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Prediction of Cavitation Bubble Dynamics in a Tube with High-Speed Internal Flow

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ABSTRACT: Cavitation in confined internal flows plays a crucial role in a wide range of natural and engineering systems. In the present work, a numerical model, based on a unified theory for bubble dynamics, is developed to simulate the morphological evolution of a single cavitation bubble traversing a microchannel nozzle. By incorporating ambient pressure and velocity fields extracted from computational fluid dynamics (CFD) simulations, the model predicts bubble growth, collapse, and rebound in non-uniform flow environments, thereby providing insight into the effects of key governing parameters on bubble dynamics within micro-channel nozzles. The results indicate that the initial radial position of the bubble, the initial ambient pressure, and the initial bubble radius all have pronounced effects on bubble morphological evolution by modifying the effective expansion distance and effective expansion duration. These changes, in turn, affect both the intensity and the spatial distribution of the transient impact loads imposed on the micro-channel walls.

KEYWORDS: Bubble dynamics; micro-channel nozzles; CFD simulations

1 Introduction

The formation of vapor-filled cavities within a liquid, a phenomenon known as cavitation, is initiated when the local pressure falls beneath the saturation vapor pressure [1]. The ubiquity of this process spans a vast range of applications, from biological systems and hydraulic engineering [2–4] to advanced technologies like ultrasonic cleaning [5,6], materials synthesis, drug delivery, and cardiovascular medicine. In many such contexts, cavitation is induced hydrodynamically as the liquid traverses geometrically confined passages with varying cross-sections, exemplified by venturi tubes [7–9], vasculature [10], and diffusers [11]. The dynamics of bubbles in these confined environments deviate significantly from those in a free field. The presence of nearby solid boundaries leads to non-spherical growth and collapse [12–15], the formation of high-velocity liquid jets [16–18], and the eventual transformation into toroidal structures [18–20]. Owing to the deleterious effects of the violent bubble collapse and the resulting jet impact, which can lead to severe material damage [21], the study of cavitation dynamics in the tube remains a subject of intense scientific investigation.

A fundamental prerequisite for predicting the hydrodynamic behavior of confined cavitation jets is the detailed characterization of bubble formation, migration, and collapse. In the field of bubble dynamics, the coupling effect of bubble oscillation and migration constitutes a central subject of investigation. Traditionally, the Rayleigh–Plesset (RP) equation [22] has served as the governing equation for predicting spherical bubble oscillations, with numerous studies extending its application to cavitation in the tube [5,23,24]. Nevertheless, this approach encounters a critical limitation: the maximum migration velocity in microchannel can attain hundreds of meters per second, comparable to the sonic speed of the fluid. This high Mach number renders the incompressibility assumption unrealistic, thereby inducing non-negligible errors. While the Gilmore [25] and Keller–Miksis [26] equations were introduced to resolve this by accounting for fluid compressibility, a more holistic approach is often required. Therefore, this study employs a unified theoretical model established by Zhang et al. [27,28] as the governing framework. Formulated from rigorous mathematical derivations, this model comprehensively accounts for compressible bubble pulsation, migration, phase transition, ambient flow conditions, fluid viscosity, and surface tension.

Nevertheless, the direct experimental characterization of cavitation in microchannel faces significant impediments [29]. The minute scales of microchannels, compounded by high flow velocities and the ephemeral nature of the jetting process, severely limit experimental observation. The flow complexity is further exacerbated by turbulence initiated by the high-speed injection, confounding direct measurements. Within intricate tubes, bubble dynamics become even more non-trivial, featuring events like bubble splitting and asymmetric jet formation [30,31]. In light of these experimental impediments, numerical simulation presents a powerful alternative. Its ascent as the preferred method for studying jet dynamics is directly linked to the rapid advancements in Computational Fluid Dynamics (CFD) [32,33]. A distinct advantage of simulation is the capacity to resolve fine-grained details of the flow field. Such details are often inaccessible to experimental measurement, providing a strong impetus for extensive research in this domain. The present study introduces a numerical model founded upon the unified bubble dynamics equation. This model explicitly considers the pressure differential inside and outside the bubble. And the equation is solved via the finite difference method. The requisite ambient velocity and pressure fields that drive the model are pre-computed using CFD. The overarching goal of this research is to furnish a holistic understanding of bubble dynamics in the microchannel, paving the way for advancements in both fundamental science and engineering applications.

2 Numerical Implementation

2.1 Equations of Bubble Motion

Within the framework of potential flow theory and law of conservation of energy, Zhang et al. [27] developed a theoretical model that couples bubble oscillation, migration, and phase transition. The resulting governing equation for bubble oscillation is expressed as follows:

$$\left(\frac{C - \dot{R}}{R} + \frac{d}{dt} \right) \left[\frac{R^2}{C} \left(\frac{1}{2} \left(\dot{R} - \frac{\dot{m}}{\rho} \right)^2 + \frac{1}{4} v^2 + H \right) \right] + \frac{d}{dt} \left(\frac{R^2 \dot{m}}{\rho} \right) = 2R\dot{R}^2 + R^2\ddot{R} \quad (1)$$

Under the assumption of no phase change, Eq. (1) can be expanded as follows:

$$\left(1 + \frac{\dot{R}}{C} \right) H + \frac{R}{C} \dot{H} + \frac{1}{4} \left(1 + \frac{\dot{R}}{C} \right) v^2 + \frac{R}{2C} v\dot{v} = \frac{3}{2} \left(1 - \frac{\dot{R}}{3C} \right) \dot{R}^2 + \left(1 - \frac{\dot{R}}{C} \right) R\ddot{R} \quad (2)$$

$$H = \int_{P_a}^{P_b} \rho^{-1} dp \approx \frac{P_b - P_\infty}{\rho} \quad (3)$$

where P_a means the ambient pressure at the bubble center; P_∞ denotes the ambient pressure surrounding the bubble, and P_b represents the pressure exerted by the flow field on the bubble surface.

$$P_b = P_0 \left(\frac{R_0}{R} \right)^{3\gamma} + P_v - \frac{2\sigma}{R} - \frac{4\mu\dot{R}}{R} \quad (4)$$

where σ means the surface tension coefficient, γ denotes the polytropic exponent, μ expresses the liquid viscosity and R_0 corresponds initial bubble radius. Consequently, the crux of solving the governing equations lies in the determination of v and P_∞ . Among these, the liquid pressure P_∞ is of paramount importance, as its variations directly dictate the dynamics of bubble growth and collapse. However, a precise analytical or empirical expression for P_∞ remains elusive. While certain simplifying assumptions have been made, such as the functional form $p_\infty = p_o(1 + A \sin \omega t)$ adopted by Shima and Tsujino [34], these approaches are neither derived from experimental measurements nor from high-fidelity simulations of realistic flow fields [35]. Instead, they typically postulate a small-amplitude harmonic pressure fluctuation. Such simplified models inherently struggle to elucidate complex cavitation phenomena. Therefore, this paper leverages the OpenFOAM platform [36] to numerically compute the pressure field P_∞ . Table 1 provides a detailed listing of these parameters and their values. All numerical simulations were conducted using CavBubbleFoam, a customized OpenFOAM solver developed based on compressibleInterFoam [16]. Turbulence was achieved using the RNG (Re-Normalisation Group) k - ϵ model. The pressure-velocity coupling was handled by the PIMPLE algorithm (PIMPLE merges the controls of PISO and SIMPLE).

Table 1: Physical properties of the working fluid and key parameters.

Parameter	Value
$\rho/(\text{kg} \cdot \text{m}^{-3})$	998.203
$\sigma/(\text{N} \cdot \text{m}^{-1})$	0.075
$C/(\text{m} \cdot \text{s}^{-1})$	1500
$\mu/(\text{Pa} \cdot \text{s})$	0.001
γ	1.4
$R_0/(\mu\text{m})$	0.5

The methodology employed in this study proceeds as follows: First, the spatial profiles of pressure, denoted as $p(x)$, and velocity, denoted as $v(x)$, are extracted from the computational fluid dynamics (CFD) results along a predetermined path within the given flow field. These spatial distributions are subsequently mapped into time-domain histories—specifically, the pressure history $p_\infty(t)$ and the velocity history $u_\infty(t)$ —via velocity integration. Next, $p_\infty(t)$ and its temporal derivative, $\dot{p}_\infty(t)$ are incorporated into the unified bubble oscillation equation as the background driving pressure, while $u_\infty(t)$ is superimposed onto the unified bubble oscillation equation as the background migration velocity. Finally, the adaptive Runge–Kutta method is employed to solve the highly nonlinear second-order ordinary differential equations, thereby yielding the time histories of the bubble radius $R(t)$, the bubble center position $x(t)$, and the pressure at the observation point $p_{out}(t)$, capturing the entire bubble lifecycle from initial growth to ultimate collapse. Consequently, this framework enables the highly nonlinear dynamics of local cavitation bubbles to be driven solely by the pressure–velocity evolution of a single-phase flow, effectively circumventing the need for explicit two-phase numerical simulations.

2.2 Numerical Implementation

As shown in Fig. 1a, the present study utilized an axisymmetric simulation to analyze the cavitation bubble dynamics in the convergent-divergent microchannel nozzle. As illustrated in Fig. 1b, a symmetry boundary condition is applied along the axis (indicated by the red horizontal dashed line). This substantially enhances computational efficiency, facilitating the subsequent high-frequency sampling of pressure and velocity data, as well as comparative analyses across various operating conditions. Simplified boundary conditions are adopted for the unidirectional flow calculations. Specifically, the front and rear surfaces are designated as axisymmetric wedge planes. A pressure inlet is applied to the left boundary, and a pressure outlet with a specified backpressure (P_{back}) is assigned to the right boundary to simulate the ambient pressure of the external chamber or test vessel. Because the internal dimensions of the nozzle are markedly smaller than those of the downstream cavity, a fully developed stable flow is assumed at the outlet cross-section. All other boundaries are defined as no-slip, adiabatic walls ($u = 0$).

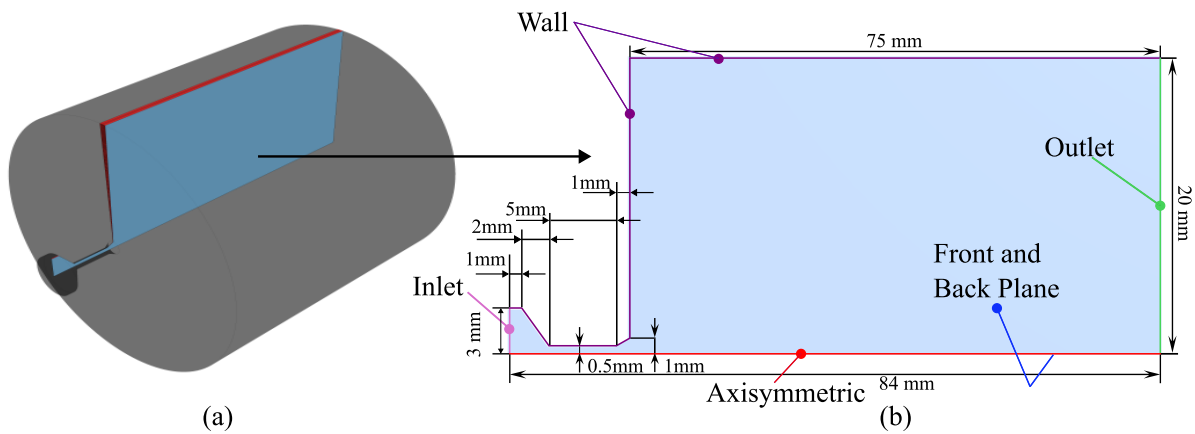


Figure 1: The computational modeling of converging-diverging microchannel nozzle: (a) Axisymmetric model with one degree in the circumferential direction, (b) Boundary conditions and geometric schematic.

Additionally, Fig. 2 illustrates the hexahedral mesh utilized throughout the computational domain, featuring local mesh refinement in the focal regions.

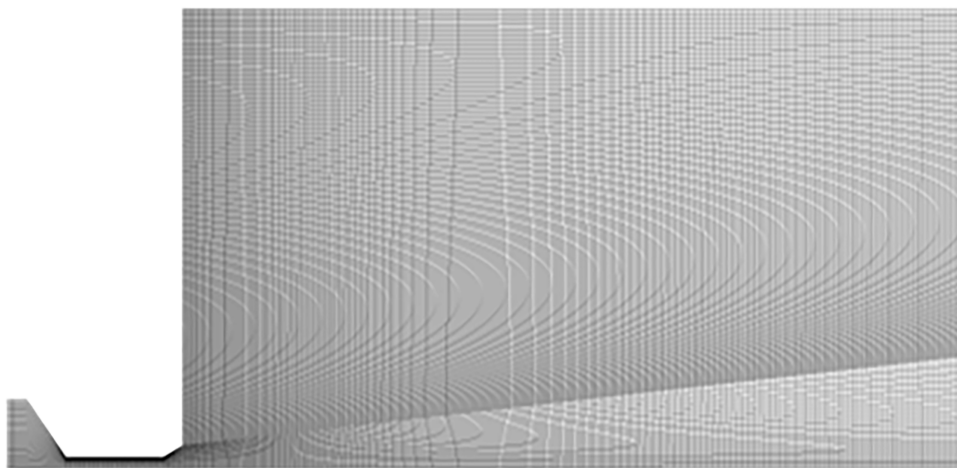


Figure 2: Meshed flow domain.

A spatial convergence study was conducted to verify grid independence. To accurately capture the steep flow gradients, a locally refined uniform mesh was employed near the throat region, transitioning to a coarser grid in the far-field. Three different grid resolutions were evaluated within the refinement zone: $\Delta x = 17$, 8, and 4 μm . As illustrated in Fig. 3, the axial static pressure distributions are highly sensitive to grid resolution in the contraction zone ($x \approx 5\text{--}10$ mm). The coarse mesh ($\Delta x = 17$ μm) severely overpredicts the magnitude of the negative pressure peak. As indicated by the arrow, progressive mesh refinement effectively mitigates this numerical over-prediction, driving the pressure minimum toward convergence. Downstream of the throat ($x > 15$ mm), all three grids produce identical pressure recovery profiles. While a minor deviation persists at the absolute minimum between the 8 and 4 μm grids, the 8 μm mesh sufficiently resolves the critical pressure gradients. Therefore, to achieve an optimal balance between numerical fidelity and computational efficiency, the $\Delta x = 8$ μm grid was adopted for all subsequent simulations.

