

Odd droplets: Fluids with odd viscosity and highly deformable interfaces

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Flows with deformable interfaces are commonly controlled by applying external fields or by modifying the boundaries that interact with the fluid; however, realizing such strategies can be demanding or impractical in many settings. Here, we demonstrate that fluids with broken symmetries can self-control their mechanics. We present a continuum model of a viscous fluid with highly deformable interfaces subject to capillary stresses. Our model features odd viscosity, a parity-violating transport coefficient that emerges in chiral fluids. Using direct numerical simulations of two-dimensional two-phase flows, we focus on the impact of an odd droplet on a superhydrophobic surface. We demonstrate that odd viscosity dramatically disrupts conventional symmetric spreading by inducing asymmetric deformations and chiral flow patterns. Our analysis reveals a variety of unconventional dynamic regimes, including leftward and rightward bouncing as well as rolling, depending on the relative strength of the odd viscosity. Overall, this work presents a framework of broad interest to the control of highly deformable liquids and materials with broken symmetries, as well as to droplet dynamics, where it showcases the potential of odd viscosity as a route for controlling interfacial flows and designing functional metamaterials with tailored fluidic properties.

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I. INTRODUCTION

The principles of symmetry and symmetry breaking lie at the heart of understanding materials and phenomena across scales, both at equilibrium and far from equilibrium. Conventional fluids often exhibit time-reversal symmetry (TRS) at equilibrium. However, fluids like two-dimensional electron fluids in magnetic fields [1–6] or chiral active fluids [7–11] can break TRS, leading to nondissipative, parity-violating odd viscosity [12,13]. This unique property has been associated with various nonintuitive phenomena on different scales, such as Hall viscosity [2,4], pattern formation in turbulent flows [14,15], topological edge modes [16,17], and anisotropic mobility, drag, and lift [18–23]. Many such flows could naturally involve deformable interfaces [9,18,22,24–27], where surface tension and stress boundary conditions give rise to rich, highly nonlinear behaviors. These shape-changing interfacial dynamics are intricately coupled with bulk flow and asymmetric

stresses and are hardly accessible with analytical methods, particularly when deformations are large. Here, we present direct numerical simulations of two-dimensional (2D) flows with odd viscosity using a continuum model to study sharp deformable interfaces with surface tension, revealing unprecedented nonlinear responses and showing how odd viscosity can be harnessed as a parameter to control multiphase flows and design functional metamaterials [28–31].

To this end, we study the rich problem of droplet dynamics and, in particular, droplet impact [32]. The impact of droplets on a substrate or fluid surface has long been a cornerstone of fluid dynamics, with classical studies revealing a wide range of phenomena depending on factors such as impact velocity, surface tension, viscosity, and wettability [32,33]. Traditional results include the spreading and retracting of droplets, crownlike splashing at high impact speeds, and intricate patterns formed by capillary waves. These dynamics are controlled by the balance of inertial, viscous, and surface tension forces, often characterized by dimensionless numbers such as Reynolds, Weber, and Ohnesorge numbers. In this article, we focus on the impact of a 2D droplet on a nonwetting surface [34–37] and show that introducing odd viscosity fundamentally alters these familiar scenarios: The parity-violating stress generates nontrivial flows, which deform the droplet interface. These deformations, in turn, modify the flow field, creating a self-reinforcing interplay between the interface shape and the underlying hydrodynamics. As we will

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discuss, this emergent two-way coupling leads to nonlinear behaviors, such as asymmetric spreading, asymmetric bouncing, and rolling, all of which are inaccessible in conventional fluids.

II. PROBLEM STATEMENT

Consider a 2D droplet of an odd viscous fluid Ω_1 with diameter D_0 and constant density ρ_d . The droplet is initialized such that its center is located at height $y = D_0$ (with $y = 0$ corresponding to the solid wall), with a uniform downward velocity U , and impacts the solid surface at $t = 0$. The surrounding fluid Ω_2 is Newtonian and has a much smaller density and (even) viscosity than the droplet. No gravitational body force is included, such that the subsequent dynamics arise solely from the interplay between inertia, viscosity, odd stresses, and surface tension. Due to the absence of gravity, varying the initial droplet position relative to the wall mainly changes only the fall duration prior to impact, without significantly modifying the impact conditions or subsequent dynamics. We solve an extension of the Navier-Stokes equations for an incompressible fluid (see Appendix A for details on the equations, boundary conditions, and nondimensionalization):

$$\nabla \cdot \mathbf{u} = 0, \quad (1)$$

$$\rho(\mathbf{D}\mathbf{u}/\mathbf{D}t) = -\nabla p + \nabla \cdot \boldsymbol{\tau}^E + \nabla \cdot \boldsymbol{\tau}^O + \mathbf{F}^S, \quad (2)$$

Here, $\mathbf{D}\mathbf{u}/\mathbf{D}t = \partial/\partial t + \mathbf{u} \cdot \nabla$ is the material derivative, and $\mathbf{u}(\mathbf{x}, t)$ and $p(\mathbf{x}, t)$ are the velocity and pressure fields, respectively. The second term on the right-hand side is the divergence of the deviatoric stress tensor $\nabla \cdot \boldsymbol{\tau}^E = \nabla \cdot (\mu^E \dot{\boldsymbol{\gamma}}^E)/\text{Re}^E$ and accounts for the dissipative even viscosity

μ^E in even strain rate field $\dot{\boldsymbol{\gamma}}^E = \begin{bmatrix} 2\frac{\partial u}{\partial x} & \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \\ \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} & 2\frac{\partial v}{\partial y} \end{bmatrix}$ with

the even Reynolds number $\text{Re}^E = \rho_d U D_0 / \mu_d^E$ measuring the inertial effects. The odd effects arise from a force exhibiting the same structure with $\mathbf{F}^O \equiv \nabla \cdot \boldsymbol{\tau}^O = \nabla \cdot (\mu^O \dot{\boldsymbol{\gamma}}^O)/\text{Re}^O$,

where $\dot{\boldsymbol{\gamma}}^O = \begin{bmatrix} -(\frac{\partial u}{\partial y} + \frac{\partial v}{\partial x}) & \frac{\partial u}{\partial x} - \frac{\partial v}{\partial y} \\ \frac{\partial u}{\partial x} - \frac{\partial v}{\partial y} & \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \end{bmatrix}$ is the odd strain rate

field, and $\text{Re}^O = \rho_d U D_0 / \mu_d^O$ is the odd Reynolds number. Nonetheless, the odd stress tensor, importantly, is not dissipative, i.e., $\boldsymbol{\tau}^O : \nabla \mathbf{u} = 0$ [1,13]. This means that the odd stress tensor does not resist the flow and remains “reactive” to it. Within the bulk of the droplet, the odd force can also be rewritten as $\mathbf{F}^O = -(\nabla \omega)/\text{Re}^O$, where $\omega(\mathbf{x}, t)$ is the z component of the vorticity field $\boldsymbol{\omega} = \nabla \times \mathbf{u}$. The last term in the right-hand side of Eq. (2) accounts for surface tension forces, such that $\mathbf{F}^S = (\kappa \delta_s \mathbf{n})/\text{We}$. Here, κ is the curvature of the interface between Ω_1 and Ω_2 , and $\mathbf{n}$ is the normal vector to that interface. The Dirac distribution function δ_s is centered on the interface, expressing the fact that surface tension is concentrated only on this region [38,39]. The Weber number $\text{We} = \rho_d U^2 D_0 / \sigma$ compares the inertial and surface tension forces introduced by the surface tension σ . We use a combination of finite volume and volume of fluid (VOF) methods [39–41], implemented in the open-source code Basilisk [42] to solve these equations. To this end, a scalar color function $c(\mathbf{x}, t)$ is introduced and convected by the droplet phase, Ω_1 . The nondimensional densities and viscosities in Eq. (2)

are calculated in the computational domain as $\rho(c) = c + (1 - c)\rho_a/\rho_d$, $\mu^E(c) = c + (1 - c)\mu_a^E/\mu_d^E$, and $\mu^O(c) = c$. Throughout this paper, we take the properties of the outer fluid as $\rho_a = 10^{-2}\rho_d$ and $\mu_a^E = 10^{-2}\mu_d^E$ (see Appendixes A and B for the details of equations, numerical implementation, and validation).

