

Air cavities as the primary drag-reduction mechanism in a superhydrophobic sphere

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Abstract: Precise control over air cavity generation in liquids is critical in numerous underwater marine systems, where cavity formation can largely reduce fluid resistance and minimize energy consumption. However, how the instability of the three-phase contact line at the microscopic scale induces the formation of macroscopic air cavities and thereby achieves drag reduction has not yet been fully elucidated. In this study, superhydrophobic (SHB) spheres with contact angles as high as 152° were fabricated using a spray-coating method. The liquid film evolution, flow field structure, pressure distribution, and hydrodynamic characteristics during the vertical immersion of spheres were systematically investigated through a combination of experiments and numerical simulations. The findings indicate that surface wettability is the key factor determining air cavity formation. The results show that the hydrophilic sphere did not generate an

air cavity under identical immersion conditions, whereas the SHB sphere formed a stable encapsulated air cavity. The underlying mechanism is that, owing to the low adhesion characteristic of the superhydrophobic surface, the adverse pressure gradient induces separation of the liquid film, which triggers pinning of the contact line near the equator of the sphere, prevents the liquid film from closing toward the poles, and thereby promotes air cavity formation. This air cavity transforms the solid–liquid no-slip boundary into an air–liquid free-slip boundary, significantly altering the flow field structure at the rear of the sphere. Under the optimal conditions within the investigated Reynolds number range, this mechanism effectively suppresses flow separation and vortex shedding, yielding a maximum observed drag reduction of approximately 90%. Additionally, the drag-reduction effect intensifies as the air cavity morphology transitions toward a stable streamlined shape with increasing Reynolds number. This research offers new insights into drag reduction via air cavities and provides a reliable theoretical foundation for the design of high-performance underwater drag-reducing surfaces.

Keywords: surface wettability; air cavity during water entry; drag reduction; hydrodynamics

1 Introduction

The process of structures entering water is a ubiquitous phenomenon observed both in natural and engineering contexts, ranging from biological predatory behaviors (such as waterfowl diving for fish or water striders skimming across the surface) to advanced technological applications (such as supercavitating torpedoes or seaplanes landing on water) [1–6]. Air cavities generated by the impact between a structure and the liquid surface play a crucial role in this complex fluid–structure interaction process. The presence of air cavities can effectively modify the flow field structure and stress distribution around the structure, thereby reducing the fluid drag experienced during underwater motion and improving its dynamic performance [7,8]. Therefore, in-depth investigation of the formation mechanisms and evolution of underwater air cavities holds significant theoretical and practical value for understanding natural phenomena and developing highly efficient underwater vehicles.

In recent years, inducing and controlling air cavities in water by regulating the wettability of solid surfaces has emerged as a frontier research topic [9–11]. As a key physicochemical property governing the behavior at solid–liquid–air three-phase interfaces [12–14], surface wettability directly influences the formation, morphology, and stability of air cavities. Ji et al. observed that superhydrophobic (SHB) surfaces promote stable air-cavity encapsulation during the vertical immersion of spheres with various degrees of wettability, achieving a drag-reduction efficiency of up to 89.5% [15]. Speirs et al. further demonstrated through systematic investigation that the morphology and closure mechanism of air cavities strongly depend on the apparent contact angle of the sphere [16]. To further enhance air cavity stability and drag-reduction performance, researchers have developed a series of surface micro-/nano-structuring and wettability heterogenization strategies. Zhao et al. fabricated parallel hydrophilic stripes on the surface of SHB steel spheres, leveraging the pinning effect of hydrophilic regions on the air–liquid interface to enhance the stability of streamlined air cavities and achieve a drag-reduction efficiency of 63.3% [17]. Watson et al. designed spheres with hemispherical hydrophobic/hydrophilic heterogeneous wetting patterns [18]. They found that surface heterogeneity can induce asymmetric cavitation and splashing behavior, generating lateral hydrodynamic forces. This provides a novel approach for passive trajectory control without active propulsion. Tao et al. achieved active control over air cavity morphology and trajectory during water entry by precisely regulating the proportion of SHB regions on the sphere surface, achieving a drag-reduction efficiency as high as 94.7% [19]. Inspired by biological organisms, such as water spiders, Yao et al. constructed SHB spheres with nano-needle-like surface structures [20]. This structure effectively traps an air layer, ensuring the formation of a durable and stable streamlined air cavity upon immersion in water. These studies collectively demonstrate that regulation of wettability through surface engineering is an effective approach for improving hydrodynamic performance and exploiting the drag-reduction potential of air cavities.

It is noteworthy that drag reduction achieved through surface wettability control exhibits broad applicability. Its benefits extend beyond the initial immersion phase and persist throughout the fully submerged motion stage. During this stage, SHB surfaces reduce drag by stabilizing the air film at the solid–liquid interface [21–25]. Chen et al. fabricated periodically oriented SHB microstructures on 316L stainless steel surfaces via combined laser ablation and chemical etching [26]. These structures demonstrated

a maximum drag reduction of 29.83% in microchannels, along with excellent droplet rebound performance and hydrophobic stability. Hu et al. demonstrated that coating the inner rotor of a Taylor–Couette flow apparatus with alternating SHB/hydrophilic rings can form a stable air ring firmly pinned to the surface, achieving significant and stable drag reduction of up to 77.2% without continuous air supply [27]. Yao et al. employed femtosecond-laser processing to create surfaces alternating between SHB and hydrophobic regions, enhancing air film stability and maintaining a 37% drag reduction at speeds of 2.3 m/s [28]. Cao et al. introduced parallel hydrophilic stripes onto a SHB surface with random roughness [29]. The three-phase contact lines formed by these stripes confine the air layer and prevent its escape from the surface microstructure. This enhances the stability of the air layer while significantly improving the surface resistance to water impact and its drag-reduction performance. Fakhri et al. developed a simple, low-cost, and stable SHB coating that achieves approximately 40% drag reduction through its micro–nano composite structure [30]. Taghvaei et al. developed a dual-scale SHB coating composed of micron- and nanometer-sized alumina particles, achieving a maximum drag reduction of 77% over a broad Reynolds number range (from 4.7×10^4 to 2.35×10^5) [31].

Although substantial progress has been achieved in utilizing surface wettability for drag reduction, the current research paradigm still exhibits certain limitations. Most existing studies focus primarily on final drag-reduction performance and macroscopic air cavity morphology by modifying surface parameters, such as contact angle and microstructure. However, a clear understanding of how microscopic contact line instability leads to macroscopic air cavity formation is still lacking. Previous research has not established a clear causal link between contact line pinning and instability, liquid film separation, and air cavity generation. As a result, the intrinsic relationship between microscopic interfacial behavior and macroscopic flow characteristics remains insufficiently explained. This incomplete mechanistic understanding limits further optimization of surface design, prediction of air cavity behavior, and maximization of drag-reduction efficiency.

Accordingly, this study focuses on SHB spheres to investigate the formation mechanism of air cavities and the associated drag-reduction mechanism during vertical water entry. A combined approach integrating high-speed imaging experiments and high-precision numerical simulations is employed to establish a multi-scale, multi-physics coupled analysis method. Based on this method, we analyze the motion

characteristics of the contact line during liquid film spreading on the surface, as well as the effects of local pressure distribution and adverse pressure gradient evolution on liquid film separation and contact line instability. The physical mechanism by which surface wettability regulates air cavity generation and evolution is revealed. Furthermore, quantitative correlations between air cavity geometry, stability, and drag-reduction efficacy are established to precisely evaluate the impact of air cavities on total fluid resistance. These findings provide a mechanistic explanation for superhydrophobicity-induced air cavity drag reduction and offer a theoretical basis for the design of high-performance drag-reducing surfaces.

2 Experimental section

The spherical vertical water-entry test system employed in this study is shown in Fig. 1. It primarily consists of four main modules: the experimental water tank, the sphere release and positioning mechanism, the high-speed imaging system, and the illumination system. The water tanks are custom-made from high-transparency acrylic panels to ensure unobstructed optical observation of the entry process. To minimize boundary effects on the entry flow field and air cavity morphology while accommodating observations at different depths, two tank sizes were used. The first was a short tank (internal dimensions: 20 cm $\times$ 20 cm $\times$ 30 cm), primarily used to capture the initial interface phenomena during entry. The second was a long tank (internal dimensions: 20 cm $\times$ 20 cm $\times$ 100 cm), designed to investigate the long-term evolution of the sphere underwater motion and air cavities. All tanks were equipped with a rubber buffer layer at the bottom to prevent equipment damage from sphere impacts. The imaging system employed Revealer high-speed cameras (X213) operating at 5000 fps to resolve transient interfacial dynamics during immersion. To achieve high-contrast images, a 400 W light source combined with a diffuse softbox was used to provide uniform, shadow-free illumination in the observation area.