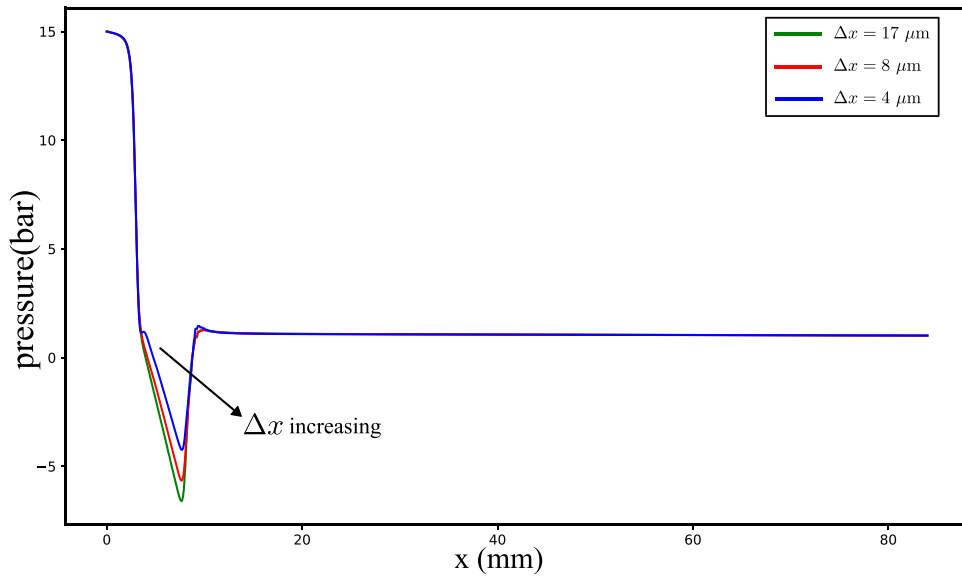


Figure 3: Mesh convergence study: Axial static pressure distributions along the microchannel for different grid resolutions ($\Delta x = 17$ μm , 8, and 4 μm). The sharp negative pressure peak near the throat region is accurately captured to verify grid independence.

As shown in Fig. 4, we compared our high-speed experimental recordings with the numerical predictions. The sequence captures the entire evolution of the bubble from its maximum equivalent radius to its violent collapse near the wall. The excellent morphological agreement between the experimental observations and the CFD results confirms the reliability of our numerical framework.

2.3 Modelling the Pressure Outside of the Bubble

In the context of the microchannel nozzle, assuming the CFD-computed pressure field accurately represents the flow domain, a moving bubble encounters continuous pressure variations. The governing bubble equation dictates that the external pressure $P_\infty(t)$ is a time-dependent function, whereas the CFD-calculated pressure $P_\infty(x)$ and velocity $P_\infty(u)$ are defined as functions of spatial coordinates. To integrate the CFD data into the unified bubble oscillation equation Eq. (1), it is imperative to convert these spatial distributions into time-dependent profiles along the bubble's path.

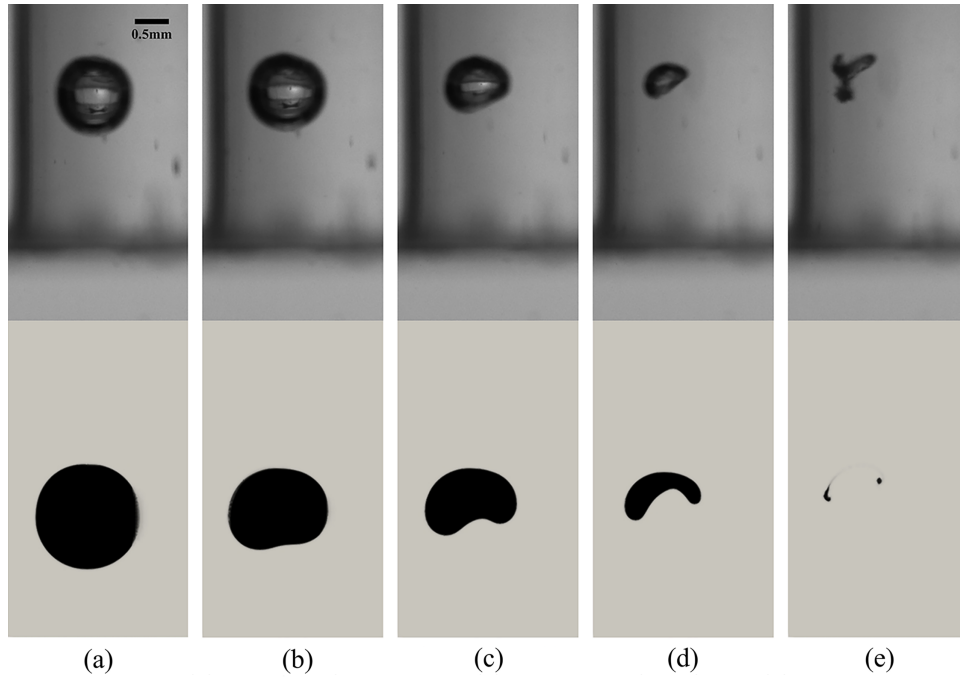


Figure 4: Validation of the numerical model against experimental data. The image sequence illustrates the evolution and collapse of a laser-induced bubble inside a circular tube. Top row: High-speed camera experimental recordings. Bottom row: Numerical phase-fraction contours obtained from the CFD simulation. The figures (a–e) correspond to instances at 80, 140, 160, 180, and 200 μs , respectively.

First, the velocity-related discrete points, $(p_i, x_i^{(p)})$ for $i = 0, \dots, N$, and the pressure-related discrete points, $(u_j, x_j^{(u)})$, are sorted in ascending order based on their spatial coordinates, such that $x_0^p < x_1^p < \dots < x_N^p$ and $x_0^u < x_1^u < \dots < x_N^u$. Subsequently, the intersection of these two coordinate sets is extracted to establish a common spatial domain, yielding (p_k, u_k, x_k) for $k = 0, \dots, M$. Concurrently, linear interpolation is applied to both the pressure and velocity data as follows:

$$p(x) = p_g + \frac{p_l - p_g}{x_l - x_g} (x - x_g), \quad x_g \leq x \leq x_l \quad (5)$$

$$u(x) = u_g + \frac{u_l - u_g}{x_l - x_g} (x - x_g), \quad x_g \leq x \leq x_l \quad (6)$$

To preclude the time step from approaching infinity as $u_k \rightarrow 0$ during the subsequent derivation of the time array, a minimum velocity threshold is imposed, defined as $u_{\min} \geq 10^{-6}$ m/s. Consequently, the corrected velocity is expressed as:

$$\tilde{u}_k = \begin{cases} \text{sign}(u_k) u_{\min}, & |u_k| < u_{\min} \\ u_k, & |u_k| \geq u_{\min} \end{cases} \quad (7)$$

In unidirectional flow, the transformation of a time-dependent function into a spatially dependent one can be achieved through the kinematic relationship between velocity and displacement.

$$t(x) = \int_{x_{\min}}^x \frac{d\zeta}{|u(\zeta)|} \quad (8)$$

In the discretized form, the temporal sequence can be evaluated recursively as follows:

$$t_0 = 0 \quad (9)$$

$$\Delta t_k = \frac{|x_k - x_{k-1}|}{|\vec{u}_k|} \quad k = 1, \dots, M \quad (10)$$

$$t_k = t_{k-1} + \Delta t \quad (11)$$

Consequently, a monotonically increasing time sequence, $0 = t_0 < t_1 < \dots < t_M = t_{\max}$, is obtained. This establishes a one-to-one correspondence between the temporal and spatial variables, $(t_k, x_k, p_k, \tilde{u}_k)$, thereby yielding the time-dependent evolution of pressure and velocity along the streamlines within the CFD flow field. To ensure continuous external excitation at any arbitrary integration time, piecewise linear interpolation is employed. Furthermore, the solution time is strictly confined to the interval $[0, t_{\max}]$ to guarantee that the bubble continuously experiences the external velocity and pressure driving forces throughout the computational process. The streamlines satisfy the following ordinary differential equation:

$$\frac{dx_i}{ds} = \frac{u(x_i)}{|u(x_i)|}, x_i(0) = x_{i,0} \quad (12)$$

To accurately capture the steep velocity and pressure gradients in regions such as the convergent section and the throat, this study employs an adaptive step-size Runge–Kutta integrator based on the cell relative velocity. The integration proceeds in the forward direction, with the primary parameters specified in [Table 2](#):

Table 2: Integration and step-size control parameters.

Parameter	Value	Description
$\Delta s_0/(\text{m})$	0.5	Initial step size
$S_{\max}/(\text{m})$	20	Maximum tracking length
$N_{\max}$	2000	Maximum number of integration steps
$\Delta S_{\min}/(\text{m})$	10^{-6}	Minimum step size
$\Delta S_{\max}/(\text{m})$	0.5	Maximum step size
$\mathcal{E}_{\max}$	10^{-6}	Truncation error tolerance

3 Results and Discussion

3.1 Characteristics of the Unidirectional Flow Field

As illustrated in [Fig. 5](#), the velocity and pressure contours reveal that the flow accelerates rapidly upon entering the convergent section from upstream. The axial velocity increases significantly within this converging region, reaching its peak near the minimum cross-section. Correspondingly, the static pressure drops sharply along the streamwise direction, forming a distinct localized low-pressure zone in the throat region.

From a radial perspective, the velocity profile at the throat cross-section is characterized by a distinct high-velocity core and a decelerated near-wall region. The shear layer expands radially and develops downstream, indicating that the momentum increment of the fluid varies across different radial positions, even at the same axial location. Similarly, the pressure field exhibits radial non-uniformity in the vicinity of the throat. The near-wall region experiences more pronounced pressure variations and is more susceptible to steep pressure drops, whereas the pressure changes within the high-velocity zone remain relatively gradual.

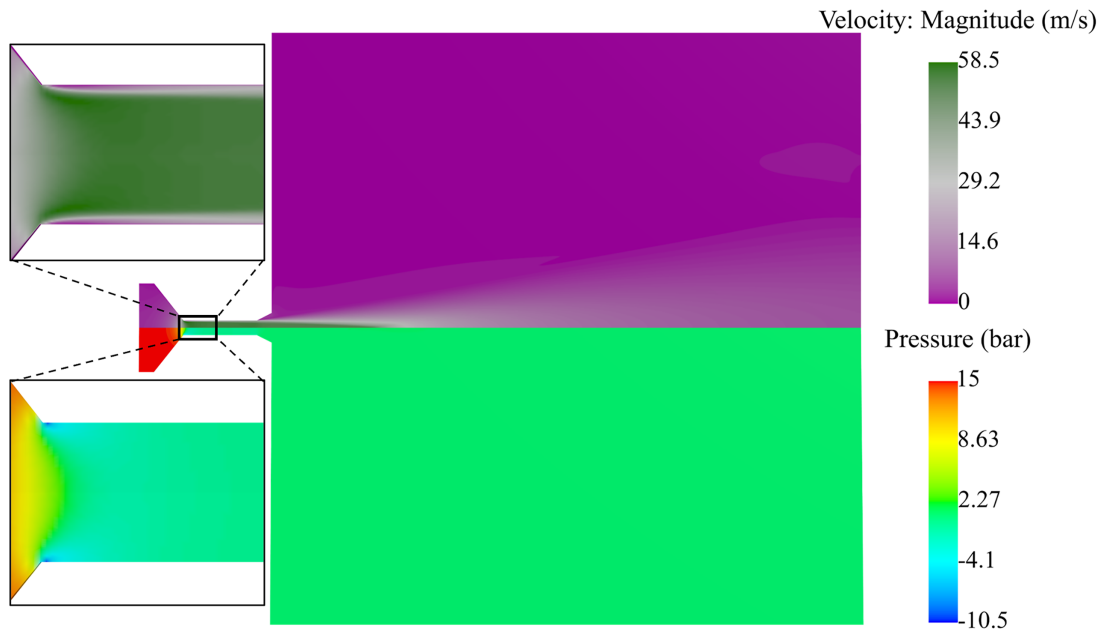


Figure 5: Velocity and pressure distributions from an Eulerian perspective.

To systematically evaluate the pressure drop and flow acceleration characteristics at various radial positions within the nozzle, five seed points are designated at distinct radial locations across the inlet cross-section. Their respective coordinates and regions are detailed in Table 3. As depicted in Fig. 6, the trajectories of the five streamlines clearly demonstrate how the fluid converges in the convergent section, traverses the throat, and enters the downstream straight section. Superimposed on the velocity magnitude contour, this provides a unified representation of the entire flow evolution—encompassing flow acceleration, pressure drop, and subsequent stabilization—experienced by the fluid at varying radial positions from upstream to downstream.

As depicted in Fig. 7, all five streamlines exhibit a sharp and distinct pressure drop within the throat region. Subsequently, the pressure gradually recovers, entering a relatively stable plateau until it reaches the downstream ambient pressure of the nozzle. However, the magnitude of this pressure drop varies significantly among the streamlines. The red streamline experiences the most drastic pressure reduction, followed by a sequential decrease in the pressure differentials for the blue, green, and yellow streamlines, while the black streamline exhibits the minimum pressure drop. These results indicate that the combined effects of local acceleration, streamline curvature, and near-wall shear within the convergent section lead to substantial variations in pressure evolution across different radial positions. Crucially, the emergence of these extreme negative pressure values at the throat represents a physical state of intense liquid tension. Driven by the severe local flow acceleration through the constriction, coupled with intense turbulent fluctuations and the slip velocity between the surrounding liquid and the bubbles, the instantaneous local pressure p_{∞} drops drastically. When p_{∞} falls not only below the thermodynamic saturated vapor pressure but deeply into the negative regime, the liquid enters a metastable state. This spatial domain, functionally defined as a volume “under tension” [37], constitutes the primary driving mechanism for phase change. Consequently, larger initial radius bubbles traversing the nozzle will inevitably be subjected to these extreme tensile stresses and pressure fluctuations, thereby triggering explosive cavitation inception and subsequent bubble collapse during the mid-to-downstream secondary pressure recovery.

Table 3: The respective coordinates of five probe points.