III. ODD VISCOSITY LEADS TO CHIRAL DROPLETS WITH ANOMALOUS HYDRODYNAMICS

A Newtonian droplet impacting a rigid, nonwetting (superhydrophobic) surface spreads horizontally, dissipating energy through shear flow. Once it reaches maximum diameter (D_{max}), surface tension retracts the droplet, which may bounce fully or partially. Throughout, the droplet remains symmetric with respect to the axis of the fall. The introduction of odd viscosity drastically breaks this established picture. Figure 1(a) shows snapshots of the impact process with parameters $\text{Re}^E = 50$, $\text{Re}^O = 10$, and $\text{We} = 10$ (see Video I in Supplemental Material [43]). At these values, inertia, viscous dissipation, and odd viscous effects all play significant roles. Prior to impact, the drag force from the surrounding fluid induces a small vorticity gradient within the droplet, giving rise to an odd force [18]. However, over the fall time considered in these simulations, these effects remain negligible. The effects of odd viscosity appear immediately upon impact as the droplet spreads over the surface, with local vorticities near the solid boundary perturbing the flow field inside the droplet. This breaks the line symmetry in the flow field, resulting in an asymmetric shape and a leftward shift of the droplet’s center of mass (see Appendixes C and D for a detailed analysis of the transient flow and odd force fields). During the transition from spreading to retraction, the flow anatomy becomes considerably more complex. The interplay of surface deformation and modulated bulk flows caused by odd viscosity produces a time-dependent, heterogeneous flow field within the droplet. Consequently, both the magnitude and direction of odd forces vary in space and time. Simply put, as retraction progresses, shear flow develops in the opposite direction to that during spreading, generating an odd force field that opposes the one generated during impact (see Appendix C for a visualization of the odd force fields). This also initiates a chiral flow within the droplet and rotates it as it rebounds from the surface. In this case, the center-of-mass velocity of the droplet drives a horizontal translation to the left as it ascends. Figure 1(b) shows an extended center-of-mass trajectory, highlighting the counterclockwise rotation of the droplet and off-center bounce.

We further quantify the effect of odd viscosity over time by three key properties [see Fig. 1(c)]: the droplet diameter (D), a shape symmetry parameter (S), and the droplet enstrophy ($e = \int_{\Omega_1} \|\omega\|^2$). Odd droplets attain a smaller maximum diameter than their even counterpart, implying that the odd viscosity partly counteracts the spreading of the droplet by generating flows that dissipate energy [also see Fig. 3(a)]. At later times, further differences are visible as both capillary-driven oscillations and solid-body rotations induced by macroscopic chirality are present. The shape symmetry parameter, which quantifies the skewness of the droplet interface [44], is defined as $S = \frac{1}{\Omega} \int_{\Omega} (\mathbf{x} - \mathbf{x}_{\text{CM}})^3 / [\frac{1}{\Omega} \int_{\Omega} (\mathbf{x} - \mathbf{x}_{\text{CM}})^2]^{3/2}$, where $\mathbf{x}_{\text{CM}}$ is

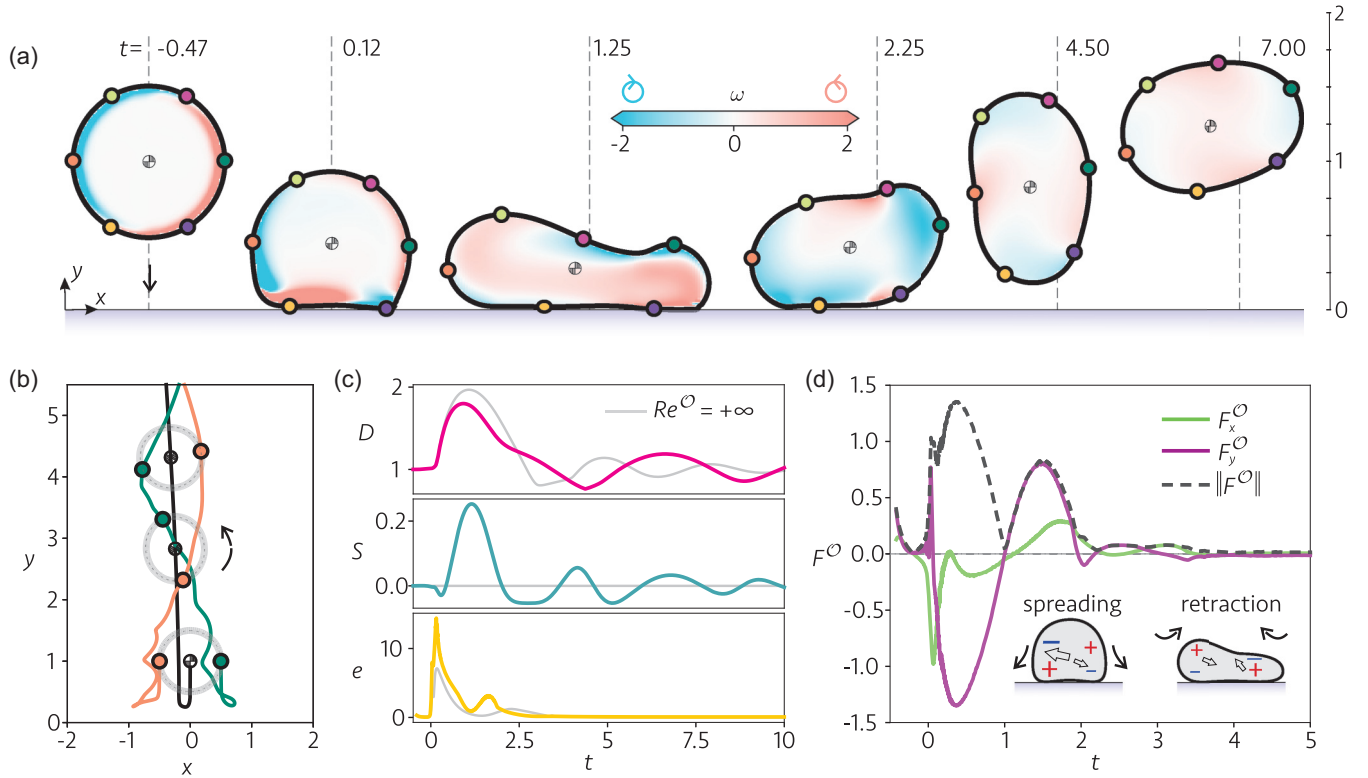


FIG. 1. Anomalous dynamics of an odd droplet. (a) An impacting two-dimensional droplet with odd viscosity spreads asymmetrically on a rigid superhydrophobic surface and bounces back off-center while rotating. The color indicates the magnitude of the vorticity ω . Lagrangian tracer particles placed on the surface of the droplet highlight the surface flows and chirality. (b) The center-of-mass trajectory of the droplet (black line) shows left/right symmetry breaking during spreading and bouncing. The rotation of the droplet is visible by tracking two initially parallel tracer particles on its sides (orange and green lines). (c) Characteristics of odd droplet impact manifest as a smaller spreading diameter D , significant shape asymmetry S , and bursts of entrophy e , particularly upon impact. The gray lines show the results for the even counterpart, where no symmetry breaking is observed. (d) The odd force field $\mathbf{F}^O$ evolves during both spreading and retraction in a nontrivial manner. In simplified terms, the vorticity quadrupole generated during spreading produces a net negative horizontal force. As retraction develops, the flow direction reverses, altering the odd effects (see Appendix C for details of the flow and odd force fields). Note that the force direction is dictated by the sign of the odd viscosity, which is determined by microscopic chirality. At long times, the odd force decays to zero as the flow field becomes nearly purely rotational. See Video I for a comparison of even and odd impacts [43].

the position vector of the center of mass of the droplet. Odd viscosity introduces a strong loss of symmetry in the droplet shape, particularly immediately after impact when shear forces are most pronounced [see Fig. 1(c)]. This is also when the entrophy is maximized. Note that the entrophy is also nonzero for a droplet without odd viscosity, since this quantity measures the magnitude of local vorticity. At later times, surface tension restores the shape symmetry by making the droplet circular again ($S \rightarrow 0$, as $t \rightarrow \infty$), while both the flow's chiral symmetry and the droplet's translational symmetry remain broken, and entrophy decays to nearly zero. To investigate the origins of the left/right symmetry breaking in more detail, we examine the magnitude of the odd force acting inside the droplet, that is, $\|\mathbf{F}^O\| = \frac{4}{\pi D_0^2} \int_{\Omega} \mathbf{F}^O$. This average force, along with its x and y components, is plotted over time in Fig. 1(d). At the early stage of interaction with the wall, negative odd forces dominate, accelerating the droplet leftward on the surface; in contrast, as retraction develops, the odd forces can change sign, guiding the droplet rightward. Note that the values shown in Fig. 1(d) are spatially averaged over a heterogeneous odd force field arising from nonlinear interactions driven by inertia and modulated by odd

viscosity, coupled with large surface deformations (see Appendix C). However, one may provide a simplified, instructive picture during spreading and retraction, where nonsymmetric vorticity quadrupoles are formed [see Figs. 1(a) and 1(d) and Appendix C]. Both left-right and top-bottom symmetries are broken in these quadrupoles, leading to net odd forces with horizontal components, leftward during spreading and rightward during retraction. In the case depicted in Fig. 1, the chiral droplet ultimately bounces leftward. However, discussed in the following, the phenomenological phase map defined by the control parameters is considerably richer, and the interplay between impact, spreading, and retraction can lead to alternative bouncing behaviors or even complete suppression of bouncing, thereby transforming droplets into *rollers*.

