly after the impulsive rise due to the expanding region behind the outwardly traveling wave (since the expansion fan is directly coupled with the shock front). When an interface is subject to these conditions, an initial “RMI-phase” sets up conditions that are further enhanced by a later “RTI-phase” which follows; this BDI naturally shares features of both RMI and RTI. Similarities and differences between RTI and RMI dynamics were previously identified numerically within the group theory approach for a broad range of conditions [1,21,41]. It is important to study fluid instabilities both individually and combined because they may not combine linearly (e.g., coupling shear and buoyancy, it is found that when shear increases, it leads to a nonmonotonic rise in growth rate [42]).

It was shown by Musci *et al.* [40] that the Georgia Tech Blast Wave Facility can achieve Euler similarity to previous HEDP experiments by employing a scaling proposed by Ryutov *et al.* [43]. This relates hydrodynamic mixing behavior across vastly different spatial scales, neglecting heat transfer and other physics. The advantage here is that the physical size of the experimental facility permits the use of high-resolution optical diagnostics (while maintaining Euler similarity). This addresses several of the following open scientific questions from [44]: (i) How does a variable acceleration change the RTI development? (ii) How does a polar converging/diverging RTI behave

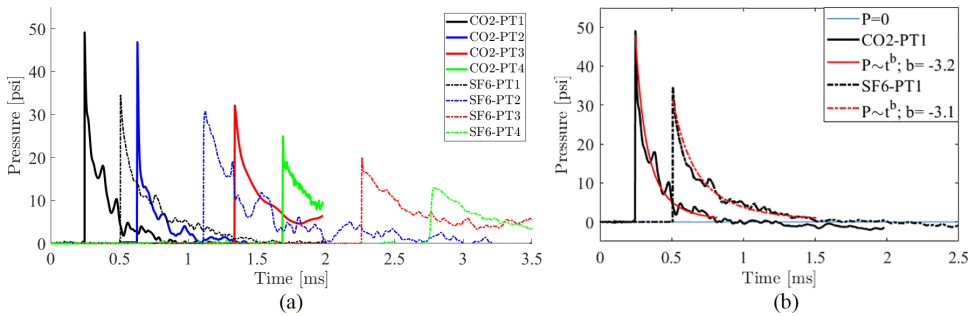


FIG. 2. (a) Pressure traces indicating the blast wave position along the facility. (b) Power-law best-fit profiles.

in contrast with the planar case? And eventually, (iii) how do drivers with different pressure profiles change the nature of RMI?

Finally, the combination of these instabilities is of interest to those studying SN and SN-type mixing, as they are thought to contribute to mass ejecta in SN explosions, and their presence in SN is not fully explained by simulations [7]. It is also of interest in ICF, as these instabilities are primary contributors to the inefficacy of the fusion reaction, thus understanding this is of key importance in HEDP [7,26,45–47].

This work presents velocity fields in the BDI using high-speed particle image velocimetry (PIV) for the first time in efforts to better understand the driving mechanisms preventing successful fusion and SN ejecta. It is a subset of experiments from a much larger experimental data set [48,49] that is part of an ongoing effort to investigate a large parameter space studying the effects of varying the Mach number, M , and the density ratio presented as the dimensionless Atwood number, using Mie scattering. The Atwood number, defined as $A = (\rho_h - \rho_l)/(\rho_h + \rho_l)$ (where ρ_h is the density of the heavy gas and ρ_l is the density of the light gas), is a dimensionless number that normalizes the density ratio from 0 to 1. It was discovered that the Atwood number has a large influence on the morphological development of perturbations. This work examines two density ratios ($A \sim 0.22$ and $A \sim 0.68$) to look closer at the dynamics of this through particle image velocimetry. Atwood numbers were selected where refractive index mismatch was not a large factor (i.e., excluding extrema in A), as well as information from preliminary results from the larger set of Mie scattering experiments, which provided information about interface location and allowed for the zoomed focus of the PIV diagnostic on finer structures. First, the growth parameter (or decay exponent), θ , is presented, and then PIV analysis allows the turbulence intensity to be evaluated. The size of the facility and experimental test time make spatial and temporal scales more manageable compared to HEDP experiments. Additionally, the aim of the non-Cartesian geometry of the experimental facility is to make progress in producing experimental data for more complex simulations, such as those in [9].

II. EXPERIMENTAL SETUP

This meter-scale blast wave facility was designed to provide highly resolved temporal and spatial measurements of hydrodynamic phenomena in blast-driven flows, which are otherwise not feasible in HEDP experiments. The facility is a vertical chamber with two large parallel plates spaced closely at 25.4 mm apart. They span a height of 2 m and diverge from a vertex forming an angle of 45° . Various inlets allow for gases to be injected into the facility. Accordingly, a heavy gas is injected into the bottom, and a light gas is injected into the top half of the chamber. This forms a stable interface between the two gases at a small slit in the interface plate. A schematic of the facility is shown in Fig. 1.

Additionally, a linear array of dynamic pressure transducers (PT 1–4) mounted on the sidewall record the time of arrival of the blast wave and the shape of the pressure profile. Figure 2(a) shows typical pressure traces measured by the transducer array that capture the blast wave progression through the test section (filled entirely with one gas, i.e., carbon dioxide, CO_2 , or sulfur hexafluoride, SF_6). This enables the calculation of M through each medium.

The blast wave is generated by detonating a small amount of explosive (RDX + PETN) in the chamber shown at the bottom of the facility. This results in an instantaneous deposition of the energy over an infinitesimally small volume, resembling a Taylor-Sedov (TS) blast wave [Fig. 2(b)], particularly how the expansion directly follows the blast front [33,40]. However, the blast waves produced are not strict TS blast waves, as they are weak blasts ($M < 5$) [50] so they do not follow the self-similar scaling as strong blasts do. (For spherical geometry, the overpressure scales as $t^{-6/5}$, and for cylindrical geometry the overpressure scales as t^{-1} .) Since the overpressure is a sharp transient, it cannot be measured reliably, so although the radial profile has previously been shown to fit well [40], when comparing it is seen to decay faster than the self-similar solution and may be better represented with a power-law fit of -3 (and at later times it is even a linearly decaying acoustic

wave). Thus, since it is a weak blast wave, it does not exactly follow the TS self-similar solution, but the immediate drop in the pressure profile resembles a TS blast wave profile. The pressure drop from the expansion is captured well by the pressure transducers, but the negative phase is not (since the negative pressure is not strong enough for the piezoelectric crystals to accurately respond). M can be controlled via the detonator strength, as demonstrated in [40]. The blast wave diverges radially and weakens (M decreases in time). The current study primarily focuses on the RP80 detonator (comprised of 123 mg of RDX and 80 mg of PETN) that yields a Mach number of ~ 1.5 at the interface when traveling through air. The blast wave propagates faster through CO_2 than SF_6 due to the different speeds of sound, resulting in Mach numbers 1.5 and 1.4, respectively (7% smaller through SF_6). A 5 V trigger signal from Teledyne RISI detonation system (that also provides a 4 kV detonating pulse for detonation) is used to synchronize the PIV acquisition system timing.

A. Initial conditions: Perturbation generation and characterization

A large port on the facility is equipped with an interface plate that spans the width of the facility. The interface is a narrow slit that spans less than 1 mm in thickness (cut with wire electric discharge machining). Gentle suction is generated by a fan connected to a small box mounted to the interface plate. This allows for a crisp interface to be sustained between the heavy (*bottom*) and light (*top*) fluids. By changing the fan speed, multimode perturbations are introduced, with amplitudes significantly smaller than the wavelength. (Note here that a slight shear is introduced at the interface here in the out-of-plane component, but its influence is not significant due to the small characteristic timescale of the instability development.) The initial perturbation amplitudes can be controlled by increasing the fan speed. However, since the location of crests and troughs is not precisely controlled, the initial perturbations for each experimental run are quasiperiodic and can be described (via wavelength and amplitude) in a stochastic sense. Statistically stationary Mie scattering snapshots of initial conditions were collected, and profiles of the initial conditions were traced using the max gradient of each column. Then, fast Fourier transforms (FFTs) of this wave in x - y space were used to get wavelengths. From here, the dominant wavelengths were averaged, giving a distribution of perturbation wavelengths ranging from 1 to 5 cm, with an average wavelength of 2 cm. The resulting blast wave from the detonation forms a curved shock that interacts with this interface. The PIV field of view (10 cm $\times$ 10 cm) enables the capture of multiple structures in each experiment.