Colors	Coordinates [m]	Regions
Red	[0, 0.0025, 0]	The region close to the wall
Blue	[0, 0.002, 0]	The outer region of the shear layer
Green	[0, 0.0015, 0]	The central region of the shear layer
Yellow	[0, 0.001, 0]	The inner region of the shear layer
Black	[0, 0.0005, 0]	The high-speed region near the axis

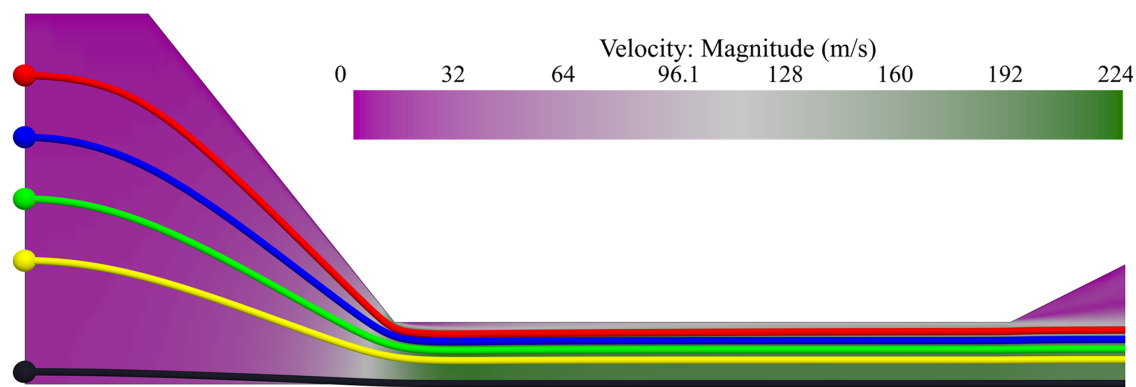


Figure 6: Schematic of the seed point locations and the resulting five streamlines and velocity profile at the throat.

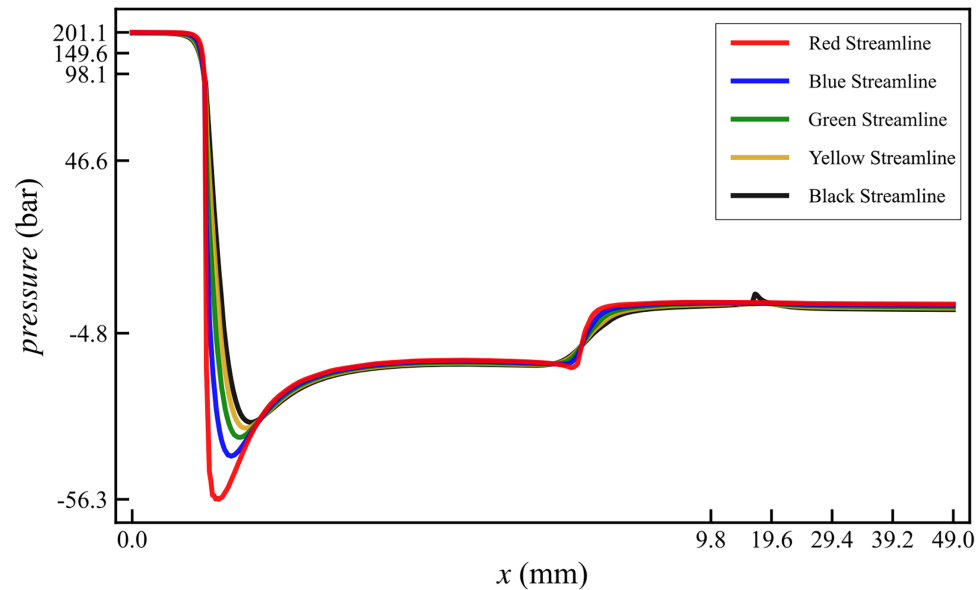


Figure 7: Spatial distribution of pressure along each streamline.

Fig. 8 presents the streamwise velocity distributions of the five streamlines. All streamlines undergo acceleration within the convergent section, with their velocity peaks concentrated at the nozzle throat. Upon entering the downstream straight section, the velocities gradually decrease and tend to converge. However, distinct differences in velocity evolution are observed among the varying radial streamlines at

the throat and its subsequent downstream region. After reaching its peak, the black streamline sustains a relatively long high-speed interval. This high-speed zone sequentially shortens for the yellow, green, and blue streamlines, while the red streamline experiences an earlier and more pronounced velocity decay. This trend perfectly aligns with their radial positions: momentum is more concentrated in the near-axis region, leading to a stronger capacity for velocity retention along the flow path. Conversely, streamlines closer to the wall are more susceptible to viscous constraints and shear dissipation, thereby exhibiting a shorter high-speed maintenance distance.

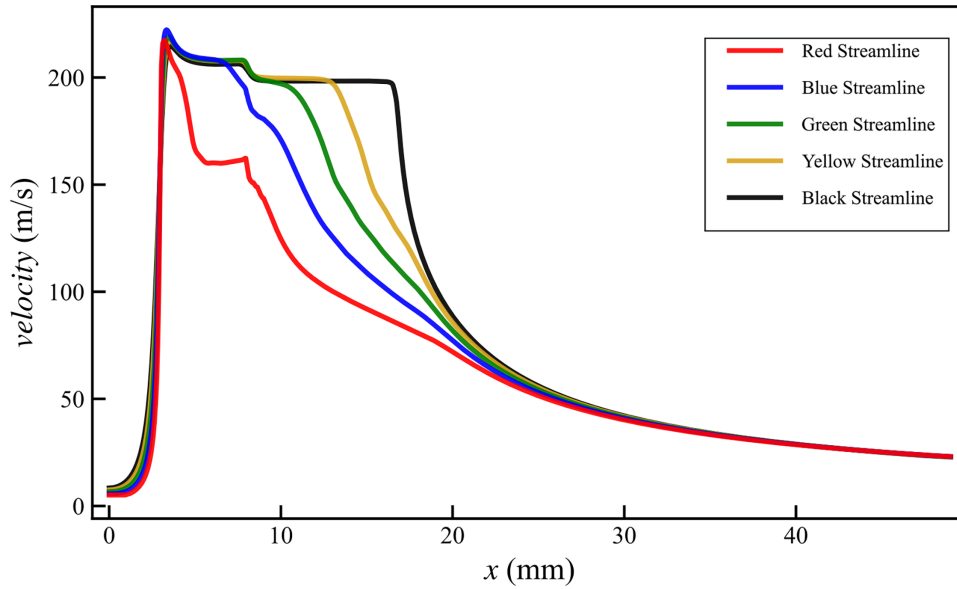


Figure 8: Spatial distribution of velocity along each streamline.

In summary, the coupled pressure-velocity analysis of the five representative streamlines in this section demonstrates that the internal single-phase flow exhibits a strong pressure drop-acceleration coupling within the throat region. Furthermore, distinct differences are observed among the varying radial streamlines concerning the magnitude of the pressure drop, peak velocity, and the extent of the high-speed retention zone.

3.2 Effect of the Initial Radial Position

To investigate the effect of the initial radial position on bubble evolution, Fig. 9a presents the evolution of the dimensionless bubble radius R/R_0 along the axial coordinate x under various initial radial locations. It can be observed that, due to the relatively high ambient pressure, the bubbles in all five cases do not expand upon entering the convergent section of the nozzle. Subsequently, as they approach the throat, the microbubbles begin to expand once the local pressure drops below the vapor pressure. Upon exiting the nozzle, although the ambient pressure has significantly exceeded the saturated vapor pressure, the bubbles do not cease growing. This is attributed to the inertial effects governed by the second derivative term in the unified bubble oscillation equation. The bubbles continue to grow, reaching their maximum radii at approximately $x = 30\sim 40$ mm from the inlet, followed by violent collapses and multiple damped oscillations in the downstream region. As r_0 increases from 0.1 to 2.5 mm, the maximum dimensionless radius $(R/R_0)_{\max}$ increases monotonically from approximately 6.0×10^3 to nearly 1.0×10^4 . Concurrently, the axial position of the primary collapse shifts significantly downstream, and the spatial lengths of the subsequent secondary

and tertiary rebounds increase accordingly. This indicates that bubbles closer to the wall remain in the depressurization zone for a longer effective duration during convective transport, thereby accumulating greater inertial energy, which ultimately results in a more intense collapse downstream. Fig. 9b depicts the converging-diverging nozzle geometry. By sharing the same axial coordinate as Fig. 9a, it directly correlates the spatial evolution of the bubble with its instantaneous physical location within the nozzle. This aligned configuration is consistently applied to all subsequent relevant figures.

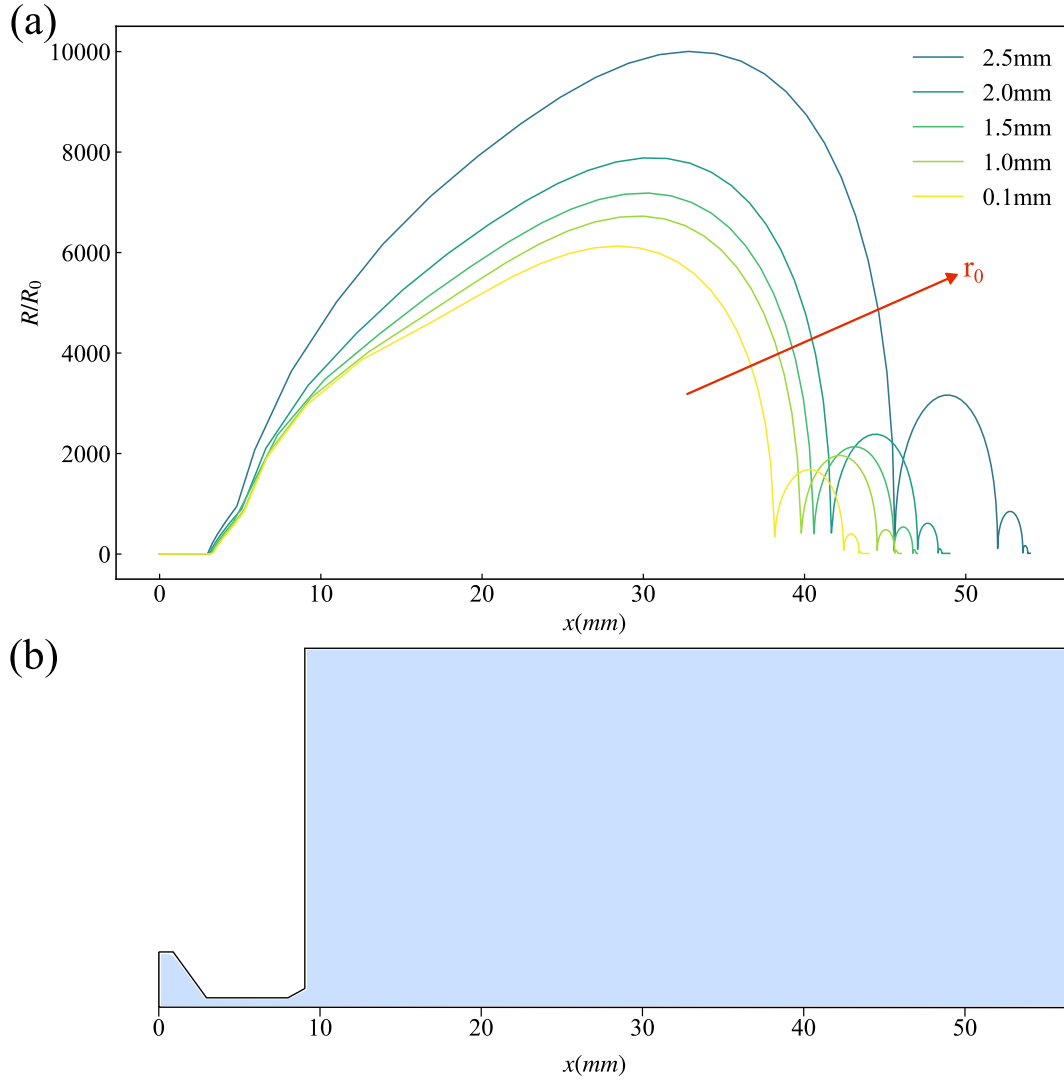


Figure 9: Spatial evolution of bubble dynamics in the microchannel. (a) Evolution of the dimensionless bubble radius along the flow direction for different initial radial positions (r_0). (b) Schematic of the converging-diverging nozzle geometry.

As illustrated in Fig. 10, bubbles from different initial radial positions all generate pressure peaks on the order of 1.4×10^4 in the vicinity of the primary collapse, followed by multiple peaks associated with subsequent re-expansions and re-collapses. A comparison with Fig. 9a reveals that the locations of the maximum pressure peaks generally coincide with the positions where the R/R_0 curve exhibits a steep decline. By comparing the results across different radial positions r_0 , it is evident that the pressure peaks for near-axis bubbles are concentrated relatively close to the nozzle exit, with a more rapid attenuation during subsequent

rebounds. In contrast, the pressure peaks for near-wall bubbles are distributed further downstream, and the amplitudes between successive peaks show a slight increase. This phenomenon indicates that the initial radial position not only affects the magnitude of bubble expansion but also alters the spatial distribution of the collapse impacts, which is of critical importance for assessing the cavitation erosion risk across different regions of the structural walls.