Figure 2 shows the effect of odd viscosity, varying $\text{Re}^O \in [1, +\infty]$, while keeping $\text{Re}^E = 50$ and $\text{We} = 10$ constant. Overall, three main dynamical categories emerge. In the limit of weak odd effects (i.e., $\text{Re}^O \rightarrow \infty$), the droplet bounces symmetrically, as expected for a purely even droplet. Remarkably, the direction of left-right symmetry breaking in the trajectories of odd droplets (whether they bounce leftward or rightward) depends on the specific value of Re^O . For large

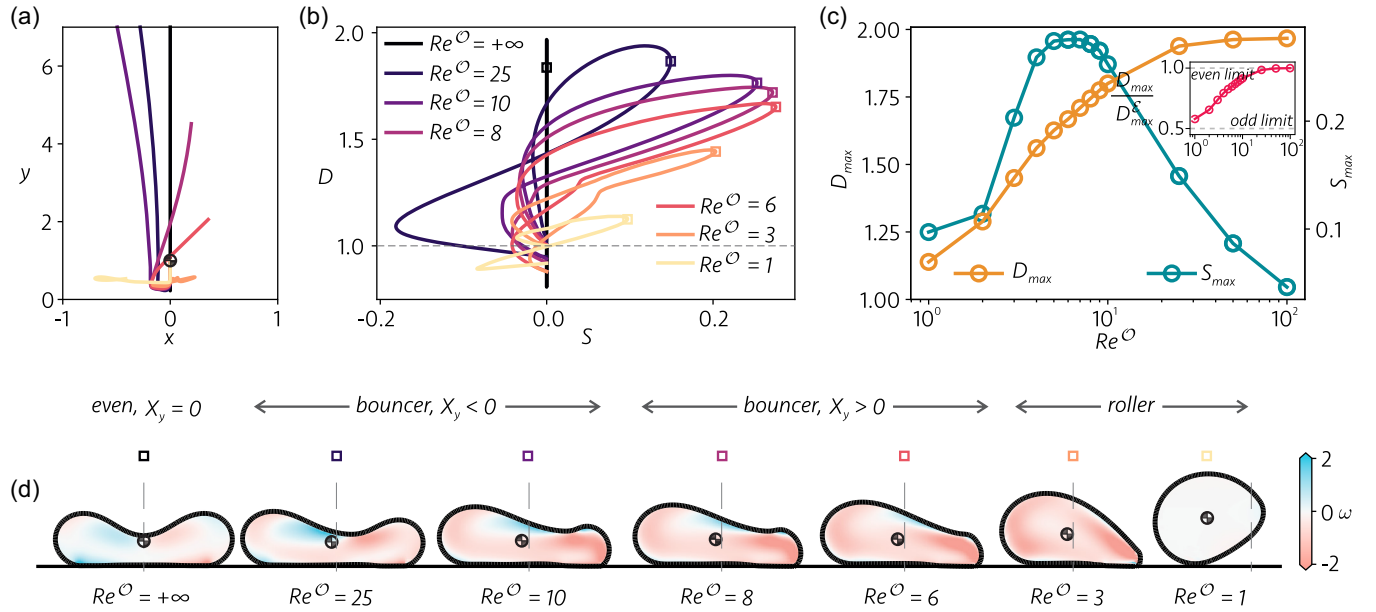


FIG. 2. Controlling droplet dynamics with odd viscosity. (a) Depending on Re^O , an odd droplet bounces leftward, rightward, or rolls on the surface, whereas an even droplet (black line) spreads, retracts, and bounces symmetrically. (b) The spreading diameter and shape symmetry parameter illustrate the influence of Re^O on droplet dynamics. As the droplet spreads and retracts, its shape changes under the combined effects of inertial, capillary, and odd forces. Decreasing Re^O always leads to a smaller maximal diameter but a nonmonotonic maximal shape asymmetry. This is highlighted in panel (c), where the maximal diameter varies between two limits: $D_{\max} \rightarrow 1$ for $Re^O \lesssim 1$ and $D_{\max} \rightarrow D_{\max}^E$ for $Re^O \gg 1$. Nonetheless, the maximal shape asymmetry $S_{\max}$ approaches zero for both large and small Re^O , corresponding to the even and purely chiral limits, respectively. (d) The shape of droplets with $S = S_{\max}$ illustrates the deviation from even-symmetric cases. As Re^O decreases (and the effects of odd viscosity become more pronounced), the droplet becomes less symmetric and bounces to the left. With a further decrease in Re^O , when odd viscosity dominates the flow, the shape asymmetry declines, causing the droplet to bounce to the right and eventually transform into a nearly circular roller. See Video II for a dynamic visualization of all cases shown in this figure [43]. We = 10 for all results in this figure.

values of Re^O (corresponding to small odd viscosity), the droplet's center of mass translates leftward during spreading and continues in that direction after bouncing. When Re^O is further reduced (i.e., odd viscosity increases), the initial leftward translation during spreading persists; however, the retraction phase induces an opposing acceleration that ultimately drives the droplet rightward upon bounce. We quantify the lateral preference of droplets using the quantity $X_y = \partial x_{cm} / \partial y_{cm}$, where x_{cm} and y_{cm} denote the coordinates of the center of mass of the droplet, measured shortly after the bounce. This quantity characterizes the slope of the center-of-mass trajectory after rebound and exhibits a transition from negative to positive values as Re^O varies. A second remarkable phenomenon is the complete suppression of bouncing by odd viscosity. In cases of very high odd viscosity [$Re^O \lesssim \mathcal{O}(1)$], the emergent odd force field results in a strong opposing flow field and effectively “lock” the droplets upon impact. This locking is associated with a rapid conversion of translational kinetic energy into a localized layer of rotational motion near the whole interface of the droplet, which leads to strong energy dissipation (due to even viscosity) and suppresses spreading, retraction, and rebound. Consequently, the chiral droplet rolls over the rigid surface with an internal rotational flow and purely horizontal translational velocity, much like a wheel (see Fig. 10 in Appendix D for an in-depth demonstration of the flow fields). The morphology of these

rollers remains nearly circular with surface perturbations and often initially moves leftward with a low speed.

We further quantify the dynamic morphology of the droplet using the parametric D - S curve shown in Fig. 2(b). Initially, the droplets exhibit a negative S , as they skew to the left during the spreading. Later, during retraction, the skewness becomes positive, with maximum values of S (indicative of the most asymmetric shape) occurring in this phase [see the snapshots of these moments in Fig. 2(d)]. Interestingly, the peak in S nearly coincides with the moment of maximum droplet diameter $D_{\max}$, as the highly deformed droplet at that instant allows for strong asymmetry. Subsequently, surface tension retracts the droplet and promotes circularity, hence weakening these asymmetries ($[D, S] \rightarrow [1, 0]$ as $t \rightarrow \infty$). The values of $D_{\max} = \max_t D(t)$ and $S_{\max} = \max_t S(t)$ are of particular interest. When considering odd viscosity as a design parameter, the maximal spreading and maximal asymmetry appear as the two most important geometrical values. Figure 2(c) shows how these values change with Re^O . Notably, odd viscosity generates overall force fields that suppress the initial spreading of the droplet, always leading to smaller diameters than those of its even counterpart. The maximal droplet spreading has two distinct limits: For negligible odd viscosity ($Re^O \rightarrow \infty$), the spreading reaches the even Newtonian values determined by the balance of surface energy, viscous dissipation, and kinetic energy [45–47]. In this limit, we define $D_{\max}^E$ as the maximal spreading obtained from a sim-

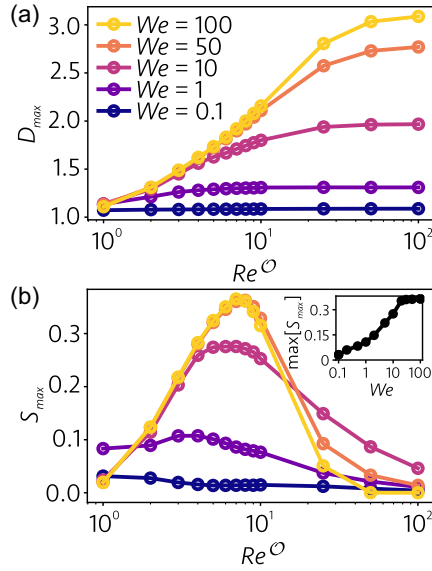


FIG. 3. Morphology dependence on surface tension and odd viscosity. Variation in (a) maximal droplet spreading and (b) maximal shape asymmetry for various values of Re^O and We . The inset in panel (b) displays the variation of the peak in $S_{\max}$ as a function of We .

ulation performed with vanishing odd viscosity ($Re^O \rightarrow \infty$), while keeping all other parameters unchanged. For highly odd droplets ($Re^O \lesssim 1$), the locking phenomena described above lead to the limit $D_{\max} \rightarrow 1$; i.e., the droplet remains circular. These two limits are made explicit in the inset within Fig. 2(c), where the vertical axis is rescaled with $D_{\max}^E$.