B. Diagnostics

Two laser-based optical diagnostics are used in the current work: Mie scattering to study the interface dynamics and PIV for velocity and vorticity dynamics. For both methods, a 120 W Quantronix Hawk Duo laser with two cavities is used to illuminate the flow. The laser beam is collimated using optical lenses and directed into the chamber through a diverging lens, creating a laser sheet ≈ 1 mm thick. It is important to note that the measurements are in a two-dimensional center plane within the flow field, thus the presented quantities are all calculated from an estimate that the out-of-plane components are of the same order of magnitude as the spanwise (x -direction) perturbations since they are not in the direction of the blast.

Mie scattering gives a visualization of the flow field and qualitatively represents the local mixture fraction of the seeded fluid. In this diagnostic, fog is premixed with the heavy gas before injecting it into the chamber (the light gas is not seeded). Individual particles are not resolved, but instead a bulk visualization of the entire flow field is feasible. The field of view (FOV) spans a 25 cm $\times$ 25 cm square window. With this, the high-speed cameras operating at 20 kHz yield full temporal flow field development of the BDI. This is useful for estimating the camera location for PIV diagnostics, which inherently have a smaller field of view, since individual particles are resolved. (Note, however, that olive oil droplets were used for PIV seeding instead of the fog used for Mie scattering.)

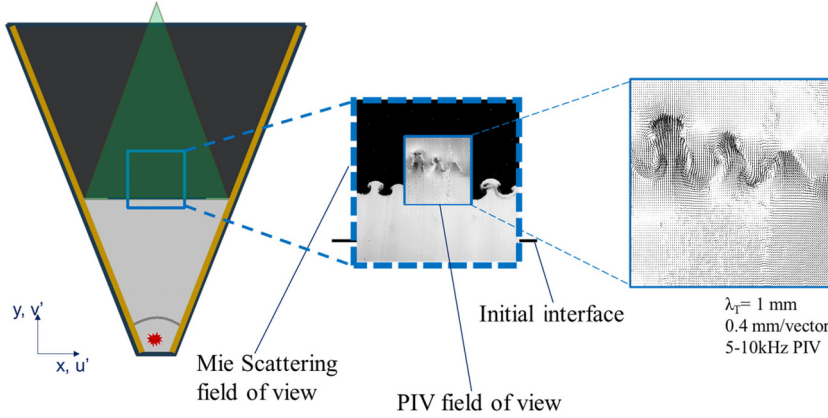


FIG. 3. Schematic showing the Mie scattering and PIV fields of view relative to the facility. The vectors presented here represent the flow field after subtracting the bulk velocity.

Particle image velocimetry (PIV) is used to resolve velocity fields, which offer a more quantitative resolution of structure development. Here, both gases are seeded with olive oil droplets in separate mixing tanks using an olive oil atomizer prior to being introduced into the blast chamber. The droplets are approximately $1\ \mu\text{m}$ in diameter [51], giving a Stokes number of 0.01 for current flow conditions and ensuring particle tracking.

A high-speed laser (10 kHz) is synchronized with a 4-megapixel camera (Phantom v2600 series). Images of particles on subsequent frames are correlated using LaVision Davis to evaluate velocity fields. An example of a resolved vector field is shown in Fig. 9. Relative to Mie scattering, the overall size of the field of view was sacrificed for increased vector resolution. The FOV for PIV presented here is $10\text{ cm} \times 10\text{ cm}$, and the resulting vector spacing is 0.5 mm per vector using an interrogation window of 24×24 square pixel correlation windows, with 50% overlap. (This resolves scales roughly half of the size of the estimated Taylor microscale, $O \sim 1\text{ mm}$, which is an intermediate length scale at which viscosity considerably affects the dynamics of turbulent eddies. Taylor microscale is derived from the estimates of the mixing width, mixing width growth rate, and Reynolds number $\text{Re} = \frac{h\dot{h}}{\nu}$ using dimensional arguments [52,53].) This FOV relative to the Mie scattering is illustrated in Fig. 3. The camera is focused on a region approximately 10 cm above the initial interface. This vertical offset accounts for the initial translation and maximizes the time that the interface is in the PIV field of view (previous Mie scattering experiments provided a good estimate of this interface position).

C. Experimental cases: Low- and high-density contrast

For the present work, two cases are considered to compare the effect of the density ratio, which is characterized by $A = (\rho_h - \rho_l)/(\rho_h + \rho_l)$. The first set ($A \sim 0.22$) is an ensemble of 10 independent runs. In these experiments, the heavy gas is CO_2 and the light gas is air. The second set ($A \sim 0.68$) is an ensemble of 14 independent runs. In these experiments, the heavy gas is SF_6 and the light gas is air. Table I summarizes the Atwood numbers of the initial conditions.

For both cases, the RP80 detonators are used to produce blast waves of consistent strength, as measured by the overpressure and time of wave arrival, $\text{Mach} \sim 1.5$ at the location of the interface. Experiments are approximately 10 ms in duration, ending when the rarefaction from the blast wave exiting the facility reflects onto the mixing region. Two example image sequences from the Mie scattering experiments corresponding to each case are shown in Figs. 4 and 5. Here, the bright region is the heavy fluid, which is seeded with fog that reflects and scatters incident laser light.

TABLE I. Summary of experimental cases.

Case		Light gas	Heavy gas	# of PIV experiments	# of Mie scattering experiments	Mach	Jump velocity	θ
A	0.22	Air	CO2	10	25	1.5	139 m/s	0.34
	0.68	Air	SF6	14	25	1.4	100 m/s	0.50

The A is seen to greatly influence the shape of the structures that form. As the blast wave passes the interface and baroclinic vorticity is deposited, the heavy fluid perturbations grow and penetrate the light fluid, forming structures that are referred to as “spikes,” and the light fluid “falls” into the heavy fluid as “bubbles.” The case where $A \sim 0.22$ demonstrates more “bulbous” bubbles and spikes that are relatively symmetric and are without much small-scale mixing until late times. On the other hand, $A \sim 0.68$ shows sharper structures, particularly spikes, and the interface distinguishing them is more unstable due to strong local shear and deposited vorticity (discussed later in Sec. VII). Furthermore, the instability features completely break up at later times as mixing transition occurs.

III. INTERFACE MOTION

At detonation, the blast wave traverses the chamber, passing through a fluid interface between the two gases. This impulsively accelerates the interface, displacing it upwards. The expansion fan directly coupled to the blast front rapidly decreases this motion, where the interface settles some position above the initial interface location. (A plot of this interface position and its velocity are shown in [40].) The initial perturbations at the interface begin to grow during this motion, causing material to mix. This continues long after the interface has been displaced (interface displacement finishes in the amount of time it takes the blast wave to traverse the entire facility, approximately 2 ms). This material mixing at the interface will be the main focus here, and not necessarily the interface position itself, which is more sensitive to Mach number/detonator strength than to density difference.

IV. PERTURBATION DEVELOPMENT

Initial imperfections are imposed on the fluid interface by perturbing the suction fan speed, which becomes unstable upon interaction with the blast wave. Vortex dipoles form, and they are

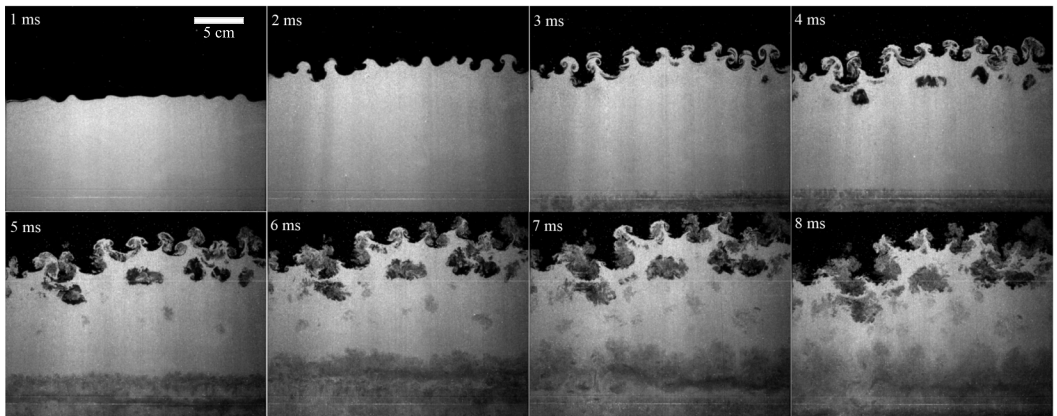


FIG. 4. Perturbation evolution from a Mie scattering image sequence with $A \sim 0.22$. Much more symmetry is observed between bubble/spike structures.