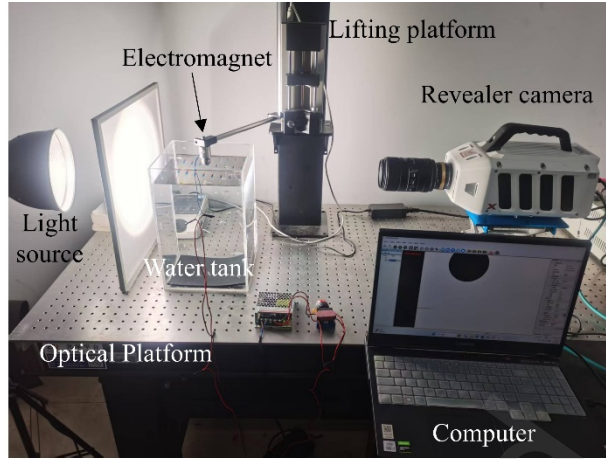


Fig. 1 Diagram of the test setup for the sphere entering water.

Table 1 Experimental Parameters

Parameters	Symbols	Value/Range	Unit
Sphere diameter	D	10, 20	mm
Sphere mass	M	0.00411, 0.0328	kg
Impact velocity	U_0	1.09—3.1	m/s
Water density	ρ	1000	kg/m ³
Dynamic viscosity	μ	0.001	Pa·s
Gravitational acceleration	g	9.8	m/s ²
Reynolds number	$Re=(\rho U_0 D)/\mu$	$1.4 \times 10^4—3.54 \times 10^4$	

Free release of the sphere without initial velocity was achieved by controlling the energization of the electromagnet. The sphere initial entry velocity U_0 was adjusted by varying the release height H , with its theoretical value calculated as $U_0 = \sqrt{2gh}$, where $g \approx 9.8 \text{ m/s}^2$. The key physical parameters involved in the experiment, including the sphere diameter D , sphere mass, impact velocity range U_0 , fluid properties, and corresponding Reynolds number Re range, are summarized in Table 1.

To elucidate the intrinsic physical mechanisms governing air cavity formation during the immersion of SHB spheres, a computational fluid dynamics (CFD) model was developed (see Supplementary Information, Section S1).

3 Results and discussion

To systematically investigate the influence of surface wettability on the mechanism

of air cavity formation during water entry, the dynamic behavior of the original (hydrophilic) spheres and the SHB spheres is compared. An SHB functional layer was fabricated on the surface of the original spheres using a spray-coating process, as shown in Fig. 2(a). Characterization of the treated sphere surface using a contact angle goniometer showed that the SHB spheres exhibited a water contact angle (WCA) of $\sim 152^\circ \pm 2^\circ$, whereas the original spheres exhibited a WCA of only $\sim 66^\circ \pm 3^\circ$. This confirms that surface modification successfully transformed the spheres from hydrophilic to SHB. Additionally, the chemical composition of the SHB coating was analyzed using energy-dispersive spectroscopy (EDS). Scanning electron microscopy (SEM) images of the coating microstructure (Fig. 2(b)) reveal a hierarchical micro–nano composite structure consisting of uniformly distributed micrometer-scale protrusions and nanoscale roughness.

Under identical experimental conditions (initial velocity $U_0 \approx 1.09$ m/s, diameter $D \approx 20$ mm), the two types of spheres exhibited distinctly different water-entry dynamics (Fig. 2(c)). Experimental observations indicate that the original sphere failed to form an effective air cavity upon impact with the liquid surface and instead generated a distinct vertical upward jet above the free liquid surface. By contrast, the SHB sphere successfully induced stable air cavity entrainment during water entry. High-speed imaging analysis indicates that the evolution of the air cavity in the SHB sphere can be systematically divided into four characteristic stages: (i) initial flow development, (ii) air cavity opening and axial stretching, (iii) air cavity neck closure, and (iv) final air cavity instability and collapse. This staged evolution clearly demonstrates the pivotal role of surface wettability in regulating the formation and evolution of air cavities.

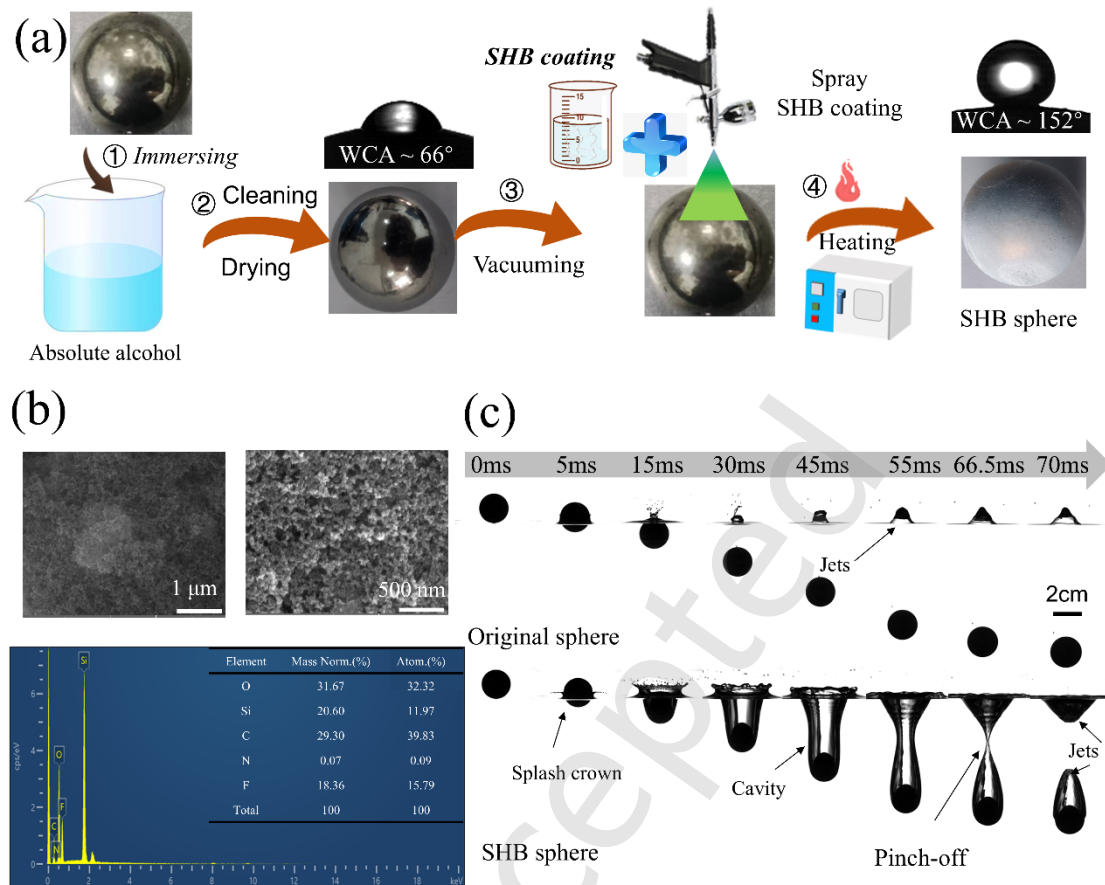


Fig. 2 Preparation process of superhydrophobic spheres and water entry behavior of spheres with different wettability. (a) Preparation process of superhydrophobic spheres; (b) Scanning electron micrograph and EDS spectra of superhydrophobic coating; (c) Water entry of original spheres and superhydrophobic spheres with an initial impact velocity of $U_0 \sim 1.09$ m/s.

During the initial impact stage, an initial liquid-film jet forms during the impact of both the original and superhydrophobic spheres on the free liquid surface. However, comparative analysis reveals that the original and SHB spheres exhibit fundamental differences in the dynamics of the subsequent liquid film evolution (Fig. 3(a)). The liquid films form and adhere to the surfaces of both types of spheres during the initial impact phase. However, as the process evolves, the liquid film on the surface of the original sphere continues to climb along the spherical surface, eventually closing at the north pole to form a complete liquid film enclosure that prevents air cavity formation. By contrast, the liquid film on the surface of the SHB sphere undergoes significant flow separation after reaching the vicinity of the equator, forming a distinctive feather-like splash crown with a central opening. This opening serves as a channel for ambient air entrainment. The external air is continuously drawn inward by the pressure difference as the sphere continues to descend, eventually forming a complete air cavity. The results indicate that the critical condition for liquid film separation is closely related to

dimensionless parameters, which can be actively controlled by increasing the Weber number ($We = \rho U_0^2 D / \sigma$) or enhancing the apparent contact angle [32–34]. These findings reveal the key mechanism through which surface wettability determines air cavity formation by modulating liquid film instability and separation. Specifically, under high-speed impact conditions (high We number), fluid inertia dominates the liquid film separation, allowing air cavities to form even on ordinary hydrophobic or even hydrophilic surfaces (Figs. S3-S5). Conversely, the superhydrophobic surface plays an enabling role in the low-speed water-entry regime, significantly lowering the critical velocity threshold for cavity generation by amplifying the dynamic contact angle hysteresis and reducing the energy barrier required for interfacial separation.