The shape symmetry parameter S , however, exhibits a nonmonotonic trend. As odd viscosity increases to moderate values, the droplet becomes more asymmetric because the velocity field is significantly altered yet not fully converted into rotation. For more extreme odd viscosity values, nearly the entire velocity field transitions to solid-body rotation so that the droplet experiences little deformation and retains a nearly circular shape with minimal asymmetry (rollers).

The variations in D and S also depend on the Weber number We , which here serves as a measure of surface tension strength (see Fig. 3). For small values of We , odd viscosity exerts only a minimal effect on these observables, as the dominant surface tension suppresses spreading and limits deformation, thereby preventing significant asymmetry. As We increases, inertial effects become more pronounced relative to capillarity, allowing the droplet to deform sufficiently for odd viscosity to influence its dynamics. Notably, $D_{\max}$ always increases monotonically with Re^O across all We values, underscoring the role of odd viscosity in suppressing droplet spreading. A plateau is always reached for high Re^O , corresponding to the even limit. In contrast, the shape symmetry parameter S exhibits a nonmonotonic trend with odd viscosity for intermediate and high We values. A peak in asymmetry is observed at intermediate Re^O , where odd viscosity is sufficiently large to disrupt conventional symmetry but not so high as to nearly eliminate droplet deformation. At large We values, this maximum asymmetry reaches a plateau as

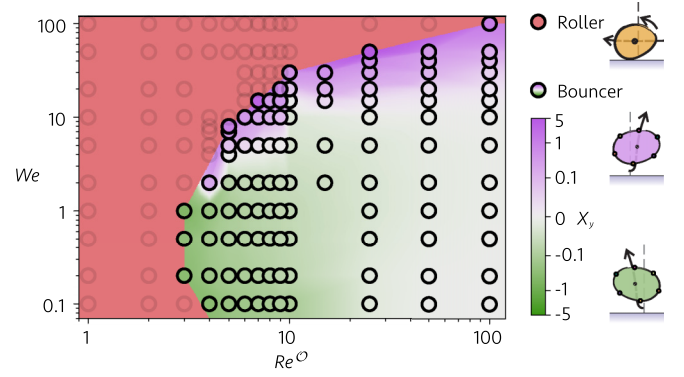


FIG. 4. Regime map of odd droplets. Left and right bouncers, as well as rollers, are classified based on the balance among inertial, viscous, capillary, and odd forces. The left/right behavior of the bouncers is determined by the bouncing angle X_y , measured immediately after the droplet leaves the surface.

surface tension effects disappear for $We \gg 1$ [see the inset in Fig. 3(b)].

A complete theoretical description of the dependence of the three odd regimes observed above, namely, left bouncers ($X_y < 0$), right bouncers ($X_y > 0$), and rollers on the control parameters (Re^O , We , and Re^E) is not trivial, particularly due to the highly nonlinear coupling among large interface deformations, inertia, and odd viscous effects and requires further investigations. In Appendixes C and D, we further inspect the internal droplet vorticity field and its corresponding odd force field F^O during the different phases of impact. Here, we proceed to perform a large-scale parametric study to identify the regime maps corresponding to these behaviors.

IV. SUMMARY AND PERSPECTIVE: ODD VISCOSITY TO CONTROL INTERFACIAL DYNAMICS AND DESIGN METAFUIDS

The results presented above suggest that odd viscosity can be used as a tool to control the shape and dynamics of droplets. In Fig. 4, we summarize the simulation results in a We - Re^O phase map. Three regimes are distinguishable: At low Re^O values, odd viscosity dominates, locks the chiral droplets, suppresses bouncing, and, regardless of We values, turns the droplets into rollers. At moderate Re^O values, all regimes may be observed, depending on We . When We is small, the odd force field during spreading dominates, causing the droplet to bounce leftward $X_y < 0$; but, as We increases, the force field emerged from a highly flattened droplet and its retraction becomes significant, and the droplet moves rightward. Surprisingly, at high We values, large surface deformations result in pronounced odd viscosity effects and the dominance of chiral flows, leading to the reappearance of rollers. As Re^O increases further, all regimes persist, albeit more weakly, as even behavior is asymptotically recovered.

Our study employs direct numerical simulations of the incompressible odd Navier-Stokes equations while using a volume-of-fluid framework to explore the highly nonlinear dynamics of droplets with deformable interfaces under cap-

illary stresses. The present approach, in principle, can be applied to a broad class of interfacial flows with deformable boundaries. In contrast to reduced descriptions based on geometrical constraints, such as lubrication or shallow-water theories (e.g., Refs. [25,48,49]), the simulations fully resolve the interface dynamics and stress fields without assuming small slopes or weak deformations. This provides access to strongly nonlinear regimes that are generally difficult to treat analytically (e.g., in Refs. [18,50]). This capability not only enables detailed comparisons with experimental observations but also offers high-resolution insight into quantities such as strain rates and stress fields, which are difficult to measure in laboratory settings.

In a broader context, we show that introducing odd viscosity offers a rich phenomenology by drastically altering the dynamics of flows with deformable interfaces, as exemplified by the dynamics of droplets. Although odd viscosity is nondissipative, it disrupts the conventional symmetric spreading observed in Newtonian fluids by modifying internal flow patterns and inducing asymmetric interface deformations, which in turn further alter the flow and the influence of odd viscosity. This nonlinear self-reinforcing loop eventually results in droplets self-control upon contact with a boundary. This control over droplet behavior, whether in the form of reduced maximal spreading or leftward/rightward bouncing or rolling motion, highlights the potential of odd viscous effects as a powerful tool for controlling multiphase flows. This ability may have promising industrial applications in the future. Consider, for example, inkjet printing, where suppressing unwanted bouncing and precisely directing droplets is of great interest. However, current droplet control methods rely on external actuation, such as temperature and electric fields, or on modifying boundary conditions like using asymmetrically structured surfaces [51–58]. Controlling droplets, and more generally interfacial flows, using odd viscosity may overcome the limitations associated with these techniques by enabling the fluid to self-regulate based on the governing physical parameters. This, nonetheless, requires the design of chiral fluids in driven or active systems. Experimental progress toward such systems has been achieved, in particular, in suspensions of driven rotors, where effective hydrodynamic descriptions incorporating odd viscosity have been used to interpret measurements of chiral wave damping and surface dynamics [9]. In these systems, an effective odd viscosity can be inferred from spectral measurements, although it remains orders of magnitude smaller than the even viscosity, with $\mu^O/\mu^E \sim 10^{-9}$ in typical realizations. In contrast, the regimes explored in our work involve comparatively strong odd-viscous effects, including cases such as the rolling droplet in Fig. 2(d), where $\mu^O/\mu^E = \text{Re}^E/\text{Re}^O = 50$. This separation highlights that the parameter ranges required to observe pronounced odd-viscous effects at the droplet scale remain far from current experimental realizations. While active and chiral fluids provide a route toward generating odd stresses, their extension to controlled macroscopic interfacial droplet dynamics remains an open challenge requiring further development. Looking even further into the future, the development of metafluidic materials that leverage activity and odd properties to interact with complex environments presents another exciting frontier. To this end, it is possible to complement the recent advances

in adaptive active solids [59], while exploring opportunities beyond the constraints of elasticity.