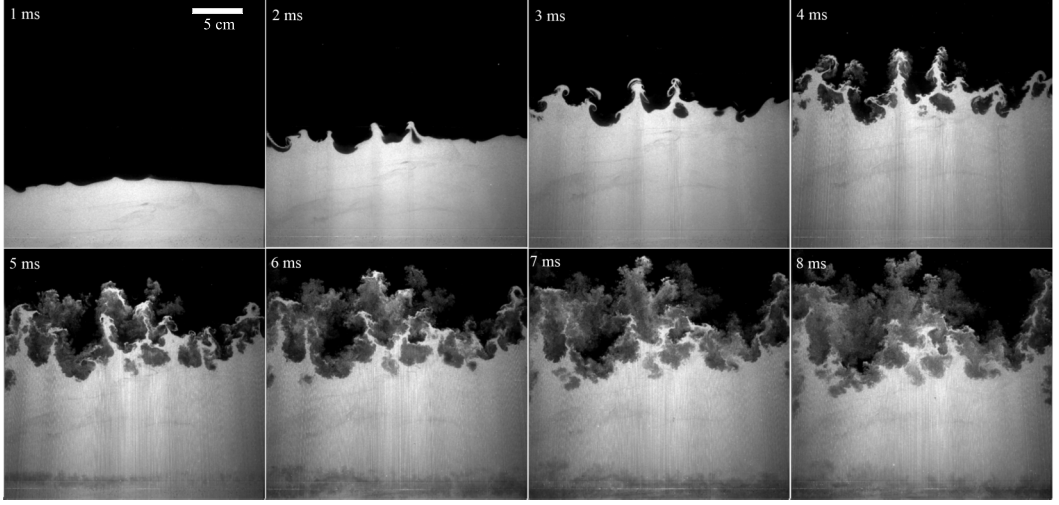


FIG. 5. Perturbation evolution from a Mie scattering image sequence with $A \sim 0.68$. Spikes are much thinner and penetrate the light fluid more.

particularly characteristic spike structures. They form from drag shearing their edges as the heavy fluid penetrates. These structures greatly enhance mixing by stretching the material interface.

The mixing width was calculated from an average of the larger experimental dataset consisting of 25 Mie scattering experiments, based on a 5 – 95 % threshold of pixel intensity. This removes the translation of the interface and just looks at the nondisplaced mixed width. This mixing width [Fig. 6(a)] and its derivative growth rate [Fig. 6(b)] are plotted for each case (shaded regions are one standard deviation). The mixing width can be thought of as a measure of the amplitude of the perturbations, and it indicates the largest length scale present [54]. Larger A leads to deeper spike penetration. The mixing width of the high- A develops to a greater extent. The red box shows the times overlap between the Mie scattering visualizations and the PIV measurements presented here. The vertical lines represent instances in time at which color maps of vorticity and turbulence intensity are presented later in Figs. 9 and 13; black lines are low- A , blue lines are high- A . The

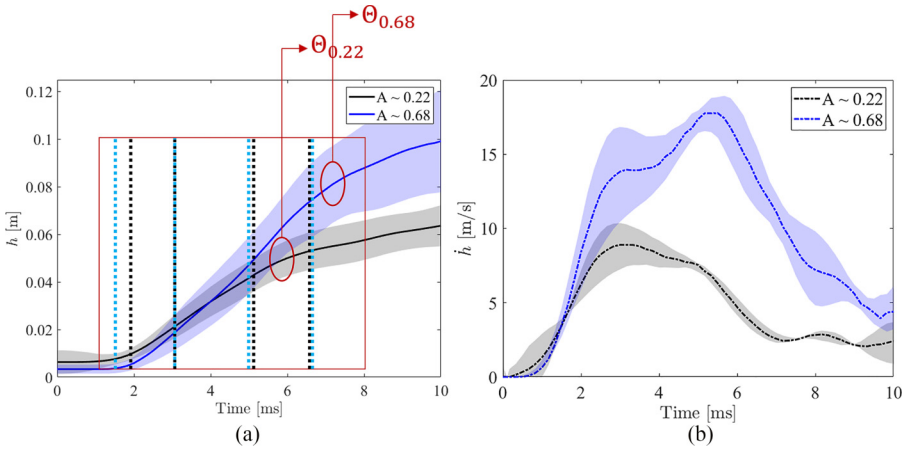


FIG. 6. (a) The mixing width, h , is plotted as a function of time (left). (b) On the right, the mixing width growth rate, $\dot{h}$, is plotted as a function of time for the two Atwood numbers.

inflection point of the mixing width [red circle in Fig. 6(a)] is the point after which the growth parameter is calculated (discussed next in Sec. V).

Because the flow is diverging, geometric effects and volumetric changes should be addressed. Due to the expanding cross-sectional area of the test section, the mass engulfed by the blast increases and the wave speed decreases. Volumetric changes, such as those caused by shock expansion, can alter the density and pressure profiles across the interface, and the geometrical divergence allows perturbations to stretch to a larger surface area. This horizontal/radial stretching decreases spike/bubble penetration. This results in slower instability development compared to an analogous planar RMI/RTI case, often referred to as the Bell-Plesset effect. (Alternatively, in a converging geometry, instability development would speed up.) In the current facility, the volume swept by the movement of the interface is only 3.6% of the total volume swept by the blast when interacting with the interface. (The geometrical divergence causes an increase in volume from 0.0038 to 0.0039 m³.) Moreover, the interface translates to the resting position before the perturbations begin to develop, so 95% of the mixing width's growth occurs at a fixed location.

Notably, at modest density ratios there is much more symmetry between spike and bubble features, but as densities of light and heavy fluids become more disparate, the spike features become much thinner, penetrating the light fluid much deeper, with less rotational energy and more translational energy. This will be seen more clearly in the turbulence intensity, where the large vertical fluctuations in velocity contribute to a larger kinetic energy for $A \sim 0.68$.

V. MEAN PERTURBATION BEHAVIOR AND GROWTH PARAMETER, θ

In pure RTI the mixing width is found to grow in the self-similar regime [55–57] as $h = \alpha A g t^2$, where h refers to the amplitude of a spike/bubble mixing layer, g is the acceleration of the system (i.e., gravity), and A is the Atwood number. In the RMI self-similar regime, the mixing layer grows proportional to t^θ as $h = (t - t_0)^\theta$, where θ is the growth parameter, though in the case of the BDI it might better be thought of as a “decay parameter.” Due to the impulsive jump velocity imparted on the interface by the blast wave (like that of a shock wave), the RMI relationship for the BDI is considered to be more appropriate [9]. Moreover, in traditional RTI, the acceleration is constant (gravity), where in the BDI the acceleration is time-varying (decreasing).

The value of θ has been found using a power-law fit to $\ln(h) = \theta * \ln(t)$ when growth has saturated, and t_0 is used to indicate a time when the blast wave has stopped interacting with the interface and the growth approaches a self-similar state (i.e., where the logarithm of the mixing width becomes linear, after the inflection point where the growth rate saturates). The growth parameter θ versus time is shown in Fig. 7. For the lower Atwood number of 0.22, $\theta = 0.34 \pm 0.02$. (Simulations from a digital twin of the current facility show reasonable agreement with experiments, $\theta \approx 0.297$ [9].) For the high Atwood number of 0.68, θ has been found to be 0.50 ± 0.05 , thus indicating more potential for mixing [58]. The planar case would develop more, where in this diverging case there is confinement in the z -direction, volumetric changes from expansions, and interface stretching from the diverging geometry. (The surface area of the interface increases by 18% through its trajectory, based on facility geometry.) The Bell-Plesset effect reduces mixing in a diverging environment [59,60], though its effect is negligible after the interface translation ceases (2 ms), long before mixing width growth saturates (6–8 ms). (This would have a stronger effect if the interface were closer to the explosive.) Aside from the decay exponent, the major outcome to note from the Mie scattering visualizations is that inherent morphological differences are present as a result of variation in density ratio [48,49]. From these estimates, comparisons of the value of the decay parameter can be compared to the time-varying theta according to the methods of Thornber *et al.* [58] as