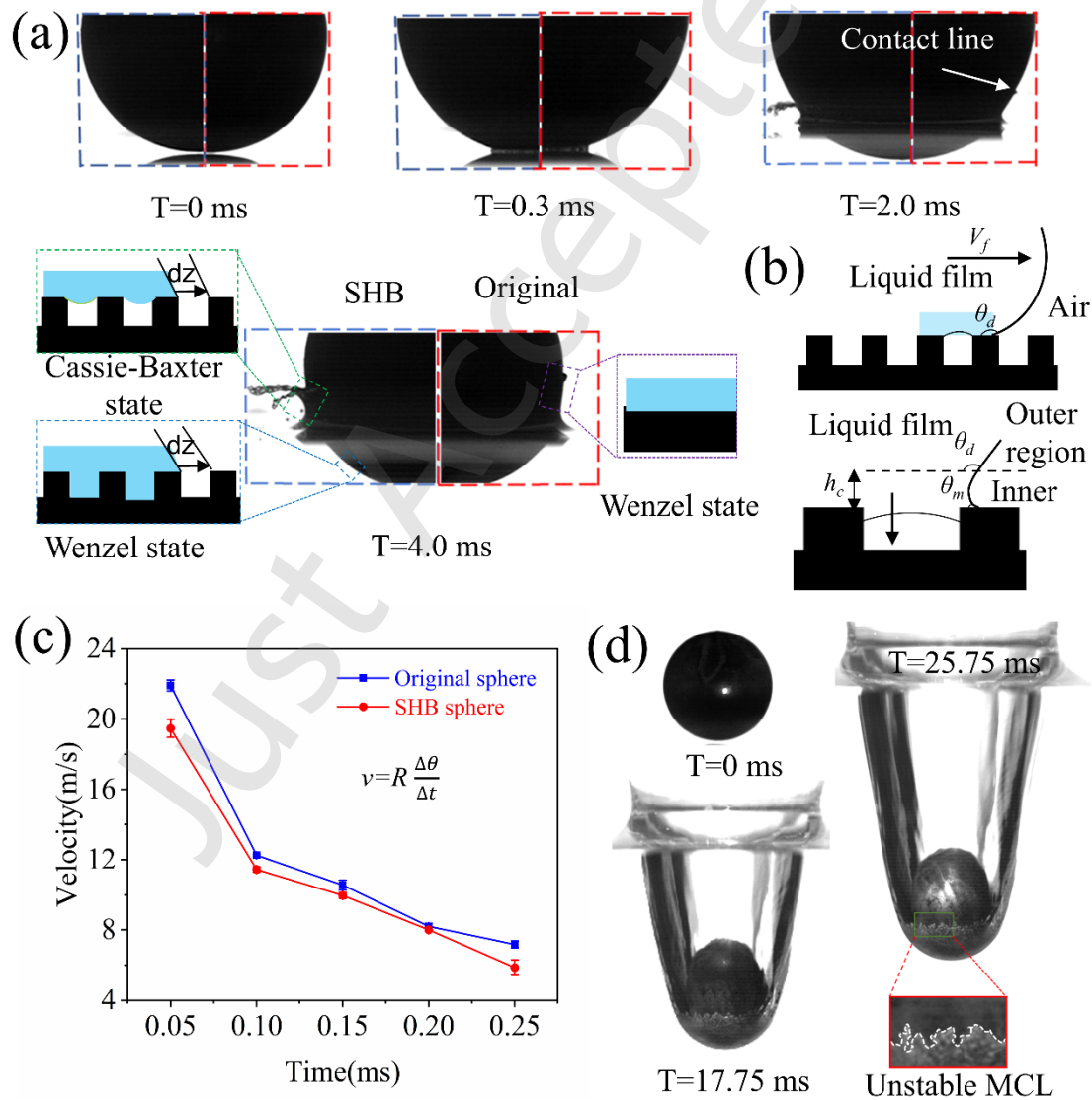


Fig. 3 Mechanism of air cavity formation under water immersion. (a) Effect of surface wetting state on liquid film climbing; (b) Dynamic wetting model of liquid films on structured surfaces, V_f is the wetting velocity of the liquid film.; (c) Relationship between the movement speed of the three-phase contact line on the sphere surface and time; (d) Formation of air cavities in superhydrophobic spheres under water immersion and instability of the three-phase contact line.

The kinetic behavior of liquid film evolution is fundamentally governed by the motion of the solid–liquid–air three-phase contact line. Experimental observations (Fig. 3(c)) indicate that, during the initial entry stage, the three-phase contact line speed on both spheres is significantly higher than the sphere descent velocity, demonstrating rapid upward migration along the spherical surface. However, the contact line on the original spheres moves significantly faster than that on the SHB sphere due to its higher surface energy. The contact line motion on the SHB sphere exhibits unique temporal evolution: its velocity gradually decreases under the combined effects of viscous dissipation and surface microstructure-induced pinning. Eventually, the contact line stagnates near the equator, forming a stable pinning state (Fig. 3(d)). This difference in contact line dynamics results in two distinct water-entry modes: the original sphere fails to form an air cavity because its contact line continues moving until it closes at the pole, whereas the SHB sphere maintains an open air cavity due to contact line pinning.

A typical sawtooth instability pattern is clearly visible upon close examination of the contact line morphology (Fig. 3(d)). This instability causes significant buckling deformation of the contact line pinned to the SHB surface, resulting in a sharp decrease in its local radius of curvature. Based on the Young–Laplace equation, an increase in local curvature directly leads to a corresponding rise in curvature-induced pressure. When the curvature pressure exceeds the immersion pressure generated by the combined effects of dynamic impact and hydrostatic pressure, liquid penetration into the surface microstructure is effectively suppressed [35]. This enables the stable entrapment of air within the micro–nano structure, forming a composite solid–liquid–air interface with a high apparent contact angle. This mechanism reveals the spatiotemporal evolution of the wetting state during sphere impact: the continuously wetted lower region primarily exhibits Wenzel wetting, while the upper region where the liquid film separates maintains Cassie–Baxter wetting. Consequently, a typical hybrid wetting state is formed, as illustrated in Fig. 3(a).

To elucidate the instability mechanism and wetting state transition of a moving three-phase contact line, an analytical model based on capillary driving force theory was developed. As depicted in Fig. 3(b), the evolution of the interfacial morphology within the inner region is primarily governed by the combined effects of capillary forces and viscous friction. Conversely, in the outer region—extending beyond the characteristic capillary length (h_c) from the solid surface—macroscopic equilibrium is dictated by capillarity coupled with external forces, including inertia and gravity. The

global stability of the advancing liquid film fundamentally depends on boundary constraints and geometric continuity between the outer interface and this inner region. While the unconstrained outer interface can continuously adjust its local curvature to preserve macroscopic balance, the spatially constrained inner region cannot readily accommodate these morphological changes. This structural mismatch causes instability to originate from perturbations within the inner region. Therefore, understanding wetting failure necessitates a focused examination of this localized domain. The dominant driving force for contact line motion in this region can be expressed as [32]:

$$F = \gamma_{lv} (\cos \theta_m - \cos \theta_d) \quad (1)$$

where F is the capillary driving force, γ_{lv} is the liquid–air surface tension, θ_d is the dynamic contact angle, and θ_m is the microscopic contact angle. This equation reveals how the surface microstructure influences the capillary driving force through modulation of θ_d . The capillary driving force impedes liquid spreading when θ_d significantly exceeds the equilibrium contact angle, thereby promoting contact line instability and transition toward the Cassie–Baxter state. Fundamentally, this microscale interfacial instability manifests as macroscopic kinematic retardation during complex fluid–solid interactions. During the entry of a superhydrophobic sphere into water, the dynamic evolution of the three-phase contact line and its eventual pinning position are governed by the competition between the spreading driving force and the pinning resistance. The spreading driving force mainly originates from the radial component of the impact momentum of the liquid film and the interfacial capillary pressure, whereas the pinning resistance arises from three aspects: the high-energy-barrier physical obstacles provided by the micro/nano hierarchical structures on the surface, the kinetic energy dissipation due to viscous shear in the three-phase region, and the dynamic contact angle hysteresis. As the contact line advances, the spreading driving force gradually decreases. When it falls to a level insufficient to overcome the local maximum pinning resistance, the critical condition for force balance is satisfied, causing the contact line to cease its evolution and form a stable pinned state locally.