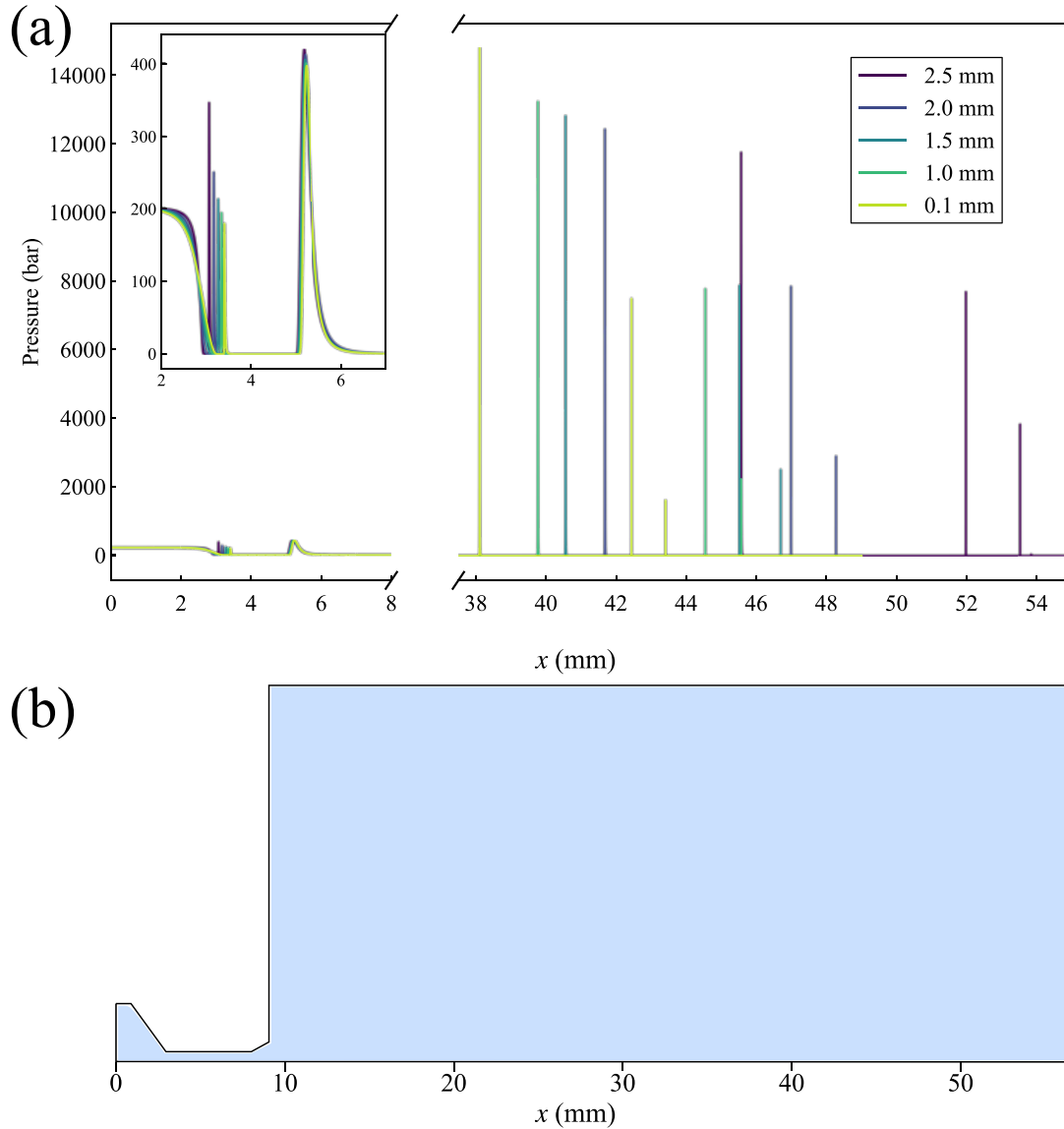


Figure 10: (a) Spatial evolution of the bubble pressure under different initial radial positions. (b) Schematic of the converging-diverging nozzle geometry.

Analysis in the time domain further reveals the underlying mechanisms of the aforementioned differences. Fig. 11 presents the temporal evolution of R_0 across the five radial positions. After experiencing a substantial negative pressure, the bubble begins to expand, collapsing rapidly upon reaching $R_{\max}$ and generating subsequent oscillations. As r_0 increases, the onset of the expansion phase is slightly delayed, but its duration is significantly prolonged; consequently, the first collapse time T_c increases from approximately 0.5 ms to nearly 0.9 ms. Meanwhile, near-wall bubbles retain stronger inertia after the primary collapse,

resulting in noticeably higher amplitudes during the second and third pulsation cycles compared to near-axis bubbles. Combined with the flow field characteristics, it can be inferred that the flow velocity near the wall is slightly lower than the axial mainstream. This allows the bubbles to spend a longer acceleration time within the depressurization region, which ultimately contributes to the near-wall bubbles generating a larger $R_{\max}$ and longer pulsation cycles.

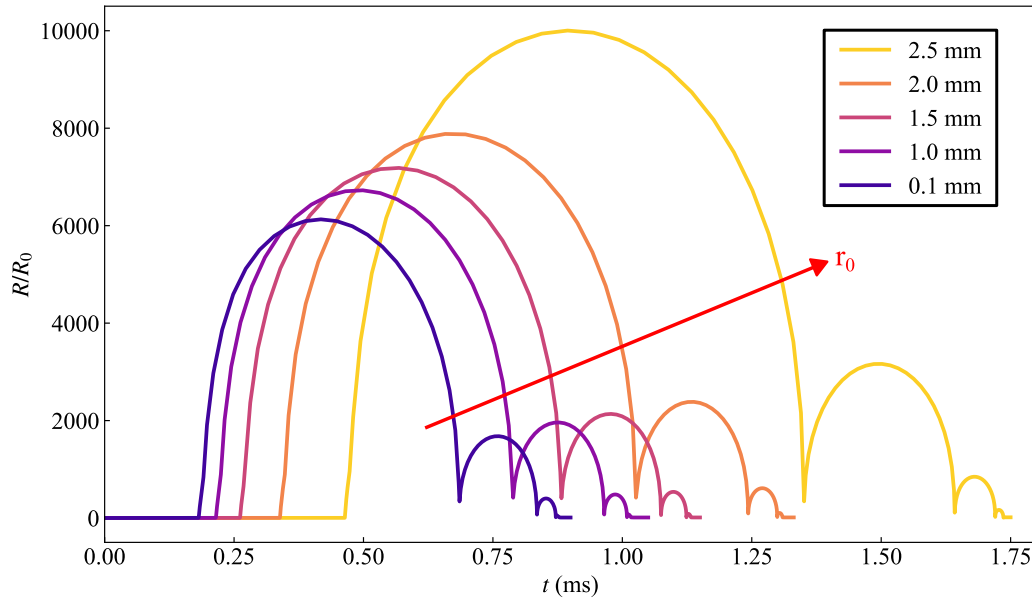


Figure 11: Temporal evolution of the bubble radius under different initial radial positions.

Fig. 12 illustrates the temporal evolution of the impact pressure generated by bubble collapse at different radial positions. Across all cases, the timing of the pressure peak coincides with the minimum bubble radius; that is, the maximum impact pressure is emitted when the bubble contracts to its minimum volume. For near-axis bubbles, the pressure peak occurs at an earlier time, whereas for near-wall bubbles, it appears later. By correlating Figs. 11 and 12, it is evident that altering the initial radial position essentially changes the duration of the negative pressure exposure. This, in turn, influences the bubble's pulsation cycle and the resulting collapse pressure, ultimately dictating the magnitude and timing of the impact peak.

As illustrated in the radial position analysis, the dynamic evolution of bubbles exhibits profound spatial heterogeneity and non-linear responses. Physically, the near-wall bubble (red line) attains the maximum non-dimensional radius due to a synergistic effect of boundary layer kinematics and local pressure gradients. The lower convective velocity near the wall extends the bubble's residence time within the extreme negative pressure zone at the throat, allowing it to absorb more mechanical energy from the flow field to drive explosive growth.

Furthermore, a significant spatial delay in the collapse point is observed for the near-wall bubble. According to classical bubble dynamics, the collapse duration is proportional to the maximum radius. The substantially larger $R_{\max}$ of the near-wall bubble dictates a prolonged collapse phase, allowing it to be convected further downstream despite its lower axial velocity. During this violent collapse, the immense centripetal inertia of the surrounding liquid extremely compresses the internal non-condensable gas, instantaneously converting liquid kinetic energy into thermodynamic internal energy, which manifests as the extreme pressure spikes exceeding 2000 bar.

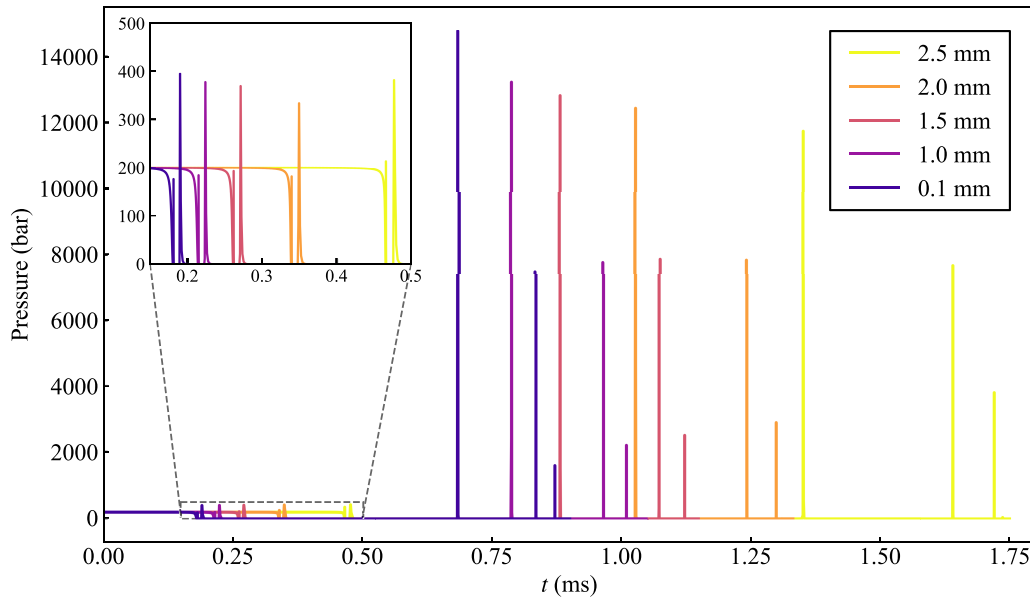


Figure 12: Temporal evolution of the bubble pressure under different initial radial positions.

Most importantly, the subsequent rebound exhibits a severely non-linear attenuation. This abrupt departure from conventional elastic rebound is mechanistically driven by intense energy dissipation. During the terminal phase of the primary collapse, the extreme interfacial acceleration induces significant acoustic radiation into the surrounding compressible liquid, coupled with local viscous dissipation. This irreversible loss of system mechanical energy deprives the bubble of the driving force required for subsequent expansions, yielding the observed highly damped, non-linear oscillatory behavior.

The complete life cycle of cavitation bubbles within this convergent-divergent micro-channel can be comprehensively understood by analyzing the spatial and temporal evolution of the non-dimensional radius alongside the pressure profiles as depicted in Figs. 8–12.

Initially, as the bubbles enter the convergent section, their growth is strictly suppressed by the relatively high local static pressure. However, as they approach the narrow throat, the severe local fluid acceleration causes the pressure to plummet deeply below the saturated vapor pressure, forming the volume of ‘under tension’. This extreme depressurization instantly triggers the onset of explosive bubble growth.

Intriguingly, as the bubbles are convected towards the mid-to-downstream region, the local ambient pressure begins to recover significantly above the vapor pressure. Despite this pressure recovery, the bubbles do not immediately cease expanding. This pronounced delay is primarily attributed to the powerful outward liquid inertia effect, which is rigorously captured by the high-order derivative terms in the unified theoretical model employed in this study. Driven by this immense inertia, the bubble continues to expand against the recovering adverse pressure gradient, ultimately reaching its $R_{\max}$ at a significant distance downstream before initiating the violent collapse phase.

Synthesizing the results across multiple cases reveals that as the initial bubble is positioned closer to the wall, both the onset of expansion and the time of the primary collapse are globally delayed. Furthermore, the maximum radius increases significantly, and the post-collapse rebounds occur more frequently. To avoid redundant presentations and to delve into the critical mechanistic question—namely, how the radial coordinate alters the entire history of the bubble from expansion to collapse, as well as the mechanism of impact load generation—this section selects two highly representative cases for a comparative analysis: the

near-wall case ($r_0 = 2.5$ mm) and the near-axis case ($r_0 = 0.1$ mm). Fig. 13 illustrates the coupled $R(x) - P(x)$ evolution along the trajectory, while Fig. 14 presents the coupled $R(t) - P(t)$ evolution. Both figures utilize superimposed layouts to compare the radius variation curves and pulsation intensities of the near-wall and near-axis bubbles within the same framework, incorporating magnified insets to clearly resolve the local features as the bubbles traverse the intense negative pressure zone.