$$\theta^{-1} = 1 - \frac{\ddot{h}h}{\dot{h}^2}. \quad (1)$$

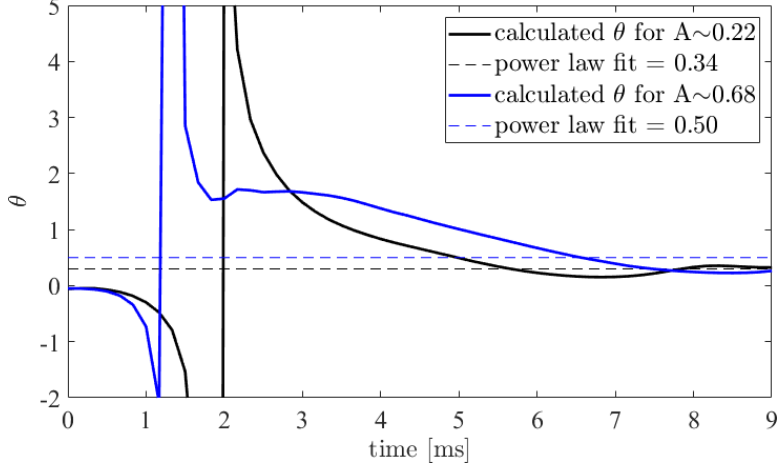
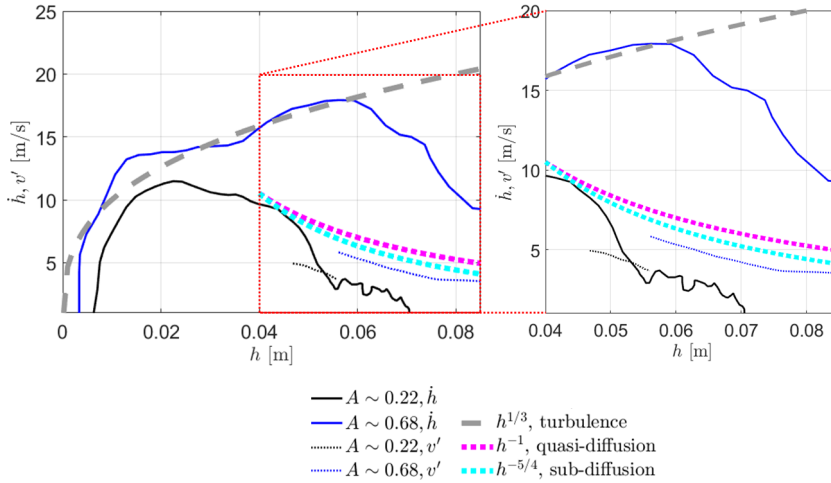

 FIG. 7. Growth parameter θ vs time.

Figure 8 considers mixing width growth rate $\dot{h}$ and velocity fluctuations v' as characteristic velocity scales, and it compares their evolution with mixing width h as a characteristic length scale. These are compared to scaling relations from group theory analysis [21,61], specifically for quasidiffusive and subdiffusive mixing, to check the nature of mixing. We find that $\dot{h}$ (velocity scale) decreases much faster with h (length scale) when compared with quasidiffusion and subdiffusive regimes. However, if we use the velocity fluctuation v' (obtained from PIV) as the velocity scale instead of $\dot{h}$, we find a better match with the subdiffusive scaling, which is usually associated with mixing driven by variable accelerations typical to explosions and blast waves, resulting in large velocities and more intense fluid motions at smaller length scales [21,61]. Velocity fluctuations are discussed in greater detail in the next section.


 FIG. 8. Characteristic velocity vs length scales, $\dot{h}$, v' vs h , along with fiducial lines for classical turbulent, quasidiffusive, and subdiffusive mixing.

VI. EVOLUTION OF VORTICITY

Fundamentally, the BDI, or the growth of the perturbations along the interface, is due to the deposition of vorticity, that is, the misalignment of the density and pressure gradients. These fluid instabilities contribute to the baroclinic production term in the vorticity transport equation:

$$\frac{D\vec{\omega}}{Dt} = \underbrace{(\vec{\omega} \cdot \nabla)\vec{u}}_{\text{Stretching}} - \underbrace{\vec{\omega}(\nabla \cdot \vec{u})}_{\text{Expansion}} + \underbrace{\nabla \times \left(\frac{1}{\rho} \div \tau \right)}_{\text{Diffusion}} + \underbrace{\frac{1}{\rho^2}(\nabla \rho \times \nabla p)}_{\text{Baroclinic}}. \quad (2)$$

In terms of the vorticity dynamics, i.e., deposition of baroclinic vorticity, the instability growth and mixing can be viewed in terms of two distinct components responsible for subsequent mixing. That is, baroclinic *deposition* of the vorticity coupled with *reorganization* and stretching of the same (due to the inherently unstable nature of the vortex sheets). Further, the vorticity deposition itself occurs in two different phases, the relative strength between which can have implications on the asymptotic interface behavior. The first phase is when the interface first interacts with the shocklike raising-pressure gradient at the leading edge of the blast-wave pressure pulse, like that of RMI.

We refer to this as the RMI-phase (equal to the decay time of the pressure wave), which is much faster than the vorticity reorganization time [$O(1 \text{ ms})$ versus $O(10 \mu\text{s})$], respectively. Further, during this shock-interface interaction, the interface is impulsively accelerated (radially outward in the current case). Immediately following this interaction with the shock-front, the blast wave embodies a rarefaction wave that then interacts with the interface, resulting in deceleration of the interface. This deceleration involves additional baroclinic deposition of the vorticity via RTI over a longer timescale than the RMI-phase. Finally, after the interaction with the trailing rarefaction, the vorticity cascades to smaller wave numbers to drive small-scale mixing of the two fluids until viscosity dissipates the kinetic energy. This last phase is truncated in the current physical setup as the rarefaction wave from the blast wave exiting the facility interacts with the interface corrupting the natural evolution.

The availability of high-speed velocity field information lets us probe this vorticity evolution in a quantitative manner. The progressive evolution of the interface vorticity from one experimental run is presented in Fig. 9. Due to the large radial translation of the interface that occurs between the RMI and RTI phases, the current PIV diagnostics were designed to only observe the RTI phase as the interface moves and stagnates into the field of view at the end of the RMI phase. In other words, only the RTI-type deposition and reorganization are observed here. The interface reaches the field of view with locally positive/negative vorticity, deposited during the RMI phase due to initial misalignment between the density and pressure gradients. With the action of the rarefaction wave on the interface, the RT instability accentuates the small perturbations, and the bubbles and spikes grow as described in Sec. V. Finally, at late times, the breakdown of the vortex sheets and the formation of vorticity clouds is evident. In the current section, we look into the details of these vorticity deposition and evolution behaviors of the BDI using the time-resolved vorticity fields.

A. Qualitative description of vorticity evolution

Figure 9 shows an example of the vorticity field evolution, ω . Since the Mie scattering and PIV measurements were not performed simultaneously (akin to Mohaghar *et al.* [54] and Carter *et al.* [62]), the interface between the heavy and light gas is not immediately evident (as it was in Figs. 4 and 5). At early times [$\sim 1 \mu\text{s}$, Figs. 9(a) and 9(e)], the interface is detectable by the sharp peaks in local enstrophy due to the vorticity deposition during the RMI phase. The sign and magnitude of the vorticity is commensurate with the local misalignment ($^\circ$ -inclination) and orientation of the interface. Additionally, the entire interface is advecting due to the shock-induced impulsive acceleration. At later times, due to the rarefaction of the blast wave, the interface decelerates to rest in the field of view, while the perturbations also grow in extent [Figs. 9(b)–9(d) and 9(f)–9(h)]. The spikes grow as the heavy fluid penetrates the lighter fluid, while the bubbles grow in the opposite