There are several immediate opportunities to extend the present work. First, developing minimal models of interfacial fluids that incorporate odd viscosity and large deformations could provide further theoretical insights into these complex behaviors. As illustrated by the odd mass-spring (OMS) model discussed in Appendix E, simple phenomenological descriptions can capture some qualitative features of odd droplet dynamics, but remain unable to accurately predict trajectories, shape asymmetries, and transition boundaries between different regimes. Developing reduced models that overcome these limitations represents a promising direction for future research. Second, the present study is restricted to two-dimensional flows. Extending this framework to fully three-dimensional simulations constitutes a natural next step, where additional rotational and translational degrees of freedom become available, potentially enabling richer modes of control and more complex odd-viscous interfacial dynamics. Third, integrating odd viscosity with other mechanical properties, particularly elasticity [59–61], may not only lead to tools for optimizing the design of functional metamaterials but also deepen our understanding of their nonlinear behavior. In fact, the intriguing similarities and differences between odd viscous fluids and odd elastic bodies with deformable interfaces [59] merit further investigation. Fourth, exploring alternative forms of odd viscous terms arising from microscopic constituents could enable tailored solutions for specific fluidic applications [62]. Finally, the implications of our results may extend beyond classical fluid mechanics and present modeling opportunities in other condensed matter problems, such as electronic fluids, where interactions with boundaries play a crucial role in system dynamics [6,16,17].

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DATA AVAILABILITY

The code used in this study is openly available in the GitHub repository [63]. All other data that support the findings of this study are available from the authors upon reasonable request.

APPENDIX A: EQUATIONS AND BOUNDARY CONDITIONS

To properly state the equations of motion used in this work, it is useful to begin by introducing the scheme used to track the droplet interface. The interface between the two fluids is tracked by the VOF scheme [40]. In this method, a tracer

$c \in [0, 1]$ stores the fraction of droplet fluid contained in each cell of the mesh. The properties ρ , $\mu^\mathcal{E}$, and $\mu^\mathcal{O}$ are then interpolated locally as

$$\rho(c) = c \rho_d + (1 - c) \rho_a, \quad (\text{A1})$$

$$\mu^\mathcal{E}(c) = c \mu_d^\mathcal{E} + (1 - c) \mu_a^\mathcal{E}, \quad (\text{A2})$$

$$\mu^\mathcal{O}(c) = c \mu_d^\mathcal{O}, \quad (\text{A3})$$

where the subscripts d and a indicate, respectively, the properties of the droplet (Ω_1) and the outer fluid (Ω_2). The volume fraction field is advected over time by solving the equation

$$\frac{\partial c}{\partial t} + \nabla \cdot (c \mathbf{u}) = 0. \quad (\text{A4})$$

The flow inside and outside the droplet is modeled by the two-phase incompressible Navier-Stokes equations, given by

$$\nabla \cdot \mathbf{u} = 0, \quad (\text{A5})$$

$$\rho \left(\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} \right) = -\nabla p + \nabla \cdot \boldsymbol{\tau}^\mathcal{E} + \nabla \cdot \boldsymbol{\tau}^\mathcal{O} + \mathbf{F}^\mathcal{S}, \quad (\text{A6})$$

where $\mathbf{u}$ and p are the velocity and pressure fields, respectively, and ρ is the fluid mass density. In our numerical method, the surface tension force is defined as a body force $\mathbf{F}^\mathcal{S} = \sigma \kappa \delta_s \mathbf{n} \approx \sigma \kappa \nabla c$, where κ is the curvature of the interface, σ the constant surface tension coefficient, $\mathbf{n}$ is the unit vector normal to the interface, and δ_s is the Dirac distribution function centered on the interface. This function expresses the fact that surface tension is concentrated only on the interface [38,39]. The deviatoric stress tensor of the fluid is $\boldsymbol{\tau} = \boldsymbol{\tau}^\mathcal{E} + \boldsymbol{\tau}^\mathcal{O}$, where the even and odd components of stress are given by

$$\boldsymbol{\tau}^\mathcal{E} = \mu^\mathcal{E} \dot{\boldsymbol{\gamma}}^\mathcal{E} = \mu^\mathcal{E} \begin{bmatrix} 2 \frac{\partial u}{\partial x} & \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \\ \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} & 2 \frac{\partial v}{\partial y} \end{bmatrix}, \quad (\text{A7})$$

$$\boldsymbol{\tau}^\mathcal{O} = \mu^\mathcal{O} \dot{\boldsymbol{\gamma}}^\mathcal{O} = \mu^\mathcal{O} \begin{bmatrix} -\left(\frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \right) & \frac{\partial u}{\partial x} - \frac{\partial v}{\partial y} \\ \frac{\partial u}{\partial x} - \frac{\partial v}{\partial y} & \frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} \end{bmatrix}, \quad (\text{A8})$$

and $\mu^\mathcal{E}$ and $\mu^\mathcal{O}$ are, respectively, the even and odd viscosity functions given in Eqs. (A2) and (A3). The structure of $\boldsymbol{\tau}^\mathcal{O}$ reflects the parity-breaking nature of odd viscosity: It is obtained from the standard strain-rate tensor ($\dot{\boldsymbol{\gamma}}^\mathcal{E}$) by a 90° rotation in the plane, leading to a stress response that is symmetric in indices but orthogonal to the local rate of deformation. This form ensures that the resulting stress is nondissipative and couples gradients of vorticity to the momentum balance.

On the rigid surface, no-slip and impermeability boundary conditions are imposed on the velocity field ($\mathbf{u} = \mathbf{0}$), while a nonwetting condition ($c = 0$) is imposed for the volume fraction field. Outflow conditions are applied for the other three boundaries of the domain.

We nondimensionalize Eqs. (A5) and (A6) by defining rescaled variables as

$$\mathbf{x} = D_0 \bar{\mathbf{x}}, \quad t = \frac{D_0}{U} \bar{t}, \quad \mathbf{u} = U \bar{\mathbf{u}}, \quad p = \rho_d U^2 \bar{p}, \quad \kappa = \frac{1}{D_0} \bar{\kappa}, \quad \delta_s = \frac{1}{D_0} \bar{\delta}_s, \quad (\text{A9})$$

where D_0 is the initial droplet diameter and U is the impact velocity with the solid boundary.

The dimensionless version of Eqs. (A5) and (A6) will then become

$$\nabla \cdot \bar{\mathbf{u}} = 0, \quad (\text{A10})$$

$$\bar{\rho} \left(\frac{\partial \bar{\mathbf{u}}}{\partial \bar{t}} + (\bar{\mathbf{u}} \cdot \nabla) \bar{\mathbf{u}} \right) = -\nabla \bar{p} + \frac{1}{\text{Re}^\mathcal{E}} \nabla \cdot (\bar{\mu}^\mathcal{E} \bar{\boldsymbol{\gamma}}^\mathcal{E}) + \frac{1}{\text{Re}^\mathcal{O}} \nabla \cdot (\bar{\mu}^\mathcal{O} \bar{\boldsymbol{\gamma}}^\mathcal{O}) + \frac{1}{\text{We}} \bar{\kappa} \bar{\delta}_s \mathbf{n}, \quad (\text{A11})$$

where the auxiliary variables $\bar{\rho}$, $\bar{\mu}^\mathcal{E}$, $\bar{\mu}^\mathcal{O}$ are the dimensionless versions of the VOF interpolating functions in Eqs. (A1) and (A3), given by

$$\bar{\rho}(c) = c + (1 - c) \rho_a / \rho_d, \quad (\text{A12})$$

$$\bar{\mu}^\mathcal{E}(c) = c + (1 - c) \mu_a^\mathcal{E} / \mu_d^\mathcal{E}, \quad (\text{A13})$$

$$\bar{\mu}^\mathcal{O}(c) = c. \quad (\text{A14})$$

Throughout this paper, we take the properties of the outer fluid as $\rho_a = 10^{-2} \rho_d$ and $\mu_a^\mathcal{E} = 10^{-2} \mu_d^\mathcal{E}$, such that our system is determined only by three parameters: the even Reynolds, odd Reynolds, and Weber numbers. These numbers are defined, respectively, by

$$\text{Re}^\mathcal{E} = \frac{\rho_d U D_0}{\mu_d^\mathcal{E}}, \quad \text{Re}^\mathcal{O} = \frac{\rho_d U D_0}{\mu_d^\mathcal{O}}, \quad \text{We} = \frac{\rho_d U^2 D_0}{\sigma}. \quad (\text{A15})$$

In all other sections of this paper, only the nondimensional variables in Eq. (A9) are used, and for convenience, we omit the upper bars.

APPENDIX B: NUMERICAL IMPLEMENTATION

We used the open-source free software language Basilisk C to solve the equations of motion (A10) and (A11). An overview of the numerical procedure will be given in this section, while detailed descriptions of Basilisk C can be found in Ref. [42]. The numerical methods used have been tested for various Newtonian and non-Newtonian fluids with deformable interfaces, including many droplet dynamics scenarios [37,39,64–68].