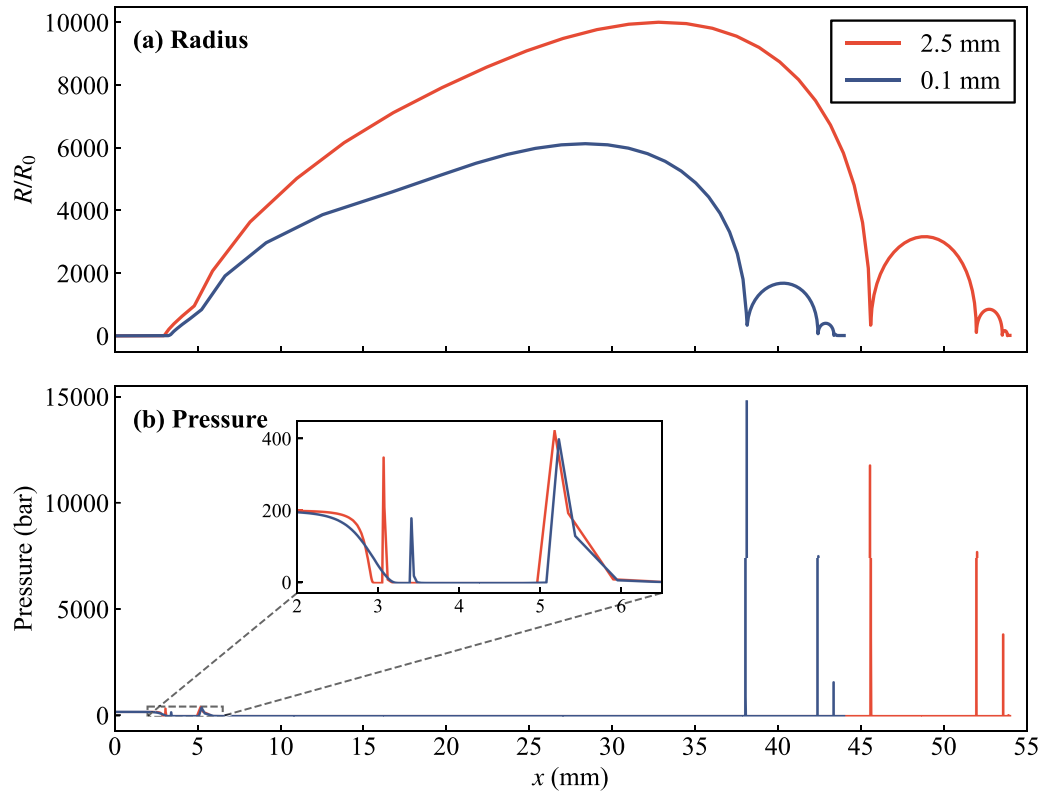


Figure 13: Comparison of the coupled axial radius-pressure evolution between near-wall and near-axis bubbles. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

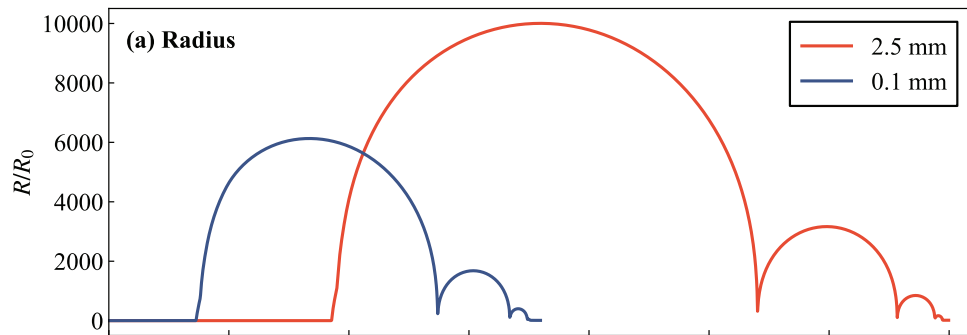


Figure 14: (Continued)

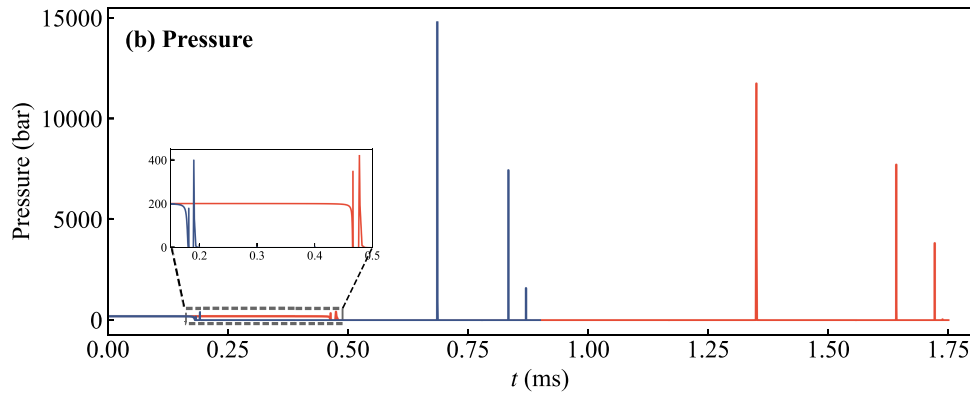


Figure 14: Comparison of the coupled temporal radius-pressure evolution between near-wall and near-axis bubbles. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

The triggering phase is characterized by a structural feature of spatial consistency coupled with temporal separation. As evident from the magnified inlet inset in Fig. 13, both cases undergo a severe negative pressure within a short distance after entering the nozzle. Subsequently, the bubble radius grows continuously, with the triggering location concentrated around $x \approx 3$ mm. This spatial location for the inception of bubble growth is predominantly governed by the nozzle-induced pressure drop, rendering the influence of r_0 negligible. In contrast, the magnified inset in Fig. 14 reveals a significant discrepancy in the expansion onset time: the near-axis bubble initiates evident growth at approximately $t \approx 0.18$ ms, whereas the near-wall bubble only begins rapid expansion after $t \approx 0.24$ ms. This temporal lag indicates that altering r_0 changes the time required for the bubble to reach the low-pressure zone, thereby significantly shifting the expansion onset time while leaving the spatial onset location virtually unchanged. Furthermore, short-term pressure fluctuations can be observed in the inset of Fig. 14, yet the bubble radius curve exhibits no discernible decline. This suggests that the early growth of the bubble radius, once initiated, is primarily driven by the substantial negative pressure, and weak perturbations are insufficient to reverse the expansion trend.

Secondly, the scale discrepancy during the expansion phase directly determines the spatiotemporal migration of the primary collapse. Fig. 13 demonstrates that the maximum dimensionless radius, $(R/R_0)_{\max}$, for the near-wall case is significantly higher than that of the near-axis case, reaching approximately $(R/R_0)_{\max} \approx 1.0 \times 10^4$ for the former and $(R/R_0)_{\max} \approx 6.0 \times 10^3$ for the latter. Meanwhile, the near-wall bubble sustains its expansion over a longer trajectory scale before completing its first collapse. This phenomenon is equally evident in the time domain: Fig. 14 shows that the near-wall bubble experiences a prolonged expansion duration, significantly delaying the time at which it reaches $R_{\max}$. Combining this with the general trends from the previous section, it can be concluded that the radial position alters the profile of the pressure history experienced by the bubble. This, in turn, modifies the effective exposure time to the negative pressure environment, thereby collectively amplifying the expansion process across both spatial and temporal dimensions.

The collapse and rebound phases manifest the magnitude of the pressure peak and the timing of the pressure spike, which are critical factors for assessing damage risk. Both figures demonstrate that at the instant of the primary violent collapse, the bubble radius sharply drops to a minimum, accompanied by a narrow-pulse-width high-pressure spike on the order of 1.5×10^4 bar and 1.2×10^4 bar. Furthermore, the occurrence time of this high-pressure spike strictly aligns with the minimum of R . This correlation confirms the origin of the load: the impact load is dominated by the inertia induced by the pressure drop and is generated by the bubble collapse, rather than being directly caused by low-frequency pressure fluctuations.

Comparing the two cases, the primary collapse of the near-axis bubble occurs earlier and exhibits a higher maximum pressure peak. Although the first collapse of the near-wall bubble is slightly delayed, it presents a larger subsequent rebound radius and a more pronounced secondary collapse, thereby generating more rebound pressure spikes. Consequently, it is evident that a single peak metric, $p_{\max}$, is insufficient to fully characterize the influence of r_0 . The impulsive load generated at the near-axis location highlights the risk of a single severe peak, whereas the near-wall location emphasizes the cumulative risk associated with the frequency and duration of impact events. This conclusion aligns consistently with the trends observed across the multiple cases discussed previously.

3.3 Effect of the Initial Ambient Pressure

To evaluate the coupling effects of the single-bubble nonlinear response, this section keeps the numerical model and the initial bubble radius ($R_0 = 0.5 \mu\text{m}$) constant, while systematically setting the initial ambient pressure (p_0) to 25, 50, 200, and 400 bar. The external pressure along the same streamline is utilized to drive the bubble dynamics equation. As observed in Fig. 15, across all cases, minor pressure perturbations initially emerge in the upstream region of the flow field. The pressure peaks are predominantly concentrated in the middle-to-downstream sections of the channel, manifesting as multiple pulses. These pulses correspond to the primary bubble collapse and the subsequent re-collapse events during rebounds, with the maximum peak pressure reaching the order of 1.6×10^4 bar.

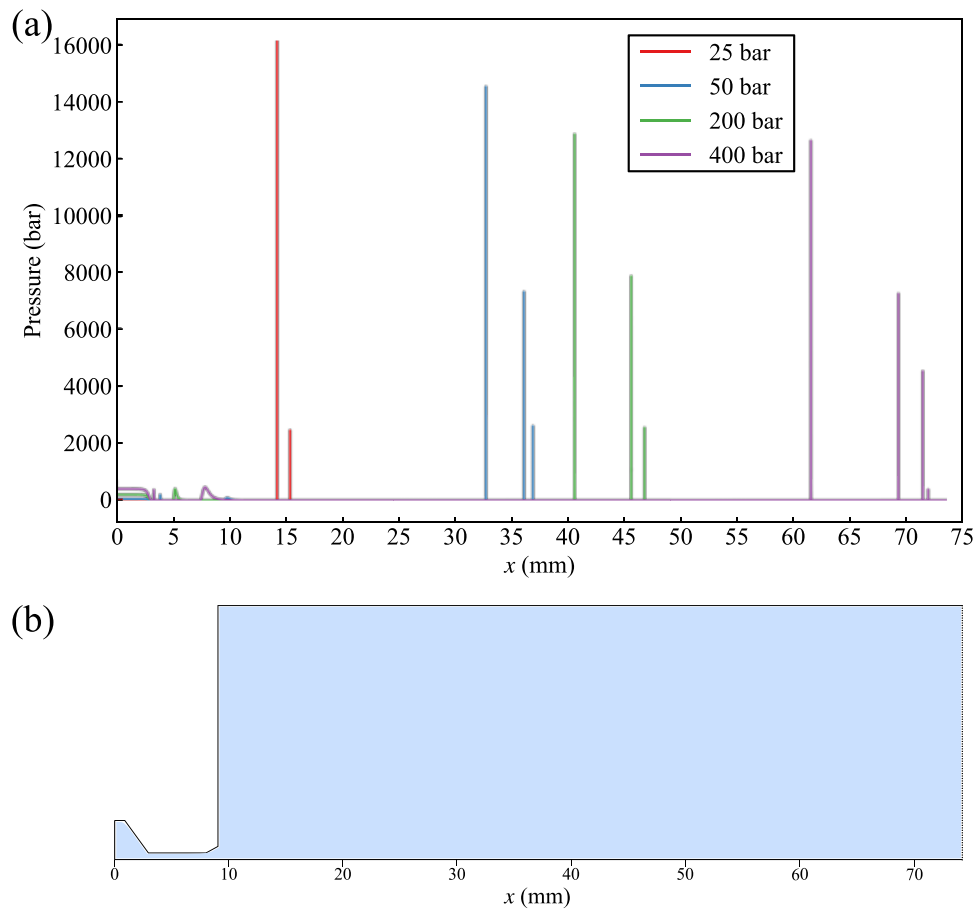


Figure 15: (a) Spatial evolution of the collapse pressure under different initial pressures. (b) Schematic of the converging-diverging nozzle geometry.

The streamwise distribution reveals that the location of the primary collapse exhibits an overall downstream migration trend as the initial pressure varies. At $p_0 = 25$ bar, the maximum pressure spike occurs near $x \approx 14$ mm. As p_0 increases to 50 bar, the primary collapse location shifts to approximately $x \approx 32$ mm. Under the conditions of $p_0 = 200$ bar and $p_0 = 400$ bar, the maximum pressure spikes shift further downstream to roughly $x \approx 41$ mm and $x \approx 62$ mm, respectively, accompanied by distinct secondary collapse pulses emerging closer to the nozzle exit. This spatial migration does not imply that the collapse occurs later in time; rather, it reflects the intensified convective transport under a higher P_0 , which sweeps the bubble further downstream before it reaches the necessary collapse conditions. This result emphasizes that the maximum impact intensity decreases monotonically with an increase in the initial pressure. Furthermore, secondary collapse pressure spikes are observed across multiple cases in the figure, located at a limited distance from the primary peak. This indicates that the multi-pulse loads induced by the rebound and re-collapse cycles must be incorporated into subsequent statistical analyses.

Fig. 16 presents the temporal evolution of wall pressure under varying initial ambient pressures p_0 . As p_0 increases, the first peak of the pressure pulse exhibits a distinct temporal delay. Specifically, the primary collapse for the $p_0 = 25$ bar case occurs at $t \approx 0.88$ ms, whereas this event is delayed to $t \approx 1.40$ ms when p_0 reaches 400 bar. This phase shift reflects the nonlinear driving effect of an increased ambient pressure on the bubble's expansion-contraction cycle. Regarding peak intensity, the maximum pressure amplitude does not increase monotonically with p_0 ; instead, it displays an inverse correlation. The highest transient impact pressure emerges at $p_0 = 25$ bar, peaking at roughly 1.6×10^4 bar. As p_0 rises, this peak pressure gradually attenuates, dropping to about 1.2×10^4 bar at $p_0 = 400$ bar. This suggests that a lower initial ambient pressure induces more violent interface motion and a highly concentrated energy focusing. Furthermore, the discrete distribution of subsequent pulses under high-pressure conditions indicates that p_0 alters the external pressure gradient. This exerts a complex regulation on the secondary bubble dynamics, causing the energy release to be more temporally dispersed.

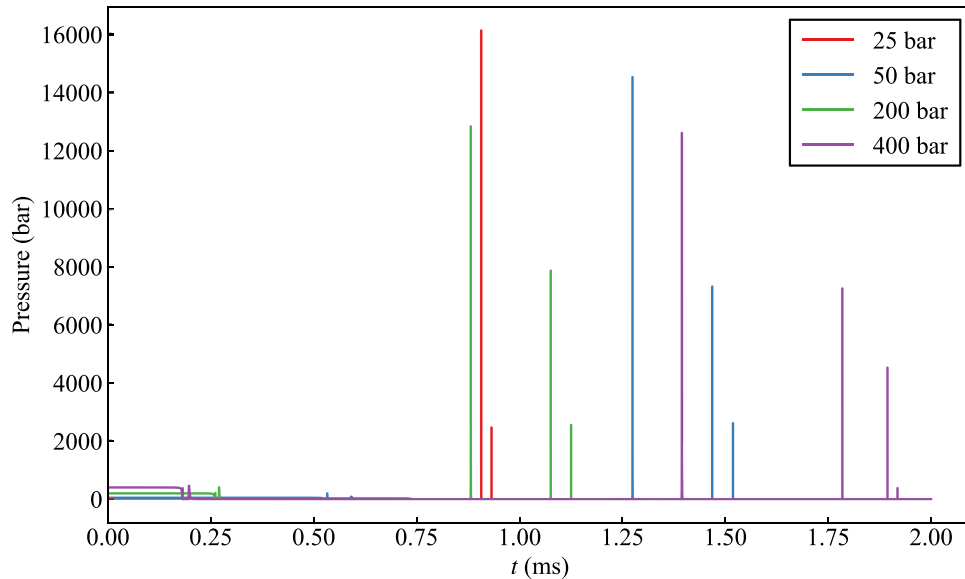


Figure 16: Temporal evolution of the collapse pressure under different initial pressures.