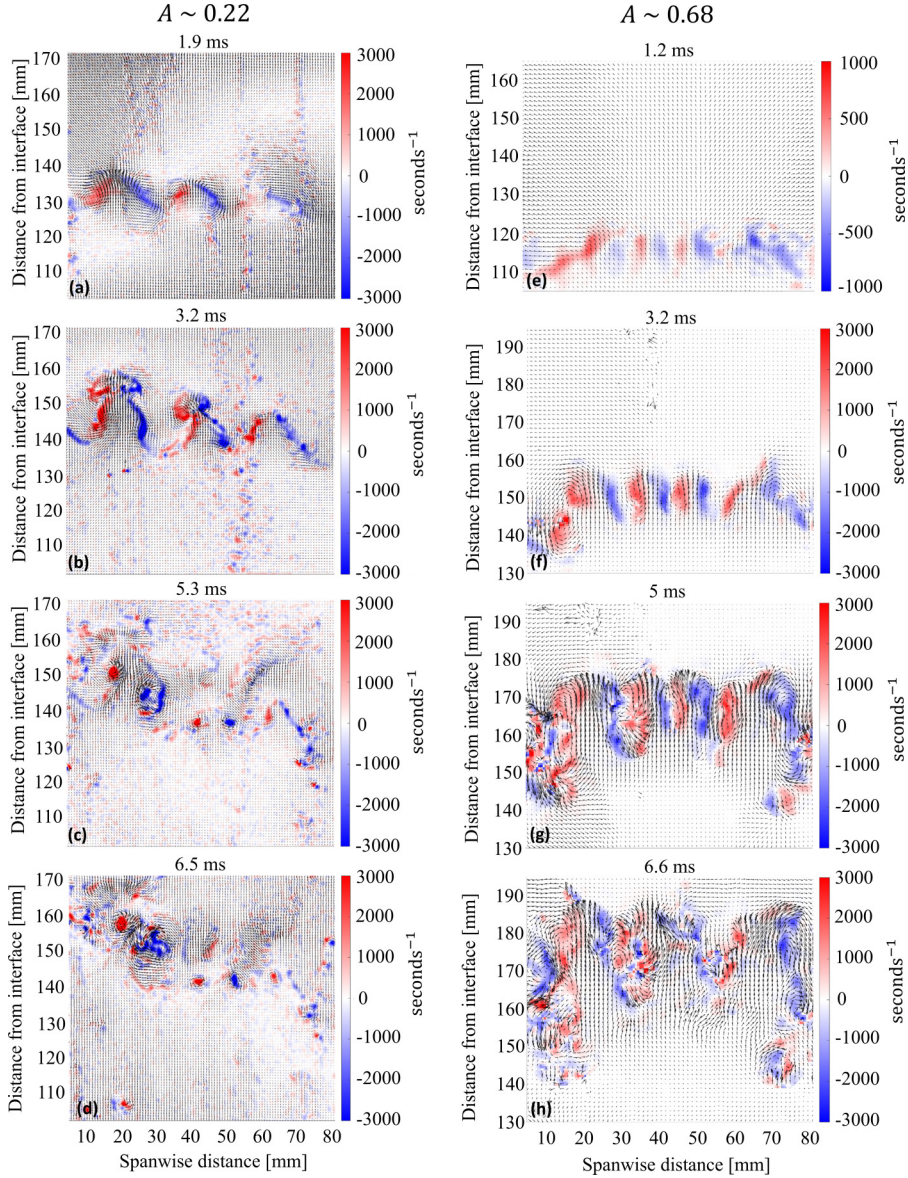


FIG. 9. Progression of perturbation development of vorticity fields for $A \sim 0.22$ (left) and $A \sim 0.68$ (right). The vectors presented here represent the flow field after subtracting the bulk velocity.

direction relative to the mean interface position. Simultaneously, the vorticity on the interface is stretched resulting in intense roll-ups typical of RTI evolution [63]. Due to the nature of these spikes, the penetrating head of the spike tends to begin to shear and a toroidal roll-up vortex forms on sides of the spike. This shows as a pair of counter-rotating dipoles when viewed in a plane, as shown in Fig. 9 [clearest in (b), (c) and (f), (g)]. This tends to "re-accelerate" the bubble and spike causing further stretching of the interface and a larger area for molecular mixing. Additionally, this self-induction can also result in spike detachment at late times, resembling pseudo-ejecta events. This material interface and vorticity breakup lasts even after the interaction with the blast wave is complete.

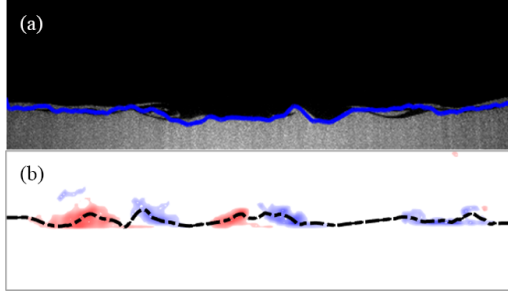


FIG. 10. Part (a) shows interface detection based on pixel intensity from Mie scattering. Part (b) shows interface detection from the peak gradient in vorticity.

B. Vorticity deposition models

Vorticity deposition models are typically used to decouple the compressible flow aspects of shock-driven instabilities (shock-interface interaction, vorticity deposition, shock-refraction, internal shock reflections, etc.) from the subsequent interface evolution that is typically divergence-free. The former can be modeled analytically/empirically [64–66] and the latter, which results in much of the perturbation growth, flow transition, and late time mixing, can be modeled/simulated more easily using incompressible multiphase flow techniques. Many of the assumptions used for these models are valid for the RMI phase of the blast-wave–interface interaction in the current experiments, such as short interaction times (i.e., impulsive acceleration), "regular" shock interaction (i.e., ignoring the reflected shocks and rarefactions), small perturbation amplitudes, impulse density gradients, and single-valued nature of the interface. We use this approach and current measurements to quantitatively investigate the relative strengths of the baroclinic vorticity deposition/generation (from RMI and RTI phases), and the utility of the existing vorticity-deposition models [65,66] for the blast-driven instability.

For this, the shape of the interface is required to deduce the misalignment between the density ($\nabla\rho$) and pressure (∇P) gradient fields. Since the initial interface is not in the FOV (as it is in Mie scattering), we can use the regions of high vorticity at early times to estimate the shape of the interface before significant development has occurred. This assumes that the initial conditions have not had time to develop in the short time it took for the interface to come into the field of view. Since the flow is irrotational everywhere except the interface, a good approximation of the initial conditions can be obtained [see Fig. 10(b)].

This indirect approach of estimating the interface shape is shown in Fig. 10(b), which traces the radial position (y –) of the maximum vorticity magnitude for each azimuthal location (x –). This snapshot corresponds to the time when the interface first enters the FOV completely ($t = 1.2$ ms). Here the vorticity is zero where the interface is parallel to the blast-wave pressure front ($\nabla\rho \times \nabla P \approx 0$). Comparing these early-time vorticity fields to the images captured using Mie scattering [Fig. 10(a)], it is seen that profiles of initial conditions are within the same range.

Two different models of circulation are used to predict vorticity along the interface, which is then compared to the experimental values: (i) a shock-polar model analysis by Samtaney *et al.* [65], hereafter referred to as SRZ98, and (ii) a 1D gas dynamics model by Weber *et al.* [66], hereafter referred to as WCB13. In both models, the estimation of circulation deposited per unit length of the interface (σ) is based on local M of the blast wave and local angle of misalignment, α' , between the interface and the blast wave. For a weak shock traveling from a heavy gas into a light gas, the SRZ98 model estimates the circulation per unit interface length, σ , to scale as

$$\sigma_{\text{SRZ98}} = \left(\frac{2c}{\gamma^{1/2}} \right) \left(1 - \frac{1}{\sqrt{\rho_l/\rho_h}} \right) \left(\frac{4\gamma}{\gamma+1} \right) (M-1) \sin(\alpha'), \quad (3)$$

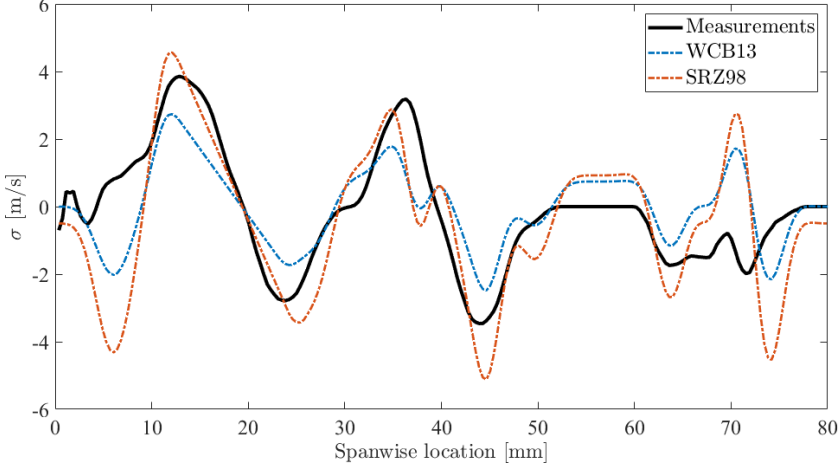


FIG. 11. Comparison of circulation (per unit interface length, σ) to predictions of vorticity from a shock-polar model (SRZ98) and a 1D gas dynamics model (WCB13).

where c is the local speed of sound, γ is the ratio of specific heats, ρ is density for light (l) and heavy (h) gas, M is the Mach number, and α' is the local misalignment between the interface and the blast wave. Thus, using only properties of the gas, the blast wave, and interface shape, the final vorticity after the RMI phase can be predicted from initial conditions. To compare with the estimates, the circulation per unit length, σ , can be calculated from the experimental vorticity field using Stokes theorem ($\sigma = \int \omega dy$).

On the other hand, the WCB13 model uses a 1D gas dynamics approach to predict the vorticity per unit length of interface, which for the current case is given by

$$\sigma_{\text{WCB13}} = \frac{1}{2} v_j A \tan(\alpha'). \quad (4)$$

where v_j is the jump velocity, A is the Atwood number, and α' is again the local misalignment between the interface and the blast wave.

Unlike SRZ98, the WCB13 model requires the jump velocity across the interface, v_j , to be prescribed to estimate the circulation deposited. This is directly measured from experimental data. The full details of both models can be found in the original works [65,66], and their application to a shock-wave can be found in Ref. [62].