The evolution dynamics of the liquid film are the key physical processes governing the formation of air cavities. However, in sphere–free-surface impact experiments, strong interfacial perturbations and capillary waves cause severe deformation of the liquid surface, which interferes with direct observations of the liquid film behavior in critical regions. To overcome this limitation, a simplified droplet–impact model

involving droplets impacting spheres and flat surfaces was employed. This design effectively avoids observational difficulties arising from complex free-surface deformations and successfully reproduces crown-like splash structures similar to those observed during water entry. More importantly, this configuration enables clear visualization and quantitative analysis of the complete evolution process, from initial impact to early liquid film development. This provides a reliable experimental foundation for elucidating the intrinsic connection between liquid film separation and air cavity formation.

To establish a dynamic correlation between droplet-on-sphere impact and sphere water entry, a controlled experiment was designed. A droplet of fixed diameter was released from a height of 6 cm to impact the surface of spheres with different wettability characteristics. By introducing the dimensionless radius ratio $\alpha = r/R$ (where r is the droplet radius and R is the sphere radius), a scaling relationship between the droplet–sphere system and the sphere–water-entry system was established. Upon impact, the droplet undergoes rapid radial spreading along the spherical surface (Fig. 4(a)). According to the principle of volume conservation for incompressible fluids, the blue region (the initial crown of the droplet) and the pink region (the radially spreading thin film) in the figure have equal volumes, with $h \ll R$, $\theta = \sqrt{L / [(1 + \alpha)r]}$ [36]. The spreading radius of the liquid film can therefore be determined as $r_c = \sqrt{3tU / [(1 + 1/\alpha)r]}$ [37].

The surface wettability exerts a decisive influence on the dynamic behavior of liquid films during droplet impact. When a droplet impacts an original sphere, the liquid film maintains continuous contact with the spherical surface and spreads smoothly across it. By contrast, on an SHB sphere, the liquid film undergoes significant flow separation during spreading, forming a characteristic crown-like structure. Theoretical analysis indicates that the formation of such crown structures requires the radial expansion velocity at the liquid film front to exceed the rate of increase of the solid–liquid contact area. During this process, the liquid film edge experiences vertical lift generated by aerodynamic effects, including lubrication forces acting on the microstructure gaps and suction forces acting above the liquid film (Fig. 4(b)). The synergistic action of these forces drives the separation of the liquid film from the substrate and stabilizes an air layer beneath the film. This mechanism directly explains air cavity formation when SHB spheres enter water: the three-phase contact line

separates before reaching the sphere apex due to the premature formation of a corona-like structure. This maintains a stable open air cavity at the sphere top, providing suitable initial conditions for the subsequent development of a complete air cavity.

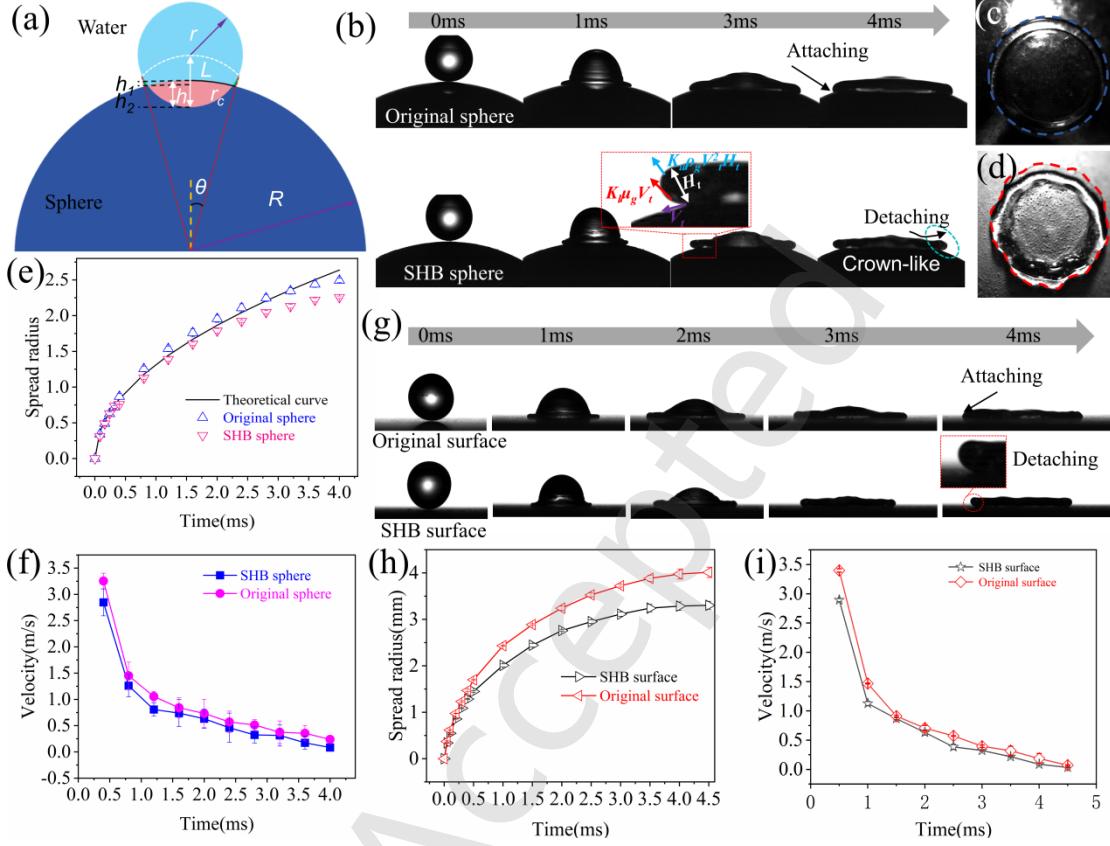


Fig. 4 Evolution of liquid films upon droplet impact on solids with different wettability. (a) Schematic of droplet impact on a sphere, h_1 is the height of the spherical cap of the sphere, and h_2 is the height of the spherical cap of the droplet; (b) Evolution of liquid films upon droplet impact on pristine and superhydrophobic spheres; Schematic top views of a droplet impacting (c) an original sphere and (d) an SHB sphere; (e) Time-dependent variation of contact line radius after droplet impact on a sphere; (f) Time-dependent variation in contact line displacement velocity after droplet impact on a sphere; (g) Evolution of liquid films upon droplet impact on pristine and superhydrophobic surfaces; (h) Time-dependent variation in contact line radius after droplet impact on a flat surface; (i) Time-dependent variation in contact line displacement velocity after droplet impact on a flat surface.

Droplet impact on a solid sphere is a complex hydrodynamic process involving multiphase interfaces, transient flow, and energy conversion. This process is governed by energy conservation and can be expressed as [38,41]:

$$E_{k0} + E_{s0} + E_{g0} = E_k + E_s + E_g + W \quad (2)$$

where E_{k0} , E_{s0} and E_{g0} represent the initial kinetic energy, surface energy, and gravitational potential energy, respectively. E_k , E_s and E_g denote the corresponding kinetic energy, surface energy, and gravitational potential energy at any given moment during the impact process, while W represents viscous energy dissipation.

The contact angle influences the surface energy. According to Young's equation [39], a smaller solid–liquid interfacial tension (γ_{sl}) on the original sphere corresponds to a larger surface energy, which facilitates droplet wetting. Conversely, the larger γ_{sl} on an SHB surface results in a smaller surface energy. This leads to an extremely low adhesion between the droplet and the surface, reducing the capacity of the droplet to wet the surface and spread.

During the impact process, the surface energy evolves over time according to [40]:

$$E_s(t) = \gamma_{gl}A_{gl}(t) + (\gamma_{sl} - \gamma_{sg})A_{sl}(t) \quad (3)$$

By substituting Young's equation into Eq. (3), the surface energy can be expressed as:

$$E_s(t) = \gamma_{gl}A_{gl}(t) - \gamma_{gl} \cos \theta A_{sl}(t) \quad (4)$$

where E_s represents the surface energy, γ_{gl} is the air–liquid interfacial tension, γ_{sl} is the solid–liquid interfacial tension, γ_{sg} is the solid–air interfacial tension, A_{gl} is the air–liquid interfacial area, A_{sl} is the solid–liquid contact area, and θ is the contact angle.