Fig. 17 illustrates the spatial evolution of the dimensionless bubble radius under four initial pressures p_0 . In all cases, the bubbles undergo an initial expansion followed by a collapse, generating one or two subsequent

rebounds with decaying amplitudes after the primary collapse. Although the curves for different p_0 exhibit consistent deviations throughout the evolution, this variation is not strictly monotonic.

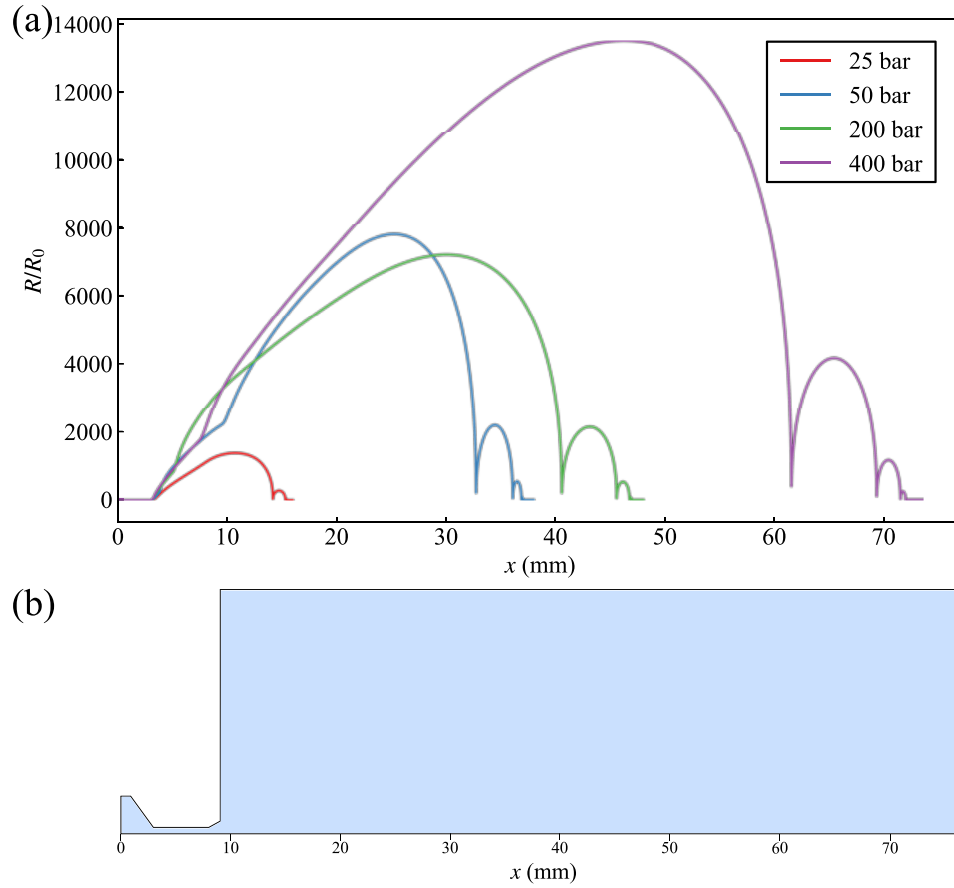


Figure 17: (a) Spatial evolution of bubble radius under different initial pressures. (b) Schematic of the converging-diverging nozzle geometry.

Three key observations emerge. First, the maximum dimensionless radius, $(R/R_0)_{\max}$, varies significantly with P_0 , yielding values of approximately 1.6×10^3 , 7.8×10^3 , 7.0×10^3 , and 1.3×10^4 for $P_0 = 25, 50, 200$, and 400 bar, respectively. Second, the location of the primary collapse shifts progressively downstream as P_0 increases, occurring near $x \approx 14, 32, 41$, and 62 mm for the respective P_0 cases. Third, following the primary collapse, most cases exhibit one or two downstream rebound oscillations with diminishing amplitudes, which are spatially distinct from the primary collapse site. These results demonstrate that by modulating the depressurization window and the pressure recovery process, the initial pressure alters the bubble expansion distance within the channel, thereby directly determining the spatial location of the primary collapse and the magnitude of the maximum dimensionless radius.

Fig. 18 further elucidates the effect of the initial pressure on the temporal evolution of the bubble radius. A striking monotonic feature is the onset of expansion: as p_0 increases, the bubble initiates growth progressively earlier. However, the subsequent expansion dynamics—specifically the dimensionless maximum radius $(R/R_0)_{\max}$ and the primary collapse time t_c —demonstrate highly non-monotonic trends. For example, the 50 bar case produces a larger $(R/R_0)_{\max}$ and a significantly longer expansion duration than the 200 bar case. Consequently, their primary collapse times do not scale linearly with p_0 . Both the 25 and

200 bar cases collapse near $t_c \approx 0.9$ ms, whereas the 400 and 50 bar cases are substantially delayed to roughly 1.4 and 1.25 ms, respectively. This lack of a monotonic correspondence between $(R/R_0)_{\max}$ and t_c indicates a complex coupling mechanism between the flow velocity and the local pressure drop.

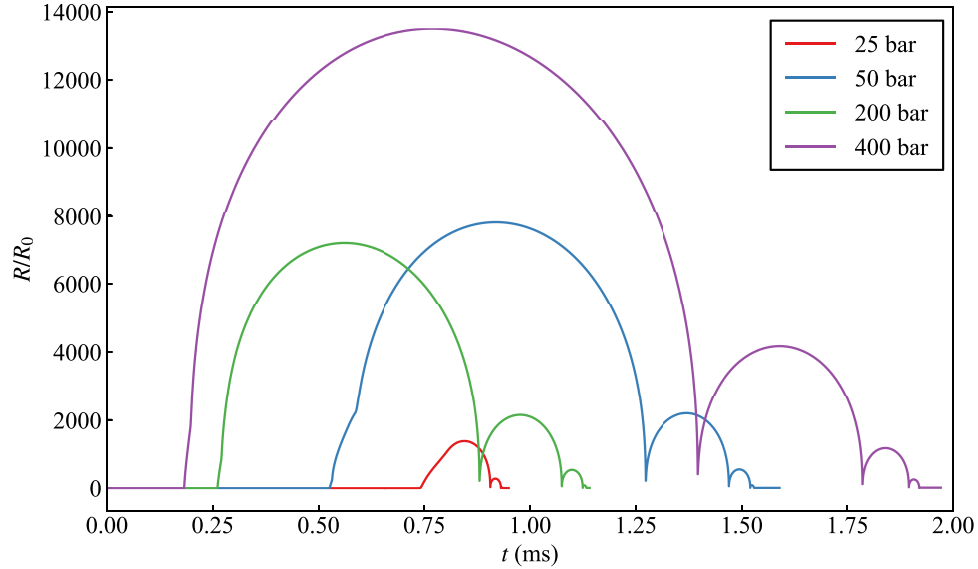


Figure 18: Temporal evolution of bubble radius under different initial pressures.

The influence of varying driving pressures on the bubble dynamics reveals highly non-linear and counter-intuitive physical phenomena, as comprehensively demonstrated in the temporal and spatial evolutions of the non-dimensional radius and internal pressure.

Primarily, a higher driving pressure significantly elevates the convective velocity of the bulk flow. Consequently, larger initial radius bubbles are accelerated into the micro-channel throat much earlier ($t \approx 0.18$ ms). There, the intensified Bernoulli effect produces extreme local liquid tension and gives rise to a characteristic volume “under tension”. This severe local depressurization acts as a massive energy source, driving the bubble to a colossal maximum dimensionless radius ($(R/R_0)_{\max} > 13000$). Governed by Rayleigh collapse kinetics, this significantly enlarged $R_{\max}$ prolongs the collapse duration. Coupled with the extreme axial velocity, the massive bubble is convected far downstream ($x > 60$ mm) before the ambient pressure recovery can induce the final collapse.

This phenomenon reveals that the triggering conditions for the primary collapse are dictated by the entire external pressure history, and cannot be explained solely by a single average pressure or maximum expansion scale. A critical conclusion drawn from this is that an increase in $(R/R_0)_{\max}$ does not necessarily correspond to a higher $p_{\max}$. For instance, The $p_0 = 25$ and 50 bar scenarios produce more extreme pressure peaks despite their smaller dimensionless radii. This demonstrates that the peak pressure is predominantly governed by the combined effect of the transient contraction intensity during the collapse phase and the external pressure recovery process, rather than by the bubble expansion scale alone.

As compared in Fig. 19, the expansion phase for both cases is triggered by the pressure drop near the inlet; however, significant discrepancies exist in the bubble growth scale and evolution distance. For the $p_0 = 400$ bar case, the radius continues to grow over an extended axial range, reaching a peak of $(R/R_0)_{\max} \approx 1.3 \times 10^4$, followed by a rebound and a distinct secondary collapse downstream. Conversely, the bubble growth in the $p_0 = 25$ bar case terminates much earlier, yielding a smaller peak of $(R/R_0)_{\max} \approx 1.6 \times 10^3$ and a

relatively limited rebound amplitude. Recalling the results from Figs. 15 and 17, the primary collapse locations for the two cases are situated at $x \approx 14$ mm (25 bar) and $x \approx 62$ mm (400 bar), respectively. By mapping the radius growth profiles and the corresponding pressure variations onto the same reference frame, Fig. 19 clearly illustrates the key differences: the high initial pressure condition features a prolonged expansion history and stronger rebounds, whereas the low initial pressure condition undergoes a smaller expansion scale with a more localized collapse position.

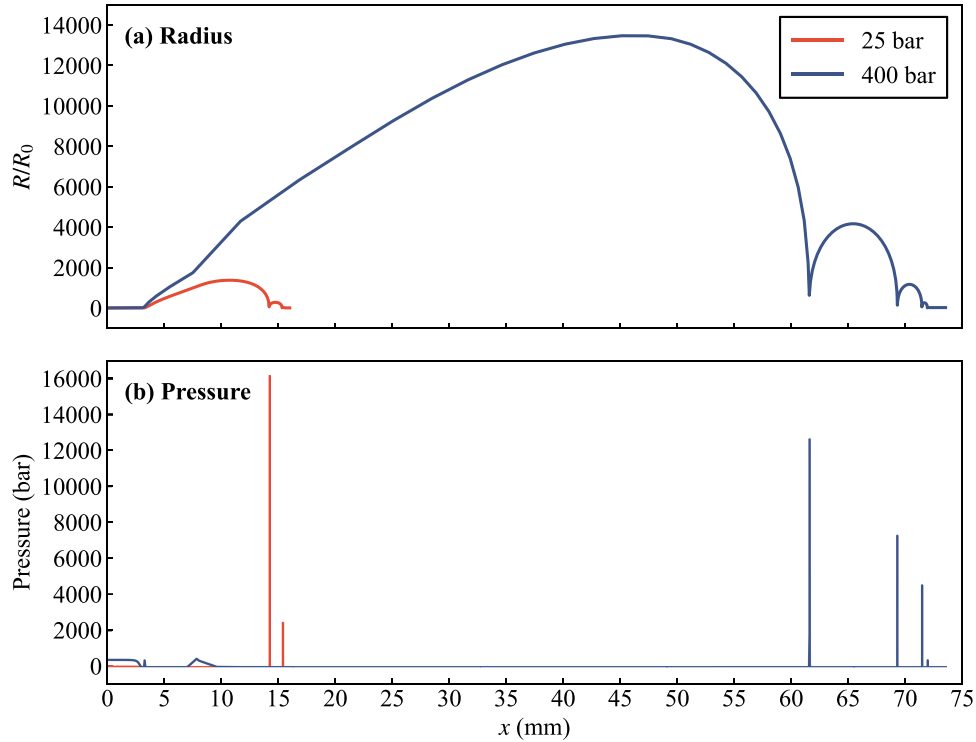


Figure 19: Comparison of the coupled spatial radius-pressure distribution between $P_0 = 25$ bar and $P_0 = 400$ bar cases. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

Further comparison in Fig. 20 highlights significant differences between the two cases regarding the onset of bubble expansion and the first oscillation cycle. The $P_0 = 400$ bar case initiates growth earlier and sustains expansion longer, with its first oscillation and subsequent rebounds spanning a broader time window. In contrast, the $P_0 = 25$ bar case exhibits a delayed growth onset, a shorter first oscillation period, and fewer rebound oscillations with attenuated amplitudes. These observations align with Figs. 16 and 18, demonstrating that the initial pressure alters the bubble evolution dynamics by modulating the external pressure history and the convective timescale.

Most intriguingly, the collapse intensity exhibits a non-linear inversion: the lowest driving pressure (25 bar), despite producing the smallest maximum radius, generates the most extreme primary collapse pressure spike exceeding 16,000 bar. This phenomenon highlights a fundamental competition between energy focusing and dissipation. For the 25 bar case, the compact collapse process facilitates a highly localized and efficient adiabatic compression of the internal gas, driven by the sharp centripetal inertia of the liquid. Conversely, the exceptionally voluminous bubble in the 400 bar case undergoes a prolonged and spatially extended contraction. During this protracted phase, substantial mechanical energy is irreversibly lost to the surrounding compressible liquid via enhanced acoustic radiation, interfacial instabilities, and turbulent

viscous shearing. This severe energy dissipation attenuates the final energy focusing efficiency, thereby reducing the peak thermodynamic pressure despite the higher initial system energy.