The simulation is setup by creating a square domain $[-4, 4] \times [0, 8]$ where a droplet of diameter 1 is placed with its center at $[0, 1]$. The velocity within the droplet region (Ω_1) is initialized as $\mathbf{u}_0 = [0, -1]$. To accurately resolve the flow structure inside the droplet and its shape, we apply increased refinement levels for the liquid phase and around it. The mesh is dynamically adapted as the droplet deforms over time. The additional refinement follows the interface and areas around it where the flow needs to be accurately resolved [41].

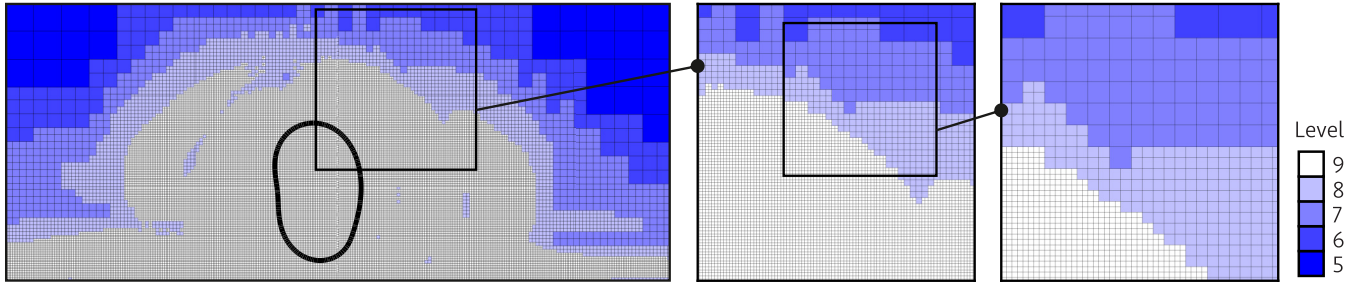


FIG. 5. Example of mesh with adaptive refinement. The mesh is based on a quadtree structure and the level of refinement is higher closer to the droplet, adapting over time as the droplet deforms and moves.

For example, the mesh corresponding to the simulation from Fig. 1 is shown in Fig. 5 for a specific time instant. This figure illustrates how the mesh is created following a quadtree structure with a maximum level of refinement 9. Considering the total size of the square domain (8), our smallest cells have size $\Delta_x = 8/2^9 = 0.015625$, such that the droplet diameter is discretized by 64 cells.

The standard Navier-Stokes solver in the Basilisk C language solves the governing equations using a projection method and a multilevel Poisson solver to treat the even viscosity term $\nabla \cdot (\mu^\varepsilon \dot{\gamma}^\varepsilon)$ (see Refs. [39,41] for more details). In our implementation, we modify this standard code in order to add the new terms associated to the odd viscous stress $\nabla \cdot (\mu^\mathcal{O} \dot{\gamma}^\mathcal{O})$.

We validate our implementation with a simple single-phase scenario that allows for an analytical solution. A two-dimensional channel is created with half-width H and length L . The parabolic Poiseuille velocity profile is imposed on the left side as a boundary condition. We nondimensionalize our equations by rescaling variables in the same way mentioned in Eq. (A9), but using H as length scale and the maximum velocity of the Poiseuille profile $U_{\max}$ as velocity scale. In this manner, we obtain the same nondimensional equations (A10) and (A11), but for a single fluid ($\bar{\rho} = \bar{\mu}^\varepsilon = \bar{\mu}^\mathcal{O} = 1$) and with no surface tension ($We \rightarrow +\infty$). The solution will then depend only on the two Reynolds numbers Re^ε and $Re^\mathcal{O}$ and the boundary conditions.

Note that we do not impose the pressure gradient as a condition, such that this gradient will be calculated by the solver from the imposed velocity field. It is, however, necessary to impose at least one boundary condition for the pressure to guarantee a unique solution. With this in mind, we impose $p = 0$ at the right wall. While this condition is not consistent at the boundary with the analytical solution presented below, we will show that further from this boundary, the solver can still recover the correct analytical solution expected for a fully developed Poiseuille profile.

In this scenario, Eqs. (A10) and (A11) can be easily solved analytically to obtain the pressure and its gradient within the domain, which are given by

$$p(x, y) = \frac{2}{Re^\varepsilon} \left(\frac{L}{H} - x \right) - \frac{2}{Re^\mathcal{O}} y, \quad (B1)$$

$$\nabla p = \left[-\frac{2}{Re^\varepsilon}, \frac{2}{Re^\mathcal{O}} \right]. \quad (B2)$$

Note that, for $Re^\mathcal{O} \rightarrow \infty$, we recover the classical pressure from a Poiseuille flow, which presents simply a constant gradient in the x direction. With the presence of odd viscosity, a constant gradient also appears in the y direction. This classical solution has also been shown in Ref. [13] and was similarly used as numerical validation in Ref. [22].

Figure 6(b) shows the pressure field within the computational domain for an even and an odd fluid. In all cases, the even Reynolds number is kept fixed as $Re^\varepsilon = 0.01$, so that no inertial effects are expected. From these fields, we can see that odd viscosity introduces a pressure gradient in the y direction as it rotates the pressure isolines from vertical to diagonal. We quantified the x and y components of this pressure gradient by taking a horizontal cross section at the center $y = 0$ and a vertical at $x = L/(2H)$. This gradient is plotted for six values of $Re^\mathcal{O}$ in Fig. 6(a), where we also show the analytical solution from Eq. (B2). We see that our solution quantitatively agrees with the expected values.

APPENDIX C: ODD FORCE FIELDS DURING SPREADING AND RETRACTION

To elucidate the time-dependent, heterogeneous odd force field during the impact, spreading, and retraction phases, we present in Fig. 7 the fields for vorticity and both components of $\mathbf{F}^\mathcal{O}$ at various time instants. Panel (a) displays the same example as in Fig. 1. We observe that near the moment of impact, the horizontal odd force $F_x^\mathcal{O}$ exhibits localized positive regions close to the solid surface due to steep vorticity gradients, while it remains negative throughout the bulk of the droplet. This explains the characteristic leftward translation during impact. As the droplet spreads and deforms, positive values of $F_x^\mathcal{O}$ begin to emerge within the bulk and along the upper surface. This shift leads to a gradual rightward acceleration that persists during the retraction. In this specific case, however, the spreading and retraction periods are too short for the rightward acceleration to completely counteract the initial leftward velocity, and the droplet ultimately bounces while still translating to the left.

Panel (b) of Fig. 7 presents a second case in which the droplet bounces and translates to the right. Near the moment of impact, the behavior is almost identical to that in panel (a). This brief timescale prevents significant viscous (both odd and even) and capillary effects from developing, so the droplet initially exhibits its characteristic odd force field with a predominantly negative $F_x^\mathcal{O}$ component. However, due to higher values of We and $Re^\mathcal{O}$, the droplet undergoes prolonged

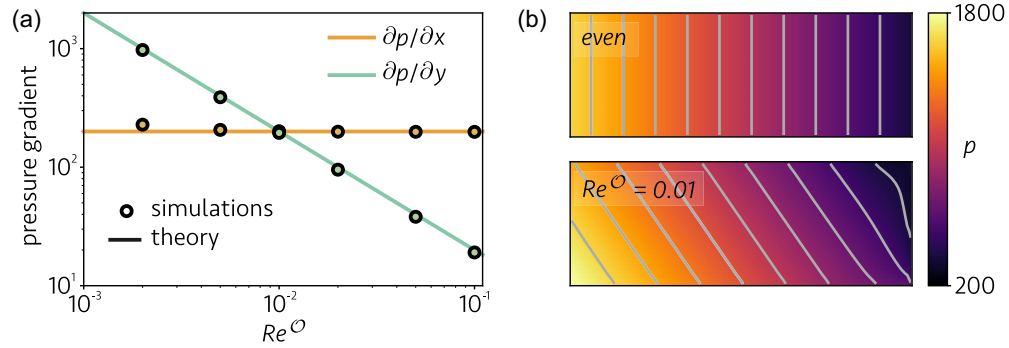


FIG. 6. Pressure field and its gradients in a Poiseuille flow for a fluid with different odd viscosity. (a) Comparison between numerical (circles) and analytical (solid lines) values for the pressure gradient as a function of Re^O . (b) The full pressure field in the channel is shown for an even and an odd fluid with $Re^O = 0.01$. The isolines of the pressure field are shown in gray.

spreading and retraction phases. During these extended stages, the emergence of positive F_x^O values eventually reverses the droplet's translation, leading to rightward motion at the moment of bouncing.