When a droplet impacts an SHB sphere, the expansion of the solid–liquid contact area increases the surface energy, raising the system energy state. According to the law of conservation of energy, the system tends to minimize this contact area, thereby promoting the separation of the liquid film from the SHB spherical surface.

In the later stage of droplet impact, the energy conservation equation becomes [40]:

$$We \left[1 - \frac{4}{\sqrt{R_e}} \frac{R(t)^2}{R_0^2} \right] = 6 \frac{R_b^2}{R_0^2} (1 - \cos \theta)(1 - \cos \theta_1) - 12 \quad (5)$$

where $R(t)$ represents the spreading radius, R_0 is the initial droplet radius, and R_b is the sphere radius. The central angle corresponding to the spreading arc is defined as $\theta_1 = R(t)/R_b$.

As derived from Eq. (5), the original sphere, with its smaller contact angle, exhibits a lower surface energy contribution, which facilitates greater droplet spreading. Furthermore, its low solid–liquid interfacial tension and high adhesion on the original sphere surface facilitate wetting, resulting in the formation of a larger solid–liquid contact area. This significantly reduces the surface energy, resulting in stable droplet

deposition, as illustrated in Fig. 4(c).

By contrast, when a droplet impacts an SHB sphere, the larger surface energy contribution limits the maximum spreading radius and promotes liquid retraction to minimize the solid–liquid contact area. Furthermore, the sphere curvature introduces a Laplace pressure difference that generates an additional normal surface tension component at the liquid film rim, increasing the possibility of the film to detach from the surface. Particularly on SHB spheres, the curvature also promotes the escape of air trapped beneath the liquid film, weakening the supporting effect of the air layer and further facilitating film separation.

When a droplet impacts an SHB sphere, shear arises between the spreading liquid film and the trapped air layer. The velocity and pressure differences across the air–liquid interface may trigger the Kelvin–Helmholtz (KH) instability, whose growth rate is given by [41]:

$$\omega^2 = \frac{\sigma}{\rho} k^3 - \frac{\Sigma_G}{\rho} k^2 \quad (6)$$

where ω is the complex growth rate, k is the wavenumber, σ is the liquid surface tension, ρ is the liquid density, and Σ_G represents the aerodynamic stress.

This instability induces rim undulations, as observed in Fig. 4(d), which promote the detachment of the film from the spherical surface. Furthermore, the microscale roughness of the SHB surface traps air to form a stable film, substantially reducing the effective solid–liquid contact area. This substantially decreases the viscous shear resistance and the associated energy dissipation, thereby facilitating droplet separation from the sphere surface.

To further quantify the influence of surface wettability on liquid film dynamics, this study systematically analyzed the temporal evolution of the liquid film spreading radius and the migration velocity of the three-phase contact line. During the initial post-impact stage, the liquid film spreading radii on the original and SHB spheres largely coincide with each other, showing strong agreement with theoretical predictions based on the inertia-dominated model (Fig. 4(e)). This indicates that liquid film expansion at this stage is primarily governed by inertial forces, while the influence of surface wettability remains negligible. However, as time progresses, the liquid film behaviors on the two surfaces diverge significantly: the spreading radius on the SHB surface lags markedly behind that on the original surface, and the corresponding contact line velocity decays more rapidly (Fig. 4(f)). This late-stage divergence highlights the

critical role of surface wettability in governing liquid film evolution through capillary forces and viscous dissipation mechanisms.

A comparison of the droplet impact processes on different surfaces (Fig. 4(g)) clearly demonstrates the significant influence of surface wettability on the liquid film evolution patterns. On the original hydrophilic surface, the droplet forms a continuously spreading liquid film that remains in constant contact with the surface after impact. By contrast, on the SHB surface, the liquid film front undergoes evident flow separation, forming a characteristic crown-like structure. This phenomenon is attributed to the aerodynamic effect induced by the microstructured SHB surface, where a stable air layer forms beneath the liquid film edge, effectively preventing direct contact between the film and the substrate.

Quantitative characterization of the spreading dynamics (Fig. 4(h)) reveals that, during the initial impact phase, the spreading radii of the liquid films on both surfaces are nearly identical, indicating that inertial forces dominate at this stage. Subsequently, the spreading radius on the SHB surface becomes significantly smaller than that on the hydrophilic surface, demonstrating pronounced spreading inhibition. Correspondingly, the migration velocity of the three-phase contact line on the SHB surface exhibits a faster decay (Fig. 4(i)), further confirming that surface wettability exerts a decisive influence on the late-stage evolution of the liquid film by modulating capillary forces and viscous dissipation.

Based on the kinetic characteristics of liquid film evolution, this study provides a theoretical interpretation of the physical mechanism underlying the formation of corona-like structures when droplets impact SHB surfaces. The formation of such structures fundamentally involves a wetting-state transition triggered by instability of the moving contact line, which represents a typical dynamic wetting transition problem. According to the moving contact line theory, the apparent contact angle and contact line velocity satisfy the following constitutive relationship [42]:

$$\theta_{ap}^3 - \theta_e^3 = 9C_a \ln \frac{al_0}{l_i} \quad (7)$$

where θ_{ap} denotes the apparent contact angle, θ_e represents the equilibrium contact angle, $C_a = \mu U / \gamma$ is the dynamic capillary number, U is the contact line velocity, γ is the interfacial tension, and a is a parameter determined by boundary conditions. Additionally, l_0 is the macroscopic characteristic length scale, and l_i is the microscopic characteristic length scale. When the contact line velocity exceeds the critical capillary

number, the advancing contact line becomes unstable. At this stage, the air phase invades the space between the liquid film and the substrate, forming an air cushion layer. This destabilization causes the leading edge of the liquid film to detach from the substrate, thereby inducing the formation of a corona-like structure.

When the droplet impacts a smooth surface, flow separation occurs when the radial expansion velocity at the liquid film front significantly exceeds the rate of increase in the solid–liquid contact area. Under this condition, the centrifugal force generated by inertial effects at the film front overcomes the adhesive force between liquid and solid, causing the film to detach from the substrate. This flow separation is sustained by the synergistic effects of aerodynamic lift and surface tension redistribution, ultimately forming a stable, crown-like geometry. The required condition can be expressed as [37]:

$$\frac{Du}{Dt} = -\frac{\partial p}{\partial x} + Re_d^{-1} \nabla^2 u \geq 0 \quad (8)$$

The dimensionless separation time t_e is given by [37]:

$$c_1 Re_d^{-1} l_e^{-1/2} + Re_d^{-2} Oh_d^{-2} = c^2 l_e^{3/2} \quad (9)$$

where u is the velocity, p is the pressure, $c_1 = \sqrt{3}/2$, $c = 1.1$, $Re_d = \rho_l r_2 U / \mu_l$, $Oh_d = \mu_l / \sqrt{\rho_l r_2 \gamma_l}$, and ρ_l , r_2 , U , μ_l , and γ_l represent the droplet density, droplet radius, impact velocity, droplet viscosity, and surface tension, respectively.

When a droplet impacts a rough surface, the mass and momentum conservation equations describing the liquid film edge are respectively [43]:

$$\begin{aligned} \gamma \frac{\pi}{4} \frac{db^2}{dt} &= [u(s,t) - v] h(s,t), \quad \frac{ds}{dt} = v \\ \gamma \frac{\pi b^2}{4} \frac{dv}{dt} &= [u(s,t) - v]^2 h(s,t) - (1 + \beta) We^{-1} \end{aligned} \quad (10)$$

where $u(s,t)$ represents the radial velocity distribution within the liquid film, $h(s,t)$ denotes the film thickness distribution, β is the capillary resistance coefficient, and γ is the shape factor. For SHB surfaces, where the liquid film edge detaches from the substrate and forms a perfect circular interface, $\beta = 1$ and $\gamma = 1$ are assumed. For hydrophilic surfaces, where the edge maintains a semicircular geometry in contact with the substrate, $\beta = 0$ and $\gamma = 1/2$ are assumed. From this formulation, it is evident that when a droplet impacts an SHB surface, the surface tension term becomes $-2We^{-1}$, whose magnitude is twice that of a hydrophilic surface ($\beta = 0$). This stronger negative term represents a significantly enhanced inward contraction force directed toward the droplet center. Physically, because the leading edge of the liquid film does not contact

the substrate on SHB surfaces, the entire perimeter of the air–liquid interface is exposed. This allows surface tension to act fully along this complete free boundary, strongly driving minimization of the interfacial area. This pronounced inward contraction not only suppresses radial spreading but also generates an upward pressure gradient at the base of the liquid film. To maintain mechanical equilibrium, the liquid ring transitions from a substrate-adhering meniscus to a detached, circular-sectioned structure. This transformation results in the highly stretched crown-like structure observed experimentally at the liquid film edge.