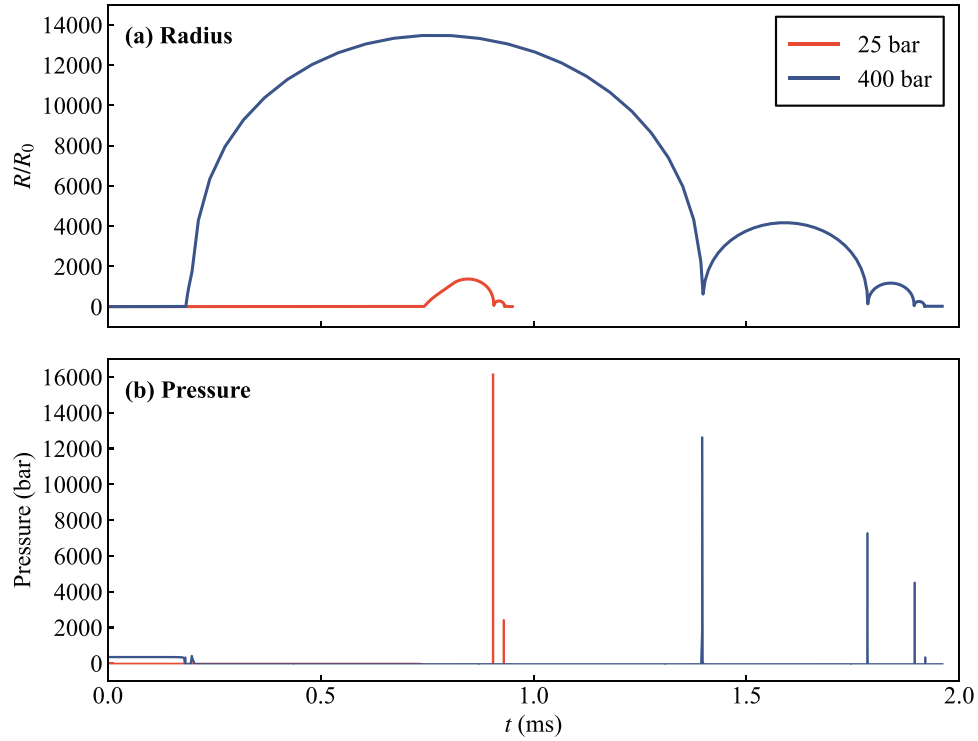


Figure 20: Comparison of the coupled temporal radius-pressure distribution between $P_0 = 25$ bar and $P_0 = 400$ bar cases. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

3.4 Effect of the Initial Bubble Radius

To elucidate scale effects, a bubble dynamics model is employed to simulate single-bubble responses under an identical external pressure history $P_\infty(t)$ for initial radii of $R_0 = 0.5, 2, 5$, and $8 \mu\text{m}$. Subsequently, the corresponding transient impact pressures are extracted from the near-field pressure.

Fig. 21 illustrates the axial spatial evolution of bubbles with different initial radii under identical flow conditions. During the initial expansion phase, the trajectories for various R_0 strongly coincide, diverging significantly only after traversing the negative-pressure zone. Both the primary collapse location and the maximum radius exhibit a strong dependence on R_0 . Specifically, R_{max} decreases as R_0 increases, dropping from approximately $6.8 \times 10^3 \mu\text{m}$ at $R_0 = 0.5 \mu\text{m}$ to merely $2.2 \times 10^3 \mu\text{m}$ at $R_0 = 8 \mu\text{m}$. Concurrently, the primary collapse shifts further downstream for smaller bubbles: the implosions for $R_0 = 8$ and $5 \mu\text{m}$ conclude near $x \approx 29$ and 41 mm, respectively, whereas for 2 and $0.5 \mu\text{m}$, they are delayed to $x \approx 42$ and 49 mm. This indicates that smaller initial bubbles experience a prolonged effective expansion distance, pushing their primary collapse toward the downstream channel or even the nozzle exit, while larger bubbles tend to implode midstream. Furthermore, the spatial distribution of the secondary oscillations confirms that smaller bubbles localize both their primary and secondary collapses near the exit region, thereby posing a significantly greater threat of cavitation damage to the downstream wall structures. Conversely, the collapse of bubbles with larger initial radii inflicts more severe damage on the midstream wall structures.

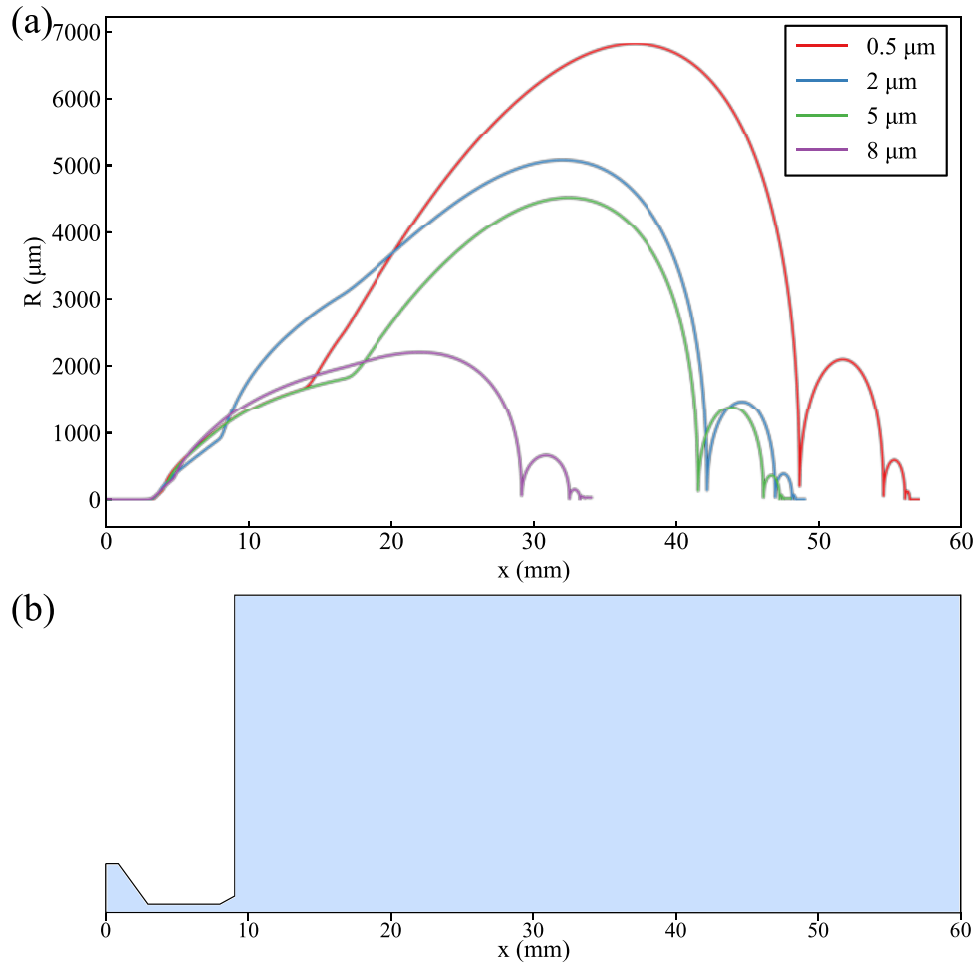


Figure 21: (a) Spatial evolution of the bubble radius for different initial radii R_0 . (b) Schematic of the converging-diverging nozzle geometry.

Fig. 22 illustrates the temporal evolution of the bubble radius for four different initial radii. During the early stage just prior to entering the negative-pressure region ($t < 0.25$ ms), the trajectories practically coincide. This overlap indicates that the external driving pressure dictates the initial phase, rendering the effect of the initial bubble size negligible. As time progresses, the influence of the initial radius becomes pronounced; specifically, a smaller R_0 results in a prolonged effective expansion duration and a substantially larger maximum radius.

Under the current conditions, the bubble with an initial radius of $R_0 = 0.5 \mu\text{m}$ reaches a maximum radius of approximately $6.8 \times 10^3 \mu\text{m}$, with its primary collapse concluding near $t \approx 1.5$ ms. Conversely, for $R_0 = 8 \mu\text{m}$, the maximum radius is only about $2.2 \times 10^3 \mu\text{m}$, and the end of the primary collapse is advanced to $t \approx 0.6$ ms. The trajectories for the intermediate initial radii fall between these two extremes, demonstrating a monotonic trend where both R_{max} and the collapse duration decrease as R_0 increases. Following the primary collapse, all cases undergo varying degrees of rebound oscillations. Notably, the amplitude and period of these secondary oscillations depend significantly on R_0 ; a smaller initial bubble radius yields a longer secondary oscillation period and a larger rebound amplitude.

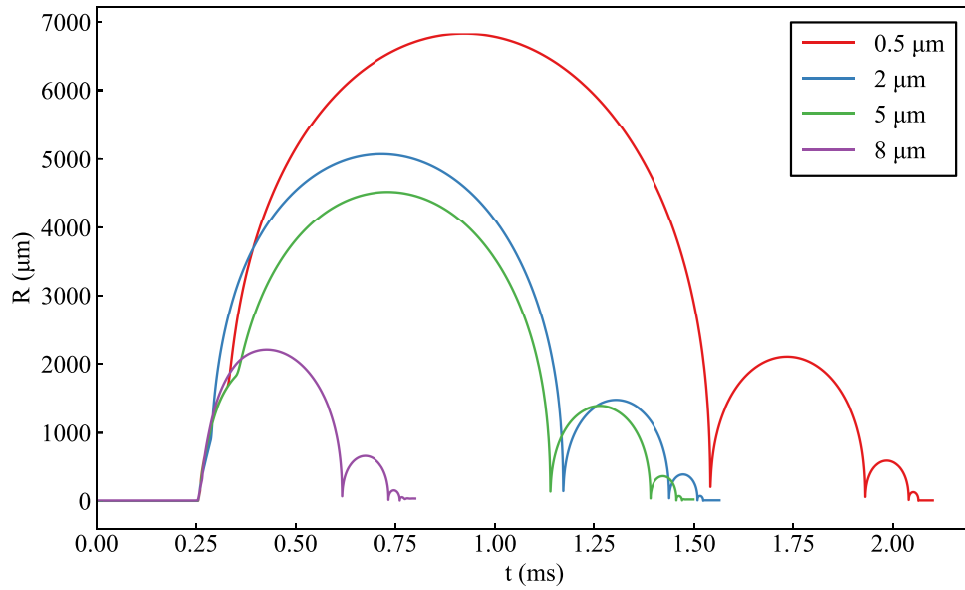


Figure 22: Temporal evolution of the bubble radius for different initial radii R_0 .

Fig. 23 illustrates the axial spatial distribution of the near-field pressure induced by bubble collapses under various initial radii. Aside from minor pressure fluctuations during the early expansion phase, the intense shock emissions are predominantly concentrated at the primary collapse and the subsequent 1–2 rebound stages, which spatially coincide with the minima of the bubble radius trajectories shown in Fig. 21. Specifically, the primary pulse peak for $R_0 = 8 \mu\text{m}$ is located at $x \approx 29 \text{ mm}$, while the peaks for the other cases shift progressively downstream: the primary peaks for 5 and $2 \mu\text{m}$ occur near $x \approx 41$ and 42 mm , respectively, and the peak for $R_0 = 0.5 \mu\text{m}$ is further delayed to the $x \approx 50 \text{ mm}$ region. In other words, under identical external flow conditions, the spatial location of the primary collapse correlates negatively with the initial bubble size; larger bubbles tend to collapse further upstream, whereas smaller initial bubbles push the collapse onset downstream. Correlating this with Fig. 24, it can be inferred that by modulating the inertial timescale and migration distance, the initial radius reshapes the spatiotemporal coordinates of the primary shock, even though the maximum peak intensities remain within a comparable order of magnitude. This finding provides direct engineering implications for identifying regional cavitation erosion susceptibility: the upstream transition section is more vulnerable to the primary collapse of larger bubbles, whereas the downstream exit region faces a higher risk from the implosion of smaller bubbles.

Fig. 24 illustrates the temporal evolution of the collapse-induced near-field pressure for various initial bubble radii. The impact spikes across the different cases exhibit a distinct temporal phase shift: the maximum pulse peak for $R_0 = 8 \mu\text{m}$ occurs earliest at $t \approx 0.6 \text{ ms}$; the peaks for $R_0 = 5 \mu\text{m}$ and $2 \mu\text{m}$ appear consecutively around $t \approx 1.15 \text{ ms}$; and the peak for $R_0 = 0.5 \mu\text{m}$ is further delayed to approximately $t \approx 1.5 \text{ ms}$. This pronounced temporal lag in the peak triggering as R_0 decreases demonstrates that the initial bubble size directly dictates the evolution cycle. Regarding the impact magnitude, the peak intensities universally reach the $O(10^4)$ bar order, yet they scale non-monotonically with the initial radius. The maximum peak is roughly $1.3 \times 10^4 \text{ bar}$ for $R_0 = 0.5 \mu\text{m}$, maximizes at over $1.4 \times 10^4 \text{ bar}$ for $R_0 = 2 \mu\text{m}$, and decreases to approximately $1.2 \times 10^4 \text{ bar}$ for both 5 and $8 \mu\text{m}$. These results indicate that the impact strength is dictated by the interplay between the maximum expansion radius and the collapse velocity. Because the initial radius modulates both of these factors, there exists an optimal initial radius range that maximizes the energy focusing efficiency.

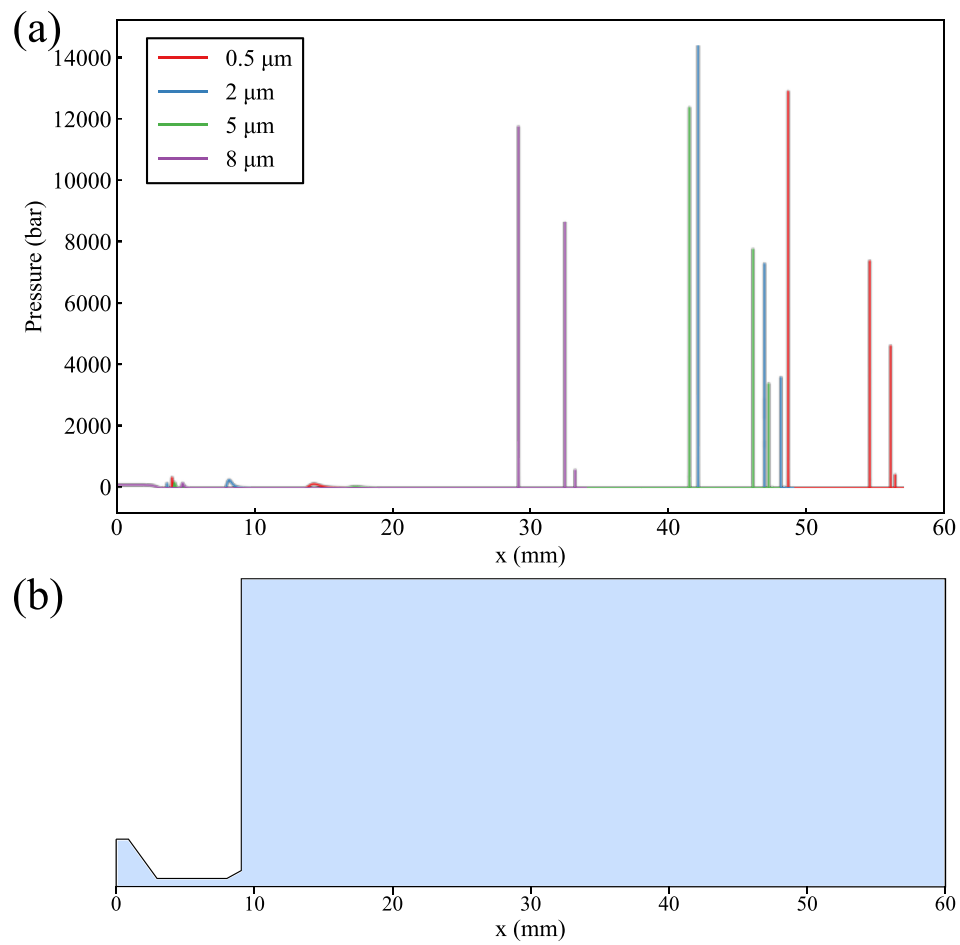


Figure 23: (a) Spatial evolution of the bubble radius for different initial radii R_0 . (b) Schematic of the converging-diverging nozzle geometry.