APPENDIX D: VELOCITY FIELD DISTRIBUTIONS FOR CHARACTERISTIC SCENARIOS

A closer inspection of the dynamics and internal flow fields of the odd droplet is provided for four characteristic scenarios: (1) leftward bouncing, (2) rightward bouncing, (3) surface rolling, and (4) an extreme case of rolling, where the droplet slides and comes to a stop. In Fig. 8, we inspect the

velocity field for an impact resulting in a droplet that bounces to the left. This case corresponds to the parameters $Re^E = 50$, $Re^O = 10$, and $We = 10$ [same as in Figs. 1 and 7(a)]. Immediately after impact, the highly inertial flow creates a strong localized shear that generates an odd force, pushing the droplet to the left. This is evidenced by the sharp negative peak in F_x^O shown in Fig. 8(a), which consistently appears shortly after impact. As the droplet spreads, the peak in F_x^O softens as the velocity field changes and its magnitude decreases. The complex interplay of inertial spreading, capillary retraction, and viscous dissipation eventually reverses the sign of the average F_x^O , leading to an acceleration of the droplet to the right that reduces the initial leftward velocity obtained during

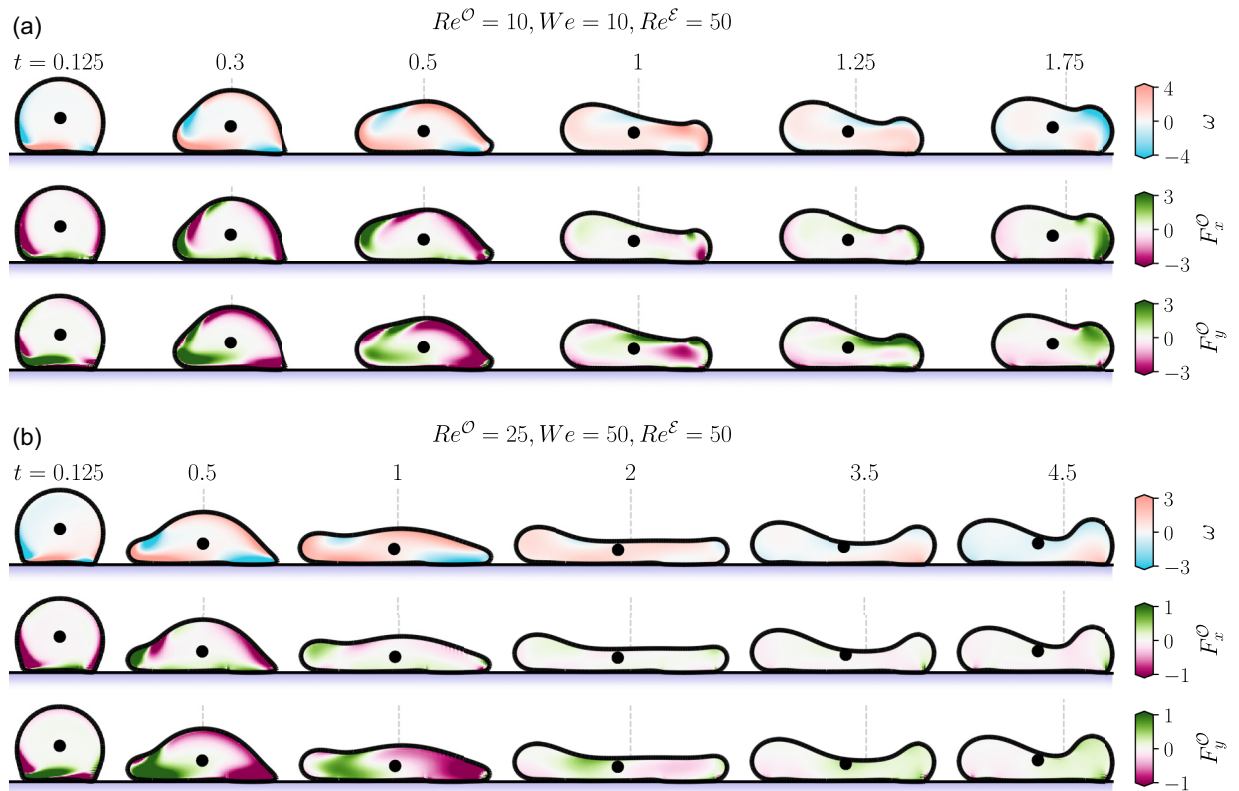


FIG. 7. Vorticity and odd force fields during spreading and retraction. (a) An example of a left bouncer ($Re^E = 50$, $Re^O = 10$, $We = 10$). (b) An example of a right bouncer ($Re^E = 50$, $Re^O = 25$, $We = 50$).

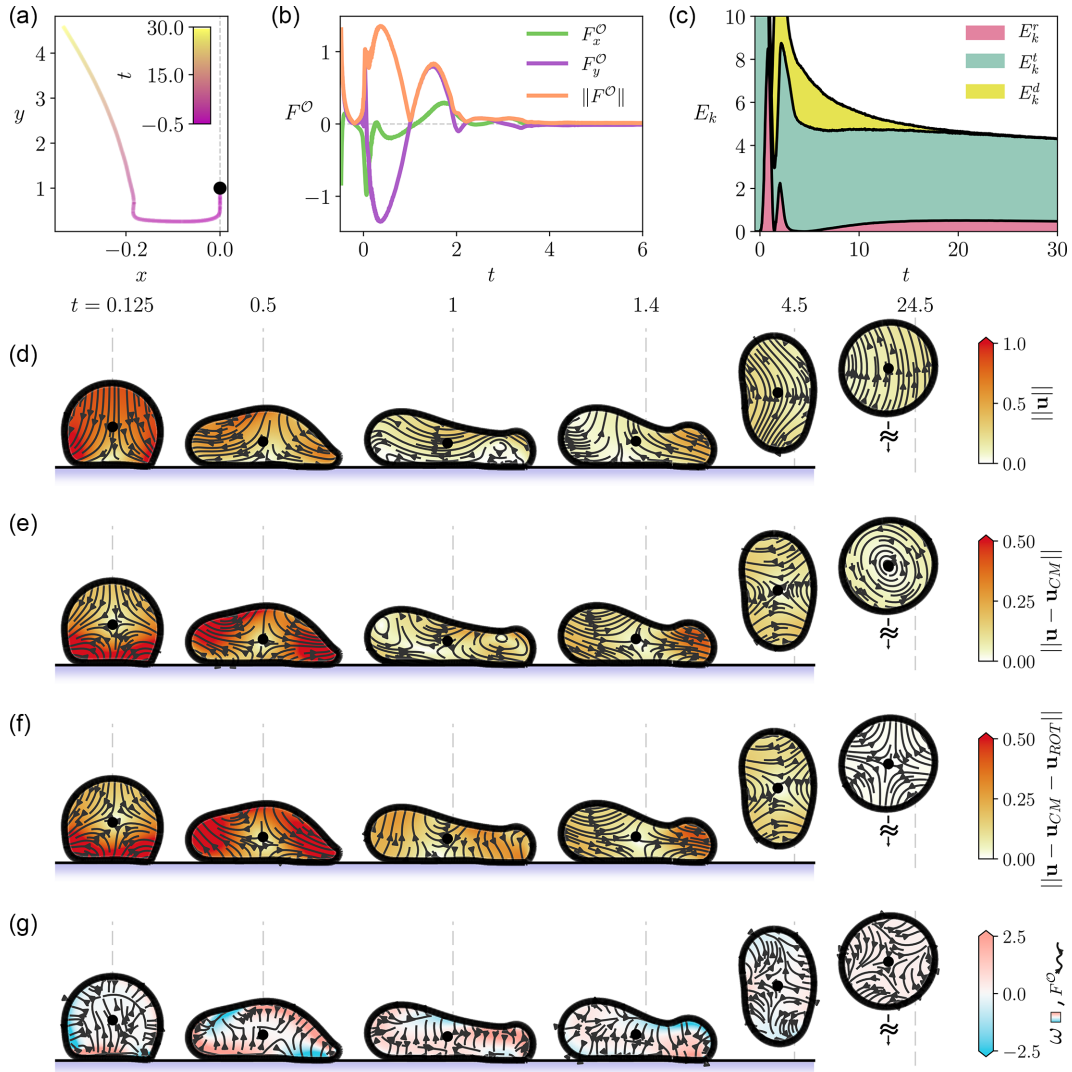


FIG. 8. Dynamics of a left-bouncing odd droplet ($\text{Re}^E = 50$, $\text{Re}^O = 10$, $\text{We} = 10$). (a) Center-of-mass trajectory, marked by time. (b) Components of the average odd force over time. (c) Kinetic energy budget over time, where superscripts r , t , and d denote rotation, translation, and deformation, respectively. The colored regions are stacked cumulatively, such that the thickness of each band corresponds to the contribution of the associated component. At late times, nearly the entire velocity field consists of translation and solid-body rotation. (d) (Lab-frame) velocity fields at selected time instants. (e) Velocity fields with the pure translation component removed, revealing the late-time solid-body rotation induced by odd viscosity. (f) Deformation component of the velocity field. (g) Vorticity fields (background colors) and the odd force (streamlines). Inside the droplet, the odd force follows the negative gradient of the vorticity ($F^O = -(\nabla \omega)/\text{Re}^O$). See Video III, for a dynamic visualization of this simulation [43].