To further elucidate the physical mechanism of air cavity formation during immersion of SHB spheres, numerical simulations were conducted to examine the transient pressure field evolution and flow field structure development (Fig. 5). The simulation results indicate that a distinct high-pressure region forms beneath the sphere during the initial impact stage (Fig. 5(a)). Simultaneously, the three-phase contact line ascends along the spherical surface under the combined effects of surface tension and hydrostatic pressure. Significant boundary layer separation occurs between the water flow and the sphere due to the low adhesion of the SHB surface. This separation process generates localized vortex structures at the leading edge of the contact line and induces the formation of a persistent low-pressure region. As the contact line continues to migrate, boundary layer separation intensifies, and the low-pressure region expands. This results in a characteristic reverse pressure gradient beneath the contact line, with the pressure decreasing away from the contact line. This reverse pressure gradient generates a fluid squeezing effect toward the low-pressure region, exerting a downward drag force on the contact line and inducing outward deformation at the air–liquid interface. As the sphere descends further, the flow field gradually stabilizes: the high-pressure region beneath the sphere expands but weakens, while the low-pressure region gradually contracts, resulting in a corresponding decrease in the reverse pressure gradient. This dynamic equilibrium ultimately causes permanent pinning of the contact line near the equator, preventing closure at the pole and thereby forming a stable air cavity.

Fundamentally, this contact line pinning position emerges from a dynamic competition between the spreading driving force and local pinning resistance. Comparative simulations demonstrate that the reduced liquid adhesion on SHB surfaces allows the adverse pressure gradient to trigger premature flow separation. This creates an 'air cavity isolation effect' that rapidly dissipates the upward spreading momentum,

firmly anchoring the contact line near the equator. Conversely, stronger liquid attachment on intermediate hydrophobic surfaces delays separation, forcing the contact line significantly higher before pinning (Figs. S6-S10). Therefore, superhydrophobicity lowers the energy barrier for early contact line separation, which is critical for sustaining a fully encapsulated cavity. Once the stable cavity is formed, further analysis of the flow field indicates that ambient air is rapidly entrained behind the descending sphere. Consequently, the air velocity inside the cavity becomes significantly higher than that of the surrounding liquid. This intense velocity shear across the air-liquid interface essentially drives the dynamic evolution of vortex structures, as illustrated in Figs. 5(b) and (c). These vortex dynamics are quantified through the spatial and temporal distribution of the vorticity ω , as shown in Fig. 5(c), which is defined as:

$$\omega = \frac{\partial u_y}{\partial x} - \frac{\partial u_x}{\partial y} \quad (11)$$

where u_x and u_y represent the velocity components in the x - and y -directions, respectively. As the cavity develops, the negative-vorticity regions along the air cavity wall exhibit a distinct two-layer structure: an inner high-vorticity zone at the air-liquid interface and an outer low-vorticity zone in the surrounding fluid. This vorticity distribution directly reflects the velocity discontinuity across the interface.

The vorticity contour map in Fig. 5(c) clearly shows that a significant negative vorticity region persists along the air cavity wall throughout the immersion process. This region reveals the presence of a persistent shear layer at the air-liquid interface, generated by the velocity gradient between the two phases at the interface: a strong velocity gradient arises between the high-speed fluid inside the air cavity and the relatively stationary external liquid. As the sphere descends, the negative vorticity region exhibits distinct spatial evolution: it initially forms a concentrated vortex layer near the sphere equator; then, it extends along the cavity wall toward the tail, ultimately forming a continuous vorticity belt enveloping the entire cavity.

This vorticity distribution pattern not only confirms the highly efficient drag-reduction effect of the air cavity interface acting as a slip boundary, but its evolution also directly reflects the dynamic stability of the air cavity morphology. When the negative-vorticity region remains continuous and stable, the air cavity maintains a streamlined configuration. Conversely, discontinuities in the vorticity band or localized alternations between positive and negative vorticity indicate the onset of cavity instability and eventual collapse.

To further characterize the flow field, the Q -criterion for identifying vortex structures is employed, expressed as:

$$Q = \frac{1}{2}(\|\Omega\|^2 - \|S\|^2) \quad (12)$$

where Q is the second invariant of the velocity gradient tensor based on Galilean transformations, Ω is the vorticity tensor, and S is the strain-rate tensor.

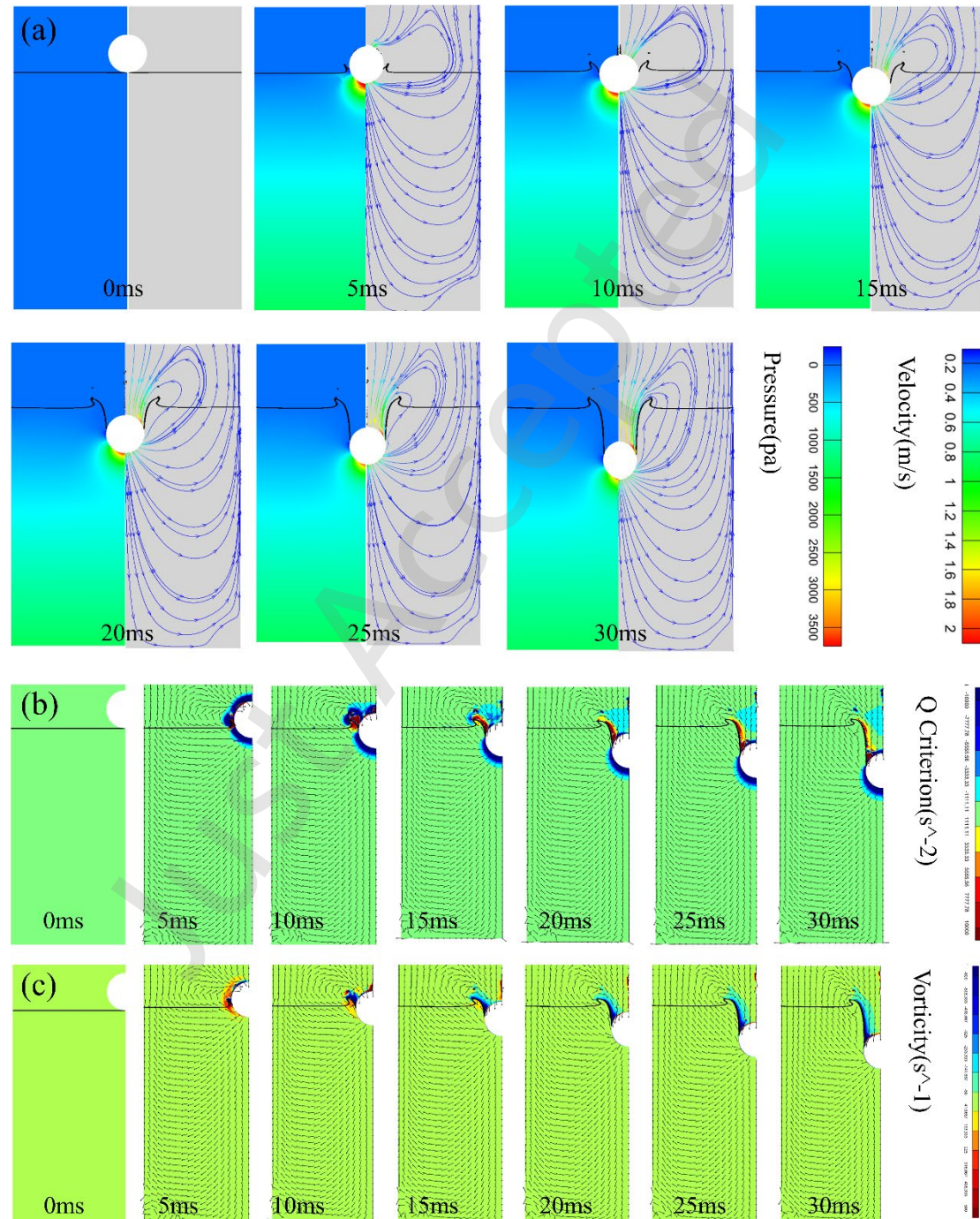


Fig. 5 Numerical simulations reveal the formation mechanism of air cavities within superhydrophobic spheres upon immersion. (a) Pressure distribution and streamline distribution during the vertical immersion of superhydrophobic spheres; (b) Q distribution during the vertical immersion of superhydrophobic spheres; (c) Vorticity distribution during the vertical immersion of superhydrophobic

spheres.