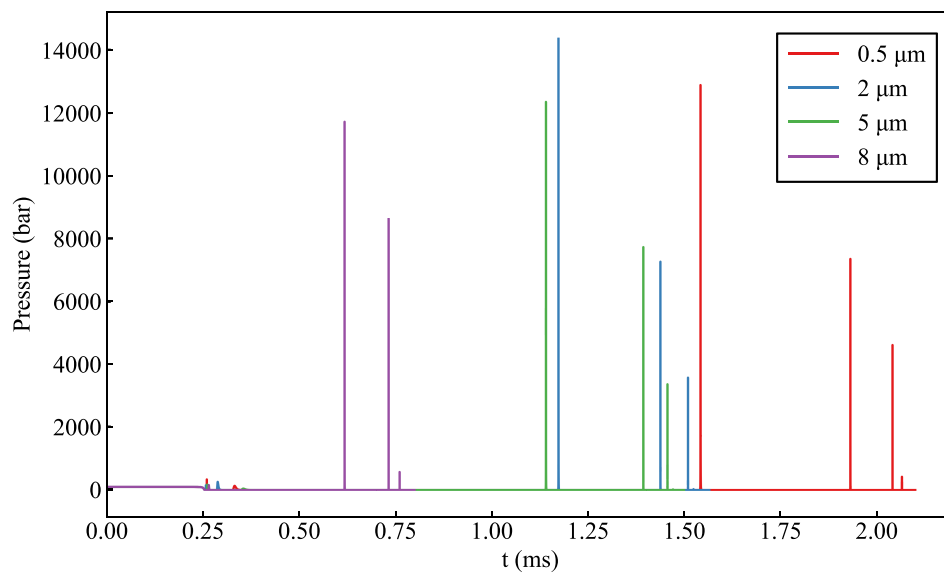


Figure 24: Temporal evolution of the bubble radius for different initial radii R_0 .

As illustrated in the parametric analysis of the initial bubble radius, the dynamic response of the bubbles exhibits a distinct non-monotonic relationship with respect to the maximum collapse intensity. Specifically, the bubble with an intermediate initial radius of $2\ \mu\text{m}$ generates the highest peak thermodynamic pressure (exceeding $14,000\ \text{bar}$), rather than the bubbles with smaller ($0.5\ \mu\text{m}$) or larger initial radii ($8\ \mu\text{m}$).

Fig. 25 presents the spatially coupled evolution of the bubble radius and pressure for two typical initial radii ($R_0 = 0.5$ and $8\ \mu\text{m}$). For the $R_0 = 8\ \mu\text{m}$ bubble, it expands to its maximum radius near $x \approx 25\ \text{mm}$ and undergoes its primary collapse at $x \approx 29\ \text{mm}$, emitting a maximum pressure pulse of approximately $1.2 \times 10^4\ \text{bar}$. Subsequently, its pressure oscillations attenuate rapidly by $x \approx 32\ \text{mm}$. In stark contrast, the $R_0 = 0.5\ \mu\text{m}$ bubble sustains its expansion over a considerably longer distance, reaching a maximum radius of $\sim 7.0 \times 10^3\ \mu\text{m}$ at $x \approx 40\ \text{mm}$. Its primary collapse is significantly delayed to the $x \approx 50\ \text{mm}$ region, generating a comparable peak pressure of roughly $1.3 \times 10^4\ \text{bar}$, with its secondary and tertiary collapse shocks localized near $x \approx 55\ \text{mm}$. This demonstrates that large and small bubbles localize their high-energy impact zones in entirely different regions—mid-channel vs. the downstream exit. Consequently, in a realistic polydisperse bubble population, the spatial superimposition of these size-dependent collapse events will generate severe shock loads over a substantially broadened axial range. This fundamental mechanism provides solid theoretical support for the expansive cavitation erosion zones frequently observed in complex cavitating flows.

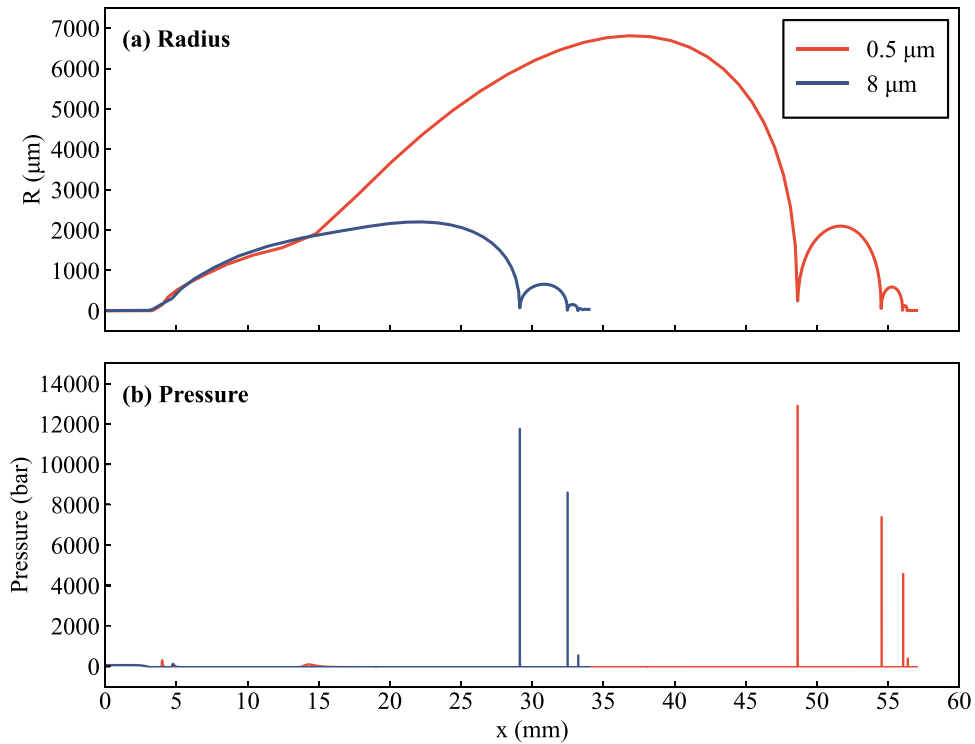


Figure 25: Comparison of the coupled spatial radius-pressure distribution between $R_0 = 0.5\ \mu\text{m}$ and $R_0 = 8\ \mu\text{m}$ cases. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

Fig. 26 presents the coupled time histories of the bubble radius and near-field pressure for two typical initial radii. The data clearly demonstrate that the $R_0 = 8\ \mu\text{m}$ bubble reaches its maximum expansion between $t \approx 0.2$ and $0.4\ \text{ms}$, and subsequently implodes at $t \approx 0.6\ \text{ms}$, emitting an intense primary pressure pulse. All subsequent rebound oscillations and their associated shock peaks are tightly confined to the $t < 0.8\ \text{ms}$

interval. In stark contrast, the $R_0 = 0.5 \mu\text{m}$ bubble undergoes a much more prolonged expansion phase, with its primary collapse substantially delayed to approximately $t \approx 1.5 \text{ ms}$. Although the primary impact peak for this smaller bubble is comparable in magnitude to that of the larger bubble (both reaching the $O(10^4)$ bar regime), its occurrence is significantly phase-shifted. Furthermore, during the $t \approx 1.5\text{--}2.1 \text{ ms}$ window, this microscopic bubble experiences two pronounced secondary rebounds, generating subsequent pressure peaks that still register in the thousands of bars. Ultimately, Fig. 26 confirms that by modulating the inertial timescale, the initial bubble radius dictates the temporal distribution of severe impact loads. Consequently, under identical external excitations, bubbles of different sizes will inflict destructive impacts on the boundary structures during entirely distinct temporal windows.

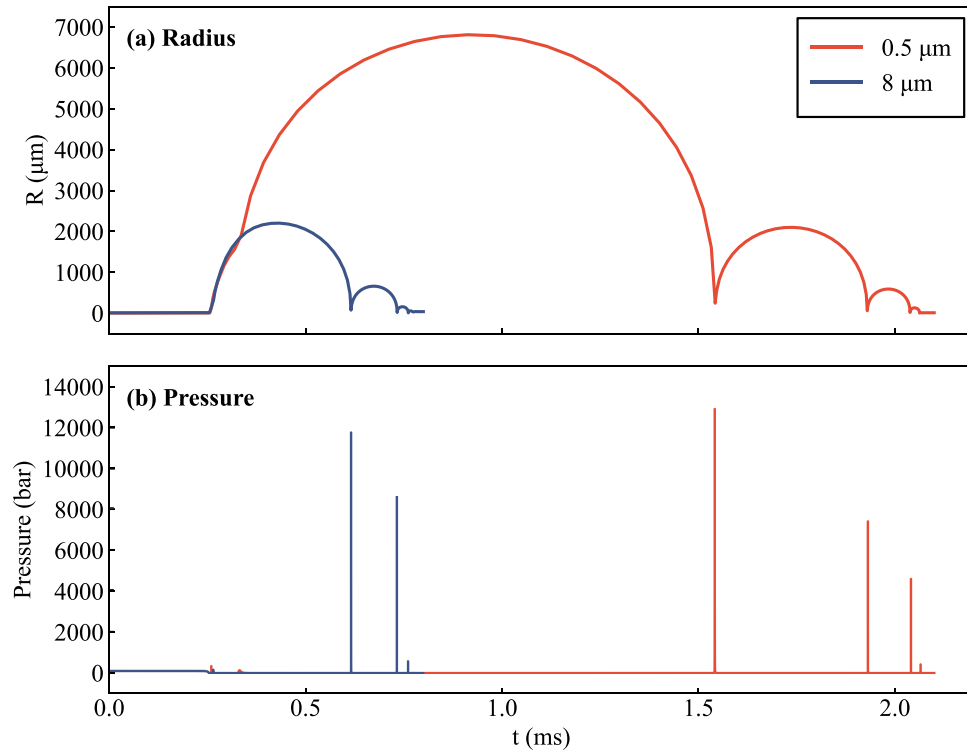


Figure 26: Comparison of the coupled temporal radius-pressure distribution between $R_0 = 0.5 \mu\text{m}$ and $R_0 = 8 \mu\text{m}$ cases. (a) Comparison of dimensionless radius; (b) Comparison of pressure.

Synthesizing the spatiotemporal comparisons from Figs. 25 and 26, the modulating effect of the initial bubble radius on the pulsation dynamics and transient impact loads manifests primarily in two dimensions. First, under identical flow field excitations, a smaller initial radius enables the bubble to achieve a vastly larger maximum radius and an extended migration distance, thereby pushing the spatiotemporal coordinates of its primary collapse closer to the downstream exit. Conversely, bubbles with a larger initial radius tend to implode rapidly within a shorter distance and duration upstream. Second, within the current parameter regime, the maximum impact peaks across disparate initial radii remain on a comparable order of magnitude. The severe shock events in the flow field are not dominated by a single bubble size; rather, they are synergistically contributed by the primary collapses of larger initial bubbles and the delayed, violent rebounds of smaller initial bubbles. Consequently, in the practical design of nozzles and microchannels, it is imperative to comprehensively consider the coupling relationship between the initial bubble size distribution and the local pressure environment to accurately evaluate regional cavitation erosion susceptibility.

4 Conclusion

This study establishes a numerical model to simulate the spatiotemporal evolution of bubbles within a converging-diverging microchannel nozzle. By utilizing CFD-computed single-phase flow data as the boundary input, the model successfully reproduces the transient bubble trajectories along the streamlines. Based on the pressure and velocity profiles extracted from five representative streamlines, the study reveals profound radial non-uniformity within the flow field: the momentum is more concentrated in the near-axis region, where fluid parcels experience a longer high-speed residence time and a shallower pressure trough.

Regarding the bubble dynamic response, a one-way coupling strategy was employed to systematically analyze the effects of the initial radial position, operating pressure level, and initial bubble radius. The following reproducible physical laws were extracted: When the initial bubble position is closer to the wall, it achieves a larger maximum radius, its primary collapse is significantly delayed, and the collapse location shifts further downstream. A higher operating ambient pressure triggers the expansion phase earlier and prolongs its duration, directly resulting in an extended first pulsation cycle and a more massive maximum bubble radius. A smaller initial bubble radius paradoxically enables the bubble to attain a larger maximum radius, driving its primary collapse closer to the nozzle exit. Fundamentally, these dynamic discrepancies originate from the variations in the duration of the external negative-pressure zones and the pressure recovery processes experienced along different streamlines. These mechanisms provide a direct theoretical basis for explaining and predicting cavitation erosion hotspots across different regions of the microchannel nozzle. It should be noted that the current model is limited to the evolution of a single bubble. Since actual cavitation in microchannels involves strong interaction effects among a bubble cloud, extending the present framework to multi-bubble cluster dynamics will be the focus of our future work.

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Availability of Data and Materials: The data that support the findings of this study are available from the Corresponding Author, [Hengzhu Bao], upon reasonable request.

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