For the current case, the blast wave is assumed to be approximately planar (ignoring blast-wave curvature effects), since the wavelength that the pressure jump would interact with is small relative to the overall curvature. With this, the local values of M , α' , σ , and v_j can be extracted as described before. Figure. 11 shows an illustrative comparison between the observed and estimated profiles of σ on the interface from a single run. The results shown are for the earliest time, $t = 1.2$ ms, where the interface is captured in the FOV, which is immediately after the RMI-interaction but shortly before the RTI-induced deceleration has significantly altered the interface. It can be seen that both the models capture the relative orientation of the local vorticity. The WCB13 model tends to generally underestimate the local vorticity while the SRZ98 model shows extrema while agreeing with the measurements elsewhere. More critically, the current analysis shows that the blast-wave vorticity deposition on the interface can indeed be decoupled into an RMI-phase and RTI-phase, where the former is more amenable to existing analytical models developed for shock-interface interaction applications. (Previous authors have also decoupled the instability phases [9].) This is due to the sensitivity of these models to α' , and a higher degree of uncertainty in the way this quantity is estimated in the current work. The WCB13 model is generally considered to be more applicable [62], since this prescribes the jump velocity without having to make assumptions on the nature of

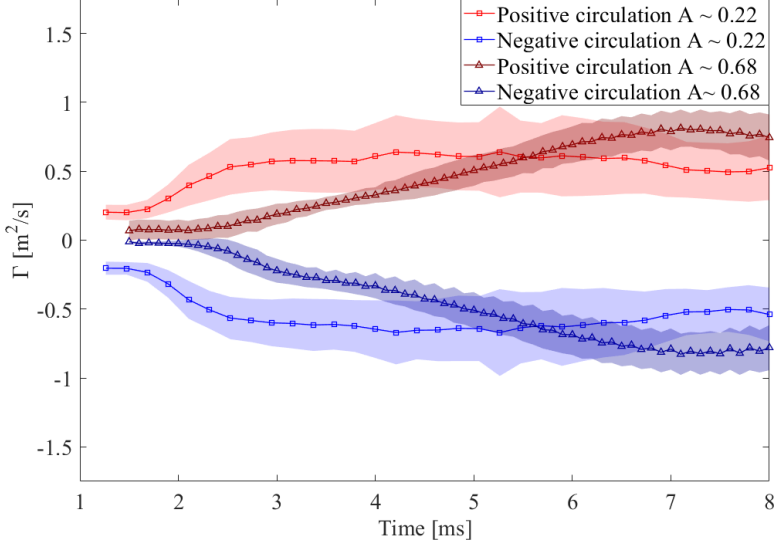


FIG. 12. Comparison of the positive and negative components of circulation for $A \sim 0.22$ and $A \sim 0.68$.

the impulsive acceleration. This additional information, however, requires the use of other models when the same is not available *a priori*.

VII. CIRCULATION AND TURBULENCE INTENSITY

This section presents the development of the flow through temporal evolution of circulation and turbulence intensity (TI) between the two cases. These quantities are of particular interest because of the implications of how energy is transferred from the momentum of the blast wave to the interface for the different density ratios. From the Mie scattering images, it was evident that smaller density ratios result in much more symmetry between the spike and bubble structures than at larger density ratios, where spikes penetrated much deeper into the lighter fluid and were much thinner than bubbles of light fluid. In this section, a more quantitative analysis is shown to support this qualitative assessment.

A. Comparison of temporal evolution of field circulation

After initial deposition from the RMI phase, the interface evolves under this vorticity while also being subjected to the deceleration from the rarefaction wave (RTI phase). Additional circulation (positive and negative orientations) is added to the interface over this duration and can be investigated using the current datasets as an ensemble average for each instance in time (phase/ensemble averaging). The total positive and negative circulations of the flow can be computed by conditionally integrating the vorticity fields across the mixing region for positive and negative vorticity, and ensemble averaging the experiments in phase, correlating run times,

$$\Gamma^\pm(t) = \left\langle \iint \omega^\pm(x, y, t) dy dx \right\rangle_i. \quad (5)$$

This evolution of $\Gamma^\pm$ is shown in Fig. 12, as the circulation evolves in time. The evolution of positive/negative circulation is equal and opposite (net circulation sums to zero), which demonstrates strong symmetry in the structures where the counter-rotating vortex sheets are equal in magnitude and opposite in orientation. The evolution starts with a nonzero initial circulation due to the aforementioned RMI interaction.

The positive and negative circulation grows until $t \approx 3$ ms for $A \sim 0.22$, and later to ≈ 7 ms for $A \sim 0.68$, indicating that the rarefaction continues to deposit additional vorticity. This circulation contribution occurs during the acceleration RTI phase, and we see that the RTI phase embodies the maximum growth in circulation. Beyond this peak, the magnitude of total interface circulation does not change significantly as the interface vorticity redistributes between different lengthscales, as was evidenced earlier in Fig. 9. The mixing continues as the vortex dipoles continue to stretch the interface, until the viscous forces dissipate the energy as heat. This is an important feature of the BDI, where further mixing development is present beyond what might be predicted by the individual RMI/RTI components.

The structures that form for $A \sim 0.68$ exhibit much less symmetry; they are much thinner and create ligament/finger type structures that do not begin to roll up until they have penetrated further than their low Atwood counterpart. Because of this, the circulation is seen to grow slower for high Atwood number case, but reaches to a similar extent by the end of the experimental time. Note that the field of view is in the same location for both data sets. Since the structures for $A \sim 0.68$ developed to a greater extent and left the field of view sooner, the data could only be evaluated up to that point. The circulation is only presented where fields from both the cases could be evaluated.

The circulation for each case can be seen in Fig. 12. Another important point to note is the significant growth in the high- A circulation, where significant vorticity is generated via secondary instabilities (as also suggested by temporal simulations [67]).

The difference in the location of this peak is due to the fact that for $A \sim 0.68$, the Mach number is $\sim 5\%$ smaller. The increased value in the peak for $A \sim 0.68$ is within one standard deviation, suggesting that the morphological shape of the structures (a strong function of A) contributes more to the turbulence intensity than to the circulation. (Turbulence intensity is discussed next.) Since the spikes from the larger Atwood number are thinner with less roll-up, the order of magnitude of the swirling strength is the same. At the low $A \sim 0.22$, features (bubbles and spikes) tend to exhibit much more symmetry, while at $A \sim 0.68$, spike features develop that were much larger than their surrounding bubble structures; some spikes were not only much larger than bubbles, but even much larger than other spikes. This is more easily visualized from the full-field Mie scattering experiments (see Figs. 4 and 5). The comparison of circulation (shown in Fig. 12) shows that there is not a large difference in swirling strength between the two different Atwood numbers. This is likely due to the fact that the Atwood number has a greater influence on the shape of the structures, but the rotational energy deposited is less affected. These structures have a larger vertical fluctuating component that is seen to contribute more to turbulence intensity, as the circulation is only 15% larger for the high Atwood case.

B. Comparison of temporal evolution of turbulence intensity

For variable density flows, several methods have been used to quantify the energy scales [68]. The fluctuating quantities indicate how much kinetic energy in the flow is decoupled from the mean [52,53]. Although this flow is highly transitional, if and when it was to become turbulent this is the energy that would lead to development of a wide range of energy scales; the velocity fluctuations are a good indicator of how active the mixing is. Strictly speaking, this is a Favre-averaged quantity, which requires information about the density. Without full knowledge of turbulence or satisfaction of this criteria, it can still provide a useful measure of turbulence intensity [69,70]. With only PIV-velocity measurements, only approximations of the kinetic energy can be made. Here this estimate will be called “turbulence intensity,” TI, so as not to confuse the two quantities. The TI is calculated from the vertical and spanwise fluctuations in velocity as

$$TI = \frac{1}{2}(v'^2 + 2u'^2). \quad (6)$$

The factor of 2 accounts for the out-of-plane fluctuations, which are estimated to be the same order of magnitude as the spanwise fluctuations, since they are both orthogonal to the propagation of the blast wave. This flow is highly anisotropic, with spanwise fluctuations averaging 10 – 30 %