From Eq. (12), it can be seen that this criterion enables precise identification of vortex core regions by quantifying the relative magnitudes of rotational and deformation energies within the flow field. When $Q > 0$, the rotation rate exceeds the deformation rate, indicating that rotational effects dominate. When $Q < 0$, the shear rate exceeds the deformation rate, indicating that shear effects dominate. As shown in Fig. 5(b), during the initial immersion stage, regions with positive and negative Q values alternate along the cavity wall. As the sphere descends and the air cavity expands, these regions with positive and negative Q values gradually diminish. Thus, during the immersion of the SHB sphere, both rotational and shear effects jointly govern the evolution of the air cavity and related flow phenomena.

To further investigate the drag-reduction characteristics of SHB spheres during underwater motion accompanied by air cavity formation, the entry dynamics of both the original and SHB spheres were studied at identical impact velocities (Fig. 6). It can be observed in Fig. 6(a) that the trajectories of the original sphere and the SHB sphere ($D \approx 20$ mm, $U_0 \approx 2.8$ m/s) are nearly indistinguishable during the initial descent stage. As the descent proceeds, however, the divergence between their trajectories becomes increasingly apparent. The SHB sphere with air cavity structures exhibits a significantly greater displacement over time than the original sphere, indicating reduced hydrodynamic drag. Simultaneously, both spheres display identical overall trends in descent velocity. There is no significant difference in the initial descent speed between the two spheres. However, as the falling distance increases, the velocity disparity becomes increasingly pronounced, with the SHB sphere exhibiting a slower rate of velocity decay. The reason why the SHB sphere moves faster than the original sphere is that the original sphere remains fully submerged in water without forming entrained air cavities, resulting in greater drag than that experienced by the SHB sphere. The primary reason for the reduced overall drag on the SHB sphere is the transformation of the no-slip boundary condition at the solid–liquid interface into an effective free-slip boundary condition at the air–liquid interface, thereby reducing shear forces at the boundary [44,45].

Furthermore, the terminal velocities of the original and SHB spheres ($D \approx 10$ mm, $U_0 \approx 2.0$ m/s) were analyzed to gain deeper insight into their motion characteristics (Fig. 6(b)). As the spheres continue to descend, their velocities gradually stabilize, achieving uniform motion. It can be observed from the figure that there is a significant difference

in terminal velocity between the SHB sphere with air cavity entrainment and the original sphere under steady-state conditions. The terminal velocity of the SHB sphere is markedly higher than that of the original sphere, further indicating that the former achieves underwater drag reduction. Theoretical analyses indicate that this increase in terminal velocity stems from the fundamental modification of the flow boundary conditions by the air cavity: the original sphere maintains a solid–liquid no-slip boundary, generating significant surface friction drag, whereas the SHB sphere forms an air–liquid free-slip interface via the air cavity while simultaneously optimizing the tail flow field structure, thereby reducing total hydrodynamic drag.

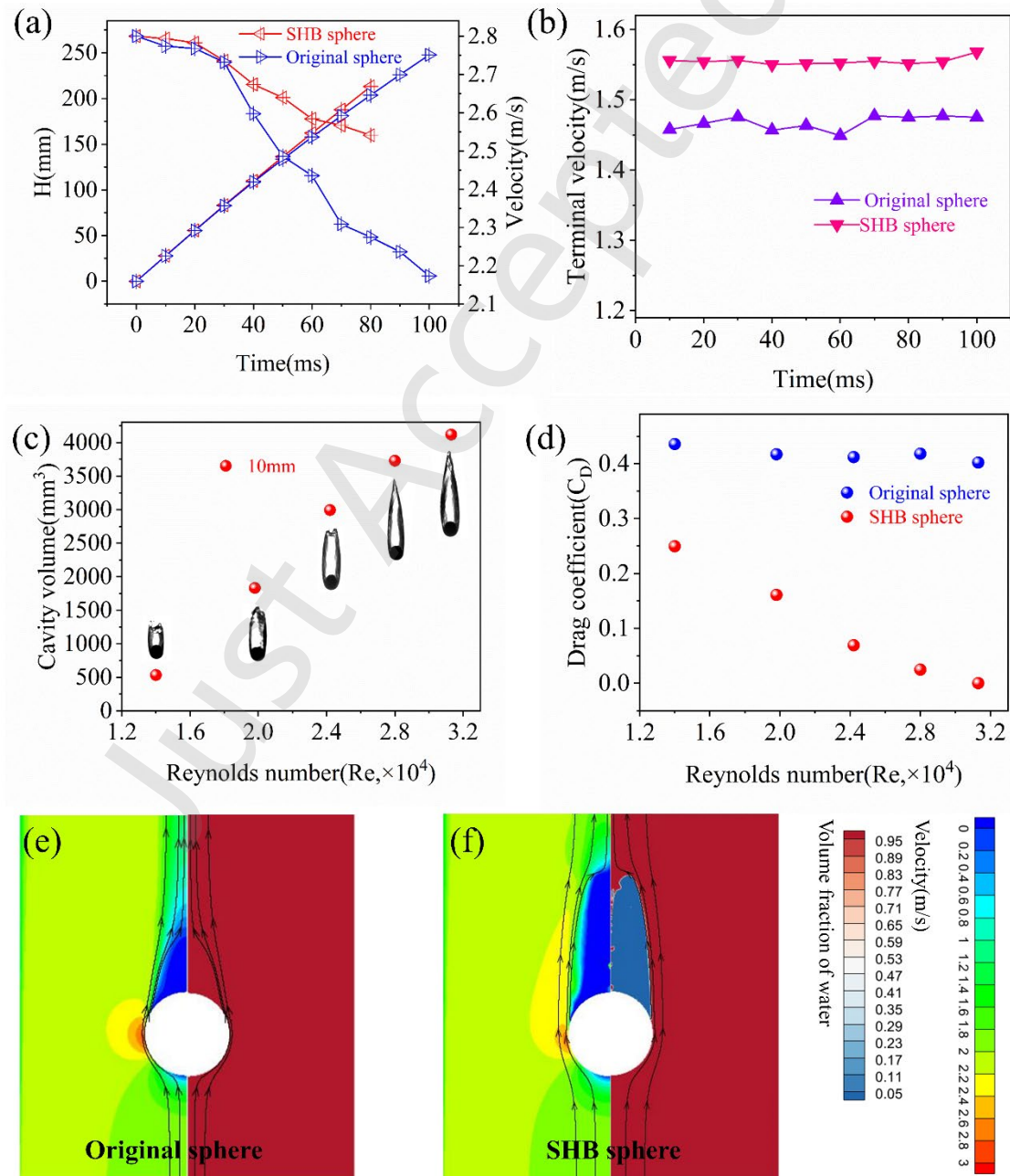


Fig. 6 Analysis of underwater fluid dynamics characteristics and drag reduction mechanisms for the

original sphere and superhydrophobic sphere. (a) Analysis of motion characteristics between the original sphere and superhydrophobic sphere; (b) Comparison of terminal velocities between the superhydrophobic sphere and original sphere; (c) Relationship between cavity volume and Reynolds number; (d) Relationship between sphere drag coefficient and Reynolds number; (e) Flow field around the original sphere; (f) Flow field around the superhydrophobic sphere.

To quantitatively evaluate the drag-reduction effect of SHB spheres during underwater motion, a dimensionless fluid drag coefficient C_D is introduced to characterize the fluid resistance associated with the air cavity structure. The drag coefficient is given by [46,47]:

$$C_D = \frac{8g(m - \rho V_C)}{\rho \pi D_s^2 U_t^2} \quad (13)$$

where g is the gravitational acceleration, m is the sphere mass, ρ is the water density, V_C is the total volume of the spherical cavity structure, D_s is the maximum diameter of the air cavity, and U_t is the terminal velocity of the sphere.

Through quantitative analysis of the steady underwater motion stage, the air cavity volume and the corresponding drag coefficient C_D were determined. The relationship between the cavity volume and the Reynolds number Re is shown in Fig. 6(c). At low Reynolds numbers, the cavity cannot maintain a stable streamlined morphology, resulting in a small cavity volume. With increasing Reynolds number, the cavity gradually transitions from an incomplete, blunt configuration to a stable streamlined structure, accompanied by a significant increase in cavity volume. Since no air cavity forms during immersion of the original sphere, its effective cavity volume corresponds directly to the geometric sphere volume. Fig. 6(d) illustrates the variation in the drag coefficient with the flow condition. The results show that the drag coefficients of the original and SHB spheres exhibit significant differences at different Reynolds numbers. The drag coefficient of the original sphere remains nearly constant, stabilizing at approximately 0.417. However, under identical Reynolds number conditions, this value is significantly higher than that of the SHB sphere, indicating that the presence of underwater air cavities effectively reduces hydrodynamic drag. Additionally, the drag coefficient of the SHB sphere decreased significantly as the Reynolds number increased from 1.4×10^4 to 3.1×10^4 , indicating enhanced drag-reduction performance at higher Reynolds numbers. This trend is closely related to the evolution of the air cavity morphology: as Re increases, the cavity transitions from an initially incomplete and unstable shape into a stable streamlined configuration, thereby significantly reducing the fluid resistance experienced by the SHB sphere.