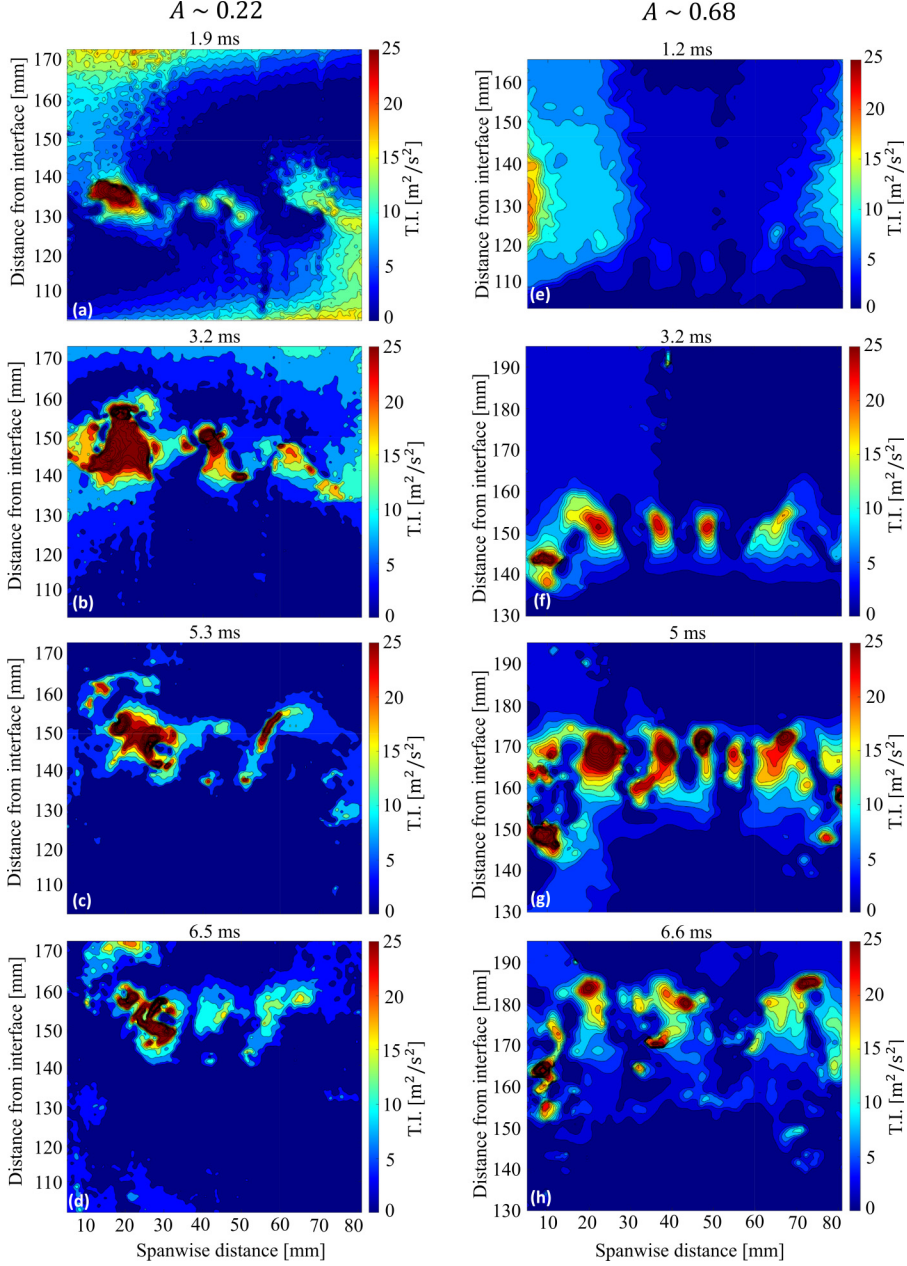


FIG. 13. Color maps with select times of the progression of turbulence intensity for $A \sim 0.22$ (left) and $A \sim 0.68$ (right).

of the vertical fluctuations. Applying Eq. (6) to the fluctuations of the velocity fields generates a matrix of TI values at each point in time. This is plotted for several times in a color-map in Fig. 13. One can clearly see the concentrated TI and where the intensity is concentrated in the penetrating spikes. These are active sites for mixing and follow the boundaries of the two fluids well (seen in Mie scattering), which identifies structures very clearly.

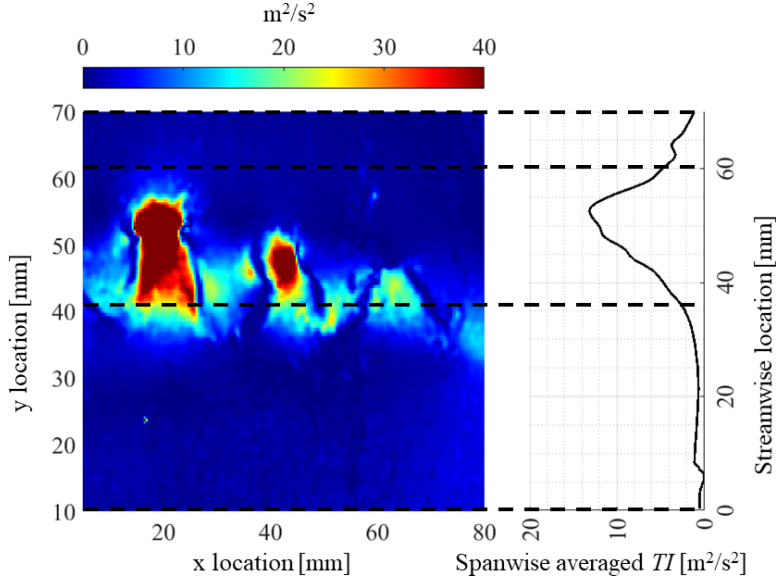


FIG. 14. On the left is a color-map of TI from $A \sim 0.22$. On the right is an example of the mixing width that comes from the resulting spanwise average of TI (averaging rows of camera pixels).

These regions of fluid penetration in Fig. 13 correspond to the spike features outlined by vorticity in Fig. 9 and show that the fluctuating quantities are significantly larger inside these spike features than any other region in the flow. Further, it is clear that the constant-density regions of the flow and regions on the interface without a mismatch between the blast wave and interface sustain no vorticity deposition and turbulence intensity from the blast wave. (Figs. 9 and 13 illustrate the bulk flow, however subsequent analysis subtracts the mean velocity to investigate the interfacial mixing [71].)

A spanwise average (averaging camera pixel rows) is shown on the right side of Fig. 14 which can be used to identify the mixing width [72]. Progressions of TI in time are plotted over the top of each other; this time development is shown in Fig. 15. In time, the peak of the TI decreases

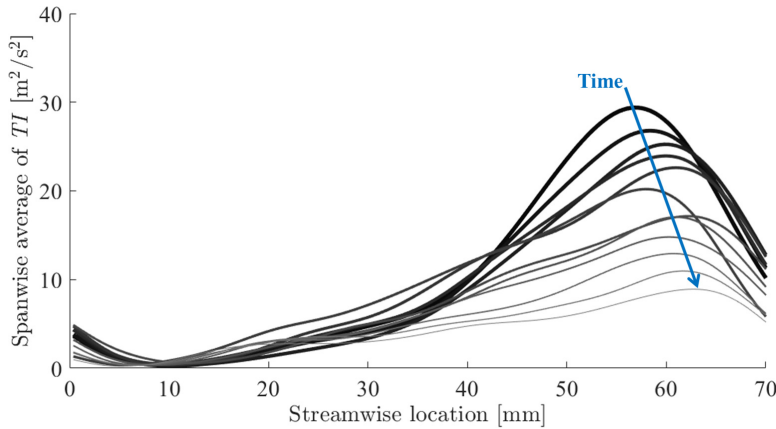


FIG. 15. Ensemble averages of spanwise averaged TI. The TI peak within the mixing region decays in time, and the width of the peak broadens (directly related to the mixed width growth). This is done for both A 's; this plot shows the evolution from $A \sim 0.22$ as a demonstration.

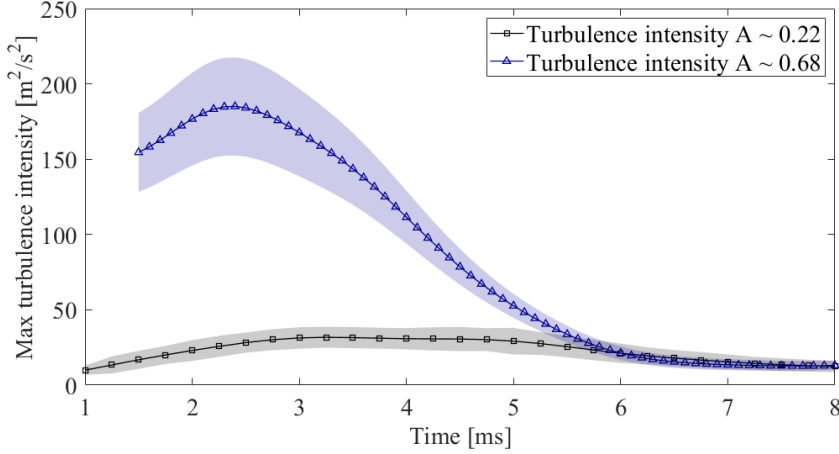


FIG. 16. Comparison of the growth in turbulence intensity for $A \sim 0.22$ and $A \sim 0.68$.

and broadens. This peak broadening exemplifies the TI dissipating through the system. As mixing develops further, kinetic energy dissipates via viscous diffusion.