To further investigate the drag-reduction mechanisms of both the original and SHB spheres underwater, numerical simulations were conducted to analyze the surrounding flow field structures. The drag-reduction mechanism of underwater air cavities can be interpreted from two perspectives: interfacial modification and flow field restructuring. On the one hand, the cavity surface forms a slip boundary, significantly reducing frictional drag. On the other hand, the cavity modifies the effective outer contour of the moving body, thereby modulating the surrounding flow field, vortex structures, and pressure distribution, which effectively diminishes pressure drag. The presence of the air cavity changes the interaction between the sphere surface and the surrounding fluid from solid–liquid contact to air–liquid contact, inducing an apparent slip effect at the air–liquid interface. According to the Navier slip condition, this interfacial slip reduces the velocity gradient within the boundary layer, thereby lowering the wall shear stress and achieving underwater drag reduction.

The flow field distribution around the original sphere without an air cavity is shown in Fig. 6(e). As the water flows over the sphere surface, the combined effects of wall friction and the reverse pressure gradient cause boundary layer separation near the sphere equator, resulting in the formation of a large recirculation zone at the sphere rear. The fluid velocity within this recirculation region is significantly lower than that of the mainstream flow, resulting in a pressure drop that creates a substantial pressure difference between the front and rear of the sphere. This pressure-difference drag substantially increases the total hydrodynamic resistance, ultimately causing the original sphere to exhibit a lower terminal velocity. By contrast, the flow field surrounding the SHB sphere in Fig. 6(f) exhibits distinct characteristics. The trailing edge of the sphere is enveloped by a streamlined air cavity, which effectively improves the wall curvature distribution. The free-slip condition at the air–liquid interface significantly suppresses boundary layer development and mitigates separation phenomena, while also weakening the reverse pressure gradient.

Fig. 7(a) shows significant boundary layer separation near the equatorial region of the original sphere without an air cavity. This phenomenon arises from a sharp increase in the adverse pressure gradient caused by the sudden curvature change on the rear surface of the sphere. Viscous effects cause the near-wall flow velocity to rapidly decay to zero, leading to premature boundary layer separation. The large-scale recirculation zone formed after separation results in a significant pressure drop at the sphere rear, creating a strong pressure contrast with the high-pressure region at the front and thereby

resulting in substantial pressure drag. Conversely, for the SHB sphere, the air cavity attached to the tail effectively extends the streamlined contour, making the wall curvature change more gradual. The pressure distribution results show that the rear pressure of the SHB sphere is higher than that of the original sphere, demonstrating mitigation of the adverse pressure gradient and a downstream shift of the boundary layer separation point. Vortex structure visualization based on the Q -criterion (Fig. 7(c)) further demonstrates that the wake vortices behind the SHB sphere are smaller in scale and more stable in structure, with a significantly reduced recirculation zone. This optimization of the flow field structure directly contributes to the reduction in pressure drag.

A comparison in Fig. 7(b) shows the presence of a high-vorticity boundary layer along the surface of the original sphere, which is indicative of strong shear effects. By contrast, the vorticity is significantly weakened in regions covered by the air film on the SHB sphere, further confirming the reduction in frictional drag.

The mobility of the air-liquid interface also plays a significant role in determining the boundary layer stability. The boundary layer on the SHB surface is thinner, with an extended laminar flow region and a noticeably delayed laminar-to-turbulent flow transition. This behavior is attributed to the air film's capacity to suppress near-wall normal velocity fluctuations and attenuate turbulent bursting events, thereby postponing the transition and reducing the turbulence intensity. Given that the frictional drag associated with a turbulent boundary layer is significantly greater than that of a laminar boundary layer, this stabilization effect directly enhances overall drag reduction. The more regular and smaller dissipative turbulent vortex structures observed in the SHB wake (Fig. 7(c)) also indirectly reflect the reduced turbulence intensity in the boundary layer.

In addition to delaying flow separation, the air cavity also modulates the shedding frequency and spatial arrangement of the wake vortices. Fig. 7(b) indicates that the wake of the original sphere features alternating vortex shedding accompanied by high energy dissipation. By contrast, due to the presence of the air cavity, the wake structure of the SHB sphere appears more diffuse and symmetric, with a lower vortex shedding frequency and reduced pressure fluctuations. This optimization of the wake structure enhances pressure recovery efficiency and reduces base drag. Pressure distribution data also indicate that the pressure recovery process downstream of the SHB sphere is more gradual.

Consequently, the flow field structure behind the SHB sphere becomes more stable, resulting in a marked reduction in the pressure-difference drag. This ultimately enables the sphere to maintain a high terminal velocity.

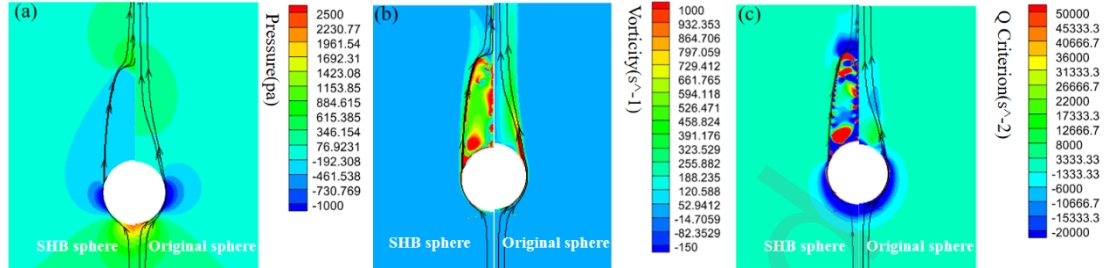


Fig. 7 Analysis of drag reduction mechanisms for the original sphere and superhydrophobic sphere. (a) Pressure distribution; (b) Vorticity distribution; (c) Q distribution.

4 Conclusions

This study systematically investigates the mechanisms of air cavity formation and drag reduction during the vertical immersion of SHB spheres into water through a combination of experimental observations and numerical simulations. The main conclusions are as follows:

(1) Revealed the dynamic critical mechanism by which surface wettability regulates air cavity formation. The results demonstrate that surface wettability is a key control parameter determining whether an air cavity can be generated. Unlike hydrophilic spheres, the liquid film on the SHB sphere surface undergoes early separation driven by an adverse pressure gradient, causing the three-phase contact line to become pinned at the equatorial region of the sphere, thereby fundamentally suppressing the closure of the liquid film toward the sphere apex. Through contact line dynamics analysis and a capillary force-driven model, it is elucidated that the wetting state transition induced by the contact line velocity exceeding the critical capillary number provides the necessary interfacial dynamic conditions for the formation of a stable air cavity.

(2) The air cavity attached to the sphere surface transforms the conventional no-slip boundary condition at the solid–liquid interface into a free-slip condition at the air–liquid interface, changing the surrounding flow field structure. Numerical simulations and Q -criterion vortex identification analyses demonstrate that the air cavity effectively suppresses flow separation in the wake region of the sphere, significantly reduces the spatial extent of the wake vortex zone, and lowers the adverse pressure

gradient, thereby substantially decreasing the pressure difference-dominated drag, achieving a maximum drag-reduction efficiency of approximately 90% under our specific test conditions ($Re \sim 1.4 \times 10^4$ to 3.1×10^4). Notably, the drag-reduction performance is closely related to the cavity morphology: stable, streamlined air cavities formed at higher Reynolds numbers exhibit superior drag-reduction performance, primarily due to further optimization of the flow structure and pressure distribution.

Overall, this study advances the mechanistic understanding of SHB-induced air cavity drag reduction from both interfacial and fluid dynamic perspectives, clarifying the fundamental role of air cavities as the core drag-reducing component. These findings provide a solid theoretical foundation and design guidance for the future development of high-performance drag-reducing surfaces and underwater devices based on stable air cavity slip mechanisms.

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Declaration of competing interest

The authors have no competing interests to declare that are relevant to the content of this article.

Electronic Supplementary Material

Supplementary material is available in the online version of this article

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