From the peak values in Fig. 15, the turbulence intensity is plotted as a function of time. Note that early times when TI is increasing have been excluded due to the extensive overlapping of plots, which introduced ambiguity in the data interpretation. Figure 16 shows the progression of turbulence intensity for the two Atwood numbers. Considering $A \sim 0.22$, a fairly gentle rise is seen, a modest peak, and then a slow decay. The $A \sim 0.68$ case behaves somewhat differently. As the interface comes into the field of view, there is a much steeper rise in turbulence intensity, followed by a much steeper decay before reaching a coasting phase. This peak in turbulence intensity is attributed to the early formation of large structures (much larger than those compared to the $A \sim 0.22$ case), whose vertical fluctuations in velocity contribute much greater to the production of turbulence, especially at early times.

Comparing turbulence intensity across the two cases, the larger Atwood number produces larger fluctuations in velocity, particularly in the vertical direction. This peak in TI is nearly three times the peak of the lower density ratio, and the main peak is sharper and more short-lived than the lower Atwood case as well. This early peak is seemingly a result of the larger structures that develop in $A \sim 0.68$, which are thinner and carry a disproportionate fluctuating component to the vertical velocity relative to the mean/bulk velocity of the interface. After this quick decay, the bulk of the dissipation begins happening at the same time (3 ms) but at a faster rate than the lower Atwood case.

VIII. TRANSITIONAL FLOWS AND THE NATURE OF THE NONLINEAR STAGE MIXING

Particularly in the context of ICF, the progression of fine perturbations to the nonlinear stage (and ultimately to turbulence) remains highly unanswered [46]. This nonlinear stage is highly dependent on initial conditions and is ultimately what leads to turbulence. The mixing transition is not universal and can manifest in a manner that is relatively inconsistent across mixing regimes depending on the local Reynolds number, and a turbulent state is not an “achieved state” in the sense that it’s sudden, but instead it is a progression that may be fast or slow, and can occur at a range of Reynolds numbers. The Reynolds number is momentum based, defined based on the mixing width as

$$\text{Re} = \frac{h\dot{h}}{\nu_{\text{mix}}}, \quad (7)$$

where ν_{mix} is the mixture viscosity, calculated as $(\mu_1 + \mu_2)/(\rho_1 + \rho_2)$ [73]. Since the Reynolds number is calculated from the developing mixing width, it changes in time.

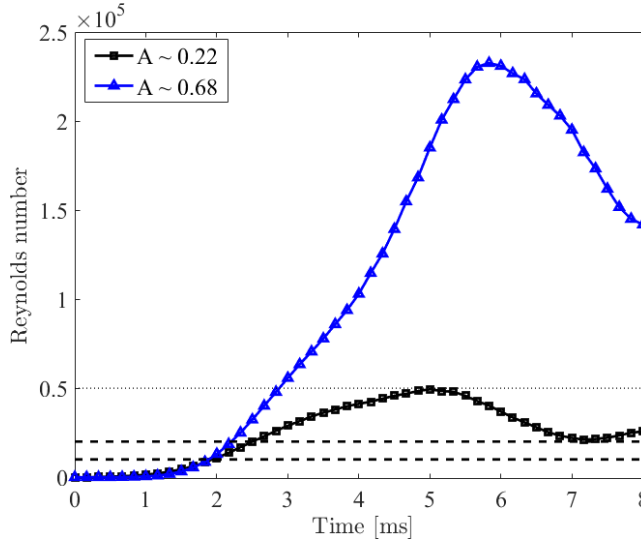


FIG. 17. The Reynolds number is plotted as a function of time. Dimotakis' criteria for transition to turbulence are shown as dark dashes for reference, and the lighter dotted line indicates where the mixing transition is seen visually in experiments.

Dimotakis [74] postulates a range of Reynolds numbers beyond which (mixing) transition to fully developed turbulent flow can occur, $\gtrsim (1 - 2) \times 10^4$. This criterion need not directly apply to the instability-driven mixing transitions being discussed in this paper. However, it remains the most widely referenced metric for flow transitions/mixing transitions in general. It is seen in several shock tube studies that mixing transition occurs at Reynolds numbers above 5×10^4 [75], which is at a significantly higher Reynolds number than Dimotakis [74]. It is especially true for transient/developing flows, particularly highly anisotropic flows, that a transition to turbulence corresponds better to a Reynolds number of 5×10^4 . Figure 17 shows the Reynolds number as a function of time for the two cases. Dashed lines outline Dimotakis' suggested range for transition, and a lighter dotted line indicates where impulsively accelerated anisotropic flows more commonly realize transition, likely due to the fact that the Reynolds number based on $h\dot{h}/\nu$ may not accurately represent the time-dependent Reynolds number, but instead represents the maximum Reynolds number that *may* be achieved [75]. It is also important to remember the time criteria required for transition. For this, enough time has to pass so that the inertial subrange develops a sufficient range of scales so that the viscous subrange is well above the Kolmogorov dissipation scale, which is established when the viscous diffusion length scale becomes larger than the inner viscous scale [29]. This criterion is purely time-dependent, however both of these criteria need be met for transition to occur.

Throughout the course of the experiments, structures in the lower Atwood case are very much preserved, and a fully turbulent state is not really achieved. This suggests that a larger Reynolds number is required for transition to turbulence in this type of flow. The larger Atwood number of 0.68 is clearly turbulent by the midpoint of the experiment, as suggested by the Reynolds number. However, $A \sim 0.22$ does not exhibit a turbulent state. A point to remember is that achieving a certain Reynolds number does not guarantee turbulence. This is especially true for transient events with time-dependent Reynolds numbers, where just because the criterion is reached, a turbulent state is not guaranteed. This is not surprising, given the nature of the transient passing blast wave. Mixing is initiated, but without a significant density gradient or more driving energy, this mixing is not sustained or increased beyond a certain point.

IX. CONCLUSIONS

The BDI is studied in a diverging cylindrical geometry using high-speed temporally resolved PIV. First, the perturbation growth was studied and characterized via the mixing width and mixing width growth rate. From this, the growth parameter (θ) was found to be consistent with that of the digital twin of the same geometry [9]. This not only validates the experimental scalability, but it also underscores the notion that the larger density contrast provides more potential for mixing, which has direct implications to ICF capsule design and mixing suppression.

From this, PIV was focused on where the interface settled after the impulsive jump, and time-resolved vector fields provided a basis for calculating turbulence and vorticity characteristics. Specifically, the vorticity imparted on the interface obtained experimentally was compared to that of two models to predict the same, finding both models to be adequate predictors, also validating the PIV vorticity calculations. Agreement with these models validates the notion that evolution of the BDI can be decoupled into RMI/RTI-phases. The long-lasting RTI-phase of the BDI exhibited the bulk of vorticity development. From these vorticity fields in time, full-field circulation was plotted as a function of time, showing prolonged growth for $A \sim 0.68$.

Next, an estimate of the turbulence intensity is presented. The kinetic energy in the flow field was highly concentrated around the spike features, showing regions of entrainment. This provided a good estimator for the mixing width (in accordance with Mie scattering estimates). The maximum/integrated kinetic energy as a function of time showed a sharp rise as the interface entered the field of view, with the high- A case rising to a much greater extent, and then decaying after the midpoint of the experiment (~ 4 ms), corresponding to the inflection point of the mixing width, and where the mixing width growth rate begins to decay. The prolonged growth associated with a larger density difference suggests that density contrast strongly influences TI, as seen where a larger A exhibited sharper peaks due to thinner, deeper penetrating spike structures.

Lastly, a mixing transition is seen to occur at Reynolds numbers significantly larger than the Reynolds number for the mixing transition in shear and other flows, and was visually upwards of 5×10^4 . This is consistent with highly anisotropic/instability studies, where fully developed mixing requires additional nonlinear processes beyond the onset of turbulence. Moreover, it is also likely due to the fact that the Reynolds number used here is a momentum-based Reynolds number, and information on the concentration may be more appropriate for estimating transition. These results provide a new benchmark for numerical and theoretical studies for the instability mixing transition.

Overall, this novel experimental investigation of the combined RMI/RTI is the first to temporally resolve velocity fields in blast environments, and it helps characterize our understanding of the BDI, which has a direct relation to ICF inefficacy and SN ejecta events. It gives insight into BDI mechanisms and clarification between the interplay of RMI and RTI, advancing the understanding of coupled instabilities.

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