impact. However, this acceleration is not strong enough to completely reverse the droplet's motion, so it moves left after bouncing. As the droplet translates upward, it experiences shape oscillations driven by surface tension acting on its non-spherical form. Since no gravitational body force is included in the simulations, the droplet is not pulled back toward the wall after rebound and can therefore continue moving upward while slowly dissipating energy through interaction with the low-density outer fluid. The velocity field generated by these oscillations and the drag produces a small odd force that, over time, modifies the droplet's internal velocity field.

At late times, as shown in panel (d), the velocity field of the bouncing droplet appears to be almost pure upward translation. However, after subtracting the center-of-mass velocity, $\mathbf{u}_{\text{CM}}$, the residual velocity field reveals an almost perfect solid-

body rotation, indicating that the droplet is both translating and rotating. By measuring the average vorticity ω_{avg} within the droplet and subtracting the solid-body rotation field,

$$\mathbf{u}_{\text{ROT}} = \frac{1}{2} \omega_{\text{avg}} [-(y - y_{\text{CM}}), x - x_{\text{CM}}], \quad (\text{D1})$$

we obtain the final residual velocity field in the droplet, which still has the potential to induce deformation. We therefore decompose the velocity field as

$$\mathbf{u} = \mathbf{u}_{\text{CM}} + \mathbf{u}_{\text{ROT}} + \mathbf{u}_{\text{def}}, \quad (\text{D2})$$

where $\mathbf{u}_{\text{def}}$ is the remaining deformation component of the flow. Note that the translational and solid-body rotational components do not trigger odd viscosity; thus, the residual

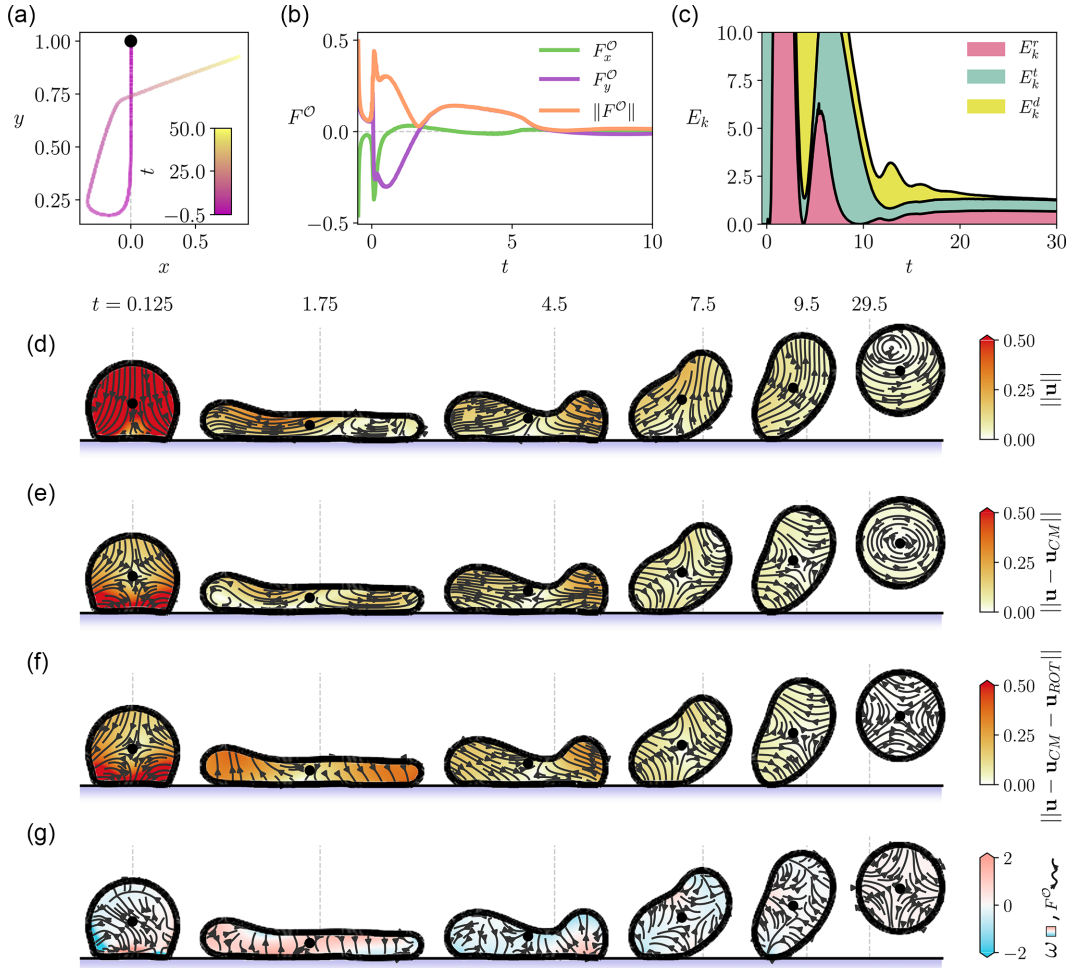


FIG. 9. Dynamics of a right-bouncing odd droplet ($\text{Re}^\varepsilon = 50$, $\text{Re}^\circ = 25$, $\text{We} = 50$). (a) Center-of-mass trajectory, marked by time. (b) Components of the average odd force over time. (c) Kinetic energy budget over time, where superscripts r , t , and d denote rotation, translation, and deformation, respectively. The colored regions are stacked cumulatively, such that the thickness of each band corresponds to the contribution of the associated component. At late times, nearly the entire velocity field consists of translation and solid-body rotation. (d) (Lab-frame) velocity fields at selected time instants. (e) Velocity fields with the pure translation component removed, revealing the late-time solid-body rotation induced by odd viscosity. (f) Deformation component of the velocity field. (g) Vorticity fields (background colors) and the odd force (streamlines). See Video IV for a dynamic visualization of this simulation [43].

velocity field shown in panel (f) is the only part capable of generating odd forces.

At very late times, nearly the entire velocity field is converted into either translation or pure rotation, resulting in a negligibly small deformation field and therefore suppressing significant odd-viscosity effects. To quantify these different contributions, panel (c) shows the total nondimensional kinetic energy within the droplet,

$$E_k = \text{Re}^\varepsilon \text{We} \int_{\Omega_1} \frac{1}{2} \|\mathbf{u}\|^2 d\Omega, \quad (\text{D3})$$

separated into its translational and rotational contributions,

$$E_k^t = \text{Re}^\varepsilon \text{We} \int_{\Omega_1} \frac{1}{2} \|\mathbf{u}_{\text{CM}}\|^2 d\Omega, \quad (\text{D4})$$

$$E_k^r = \text{Re}^\varepsilon \text{We} \int_{\Omega_1} \frac{1}{2} \|\mathbf{u}_{\text{ROT}}\|^2 d\Omega, \quad (\text{D5})$$

as well as the deformation kinetic energy, defined as the remaining contribution,

$$E_k^d = E_k - E_k^t - E_k^r. \quad (\text{D6})$$

Note that the nondimensional energy was obtained following the same rescaling presented in Eq. (A9), which leads to the correspondence $E = \frac{\mu_d^\varepsilon \sigma}{\rho_d U} \bar{E}$. For convenience, the upper bar was omitted in Eqs. (D3)–(D6) and in the rest of this article. As expected, at late times, nearly all kinetic energy is converted into translational and rotational motion. We note that, in a Newtonian (even) scenario, no rotation would occur since the droplet would maintain perfect left-right symmetry with zero average vorticity. Therefore, the presence of a rotational kinetic-energy component is a direct consequence of odd viscosity.

Figure 9 shows the velocity fields for an impact with $\text{Re}^\varepsilon = 50$, $\text{Re}^\circ = 25$, and $\text{We} = 50$, resulting in a droplet that bounces to the right. Immediately after impact, a negative

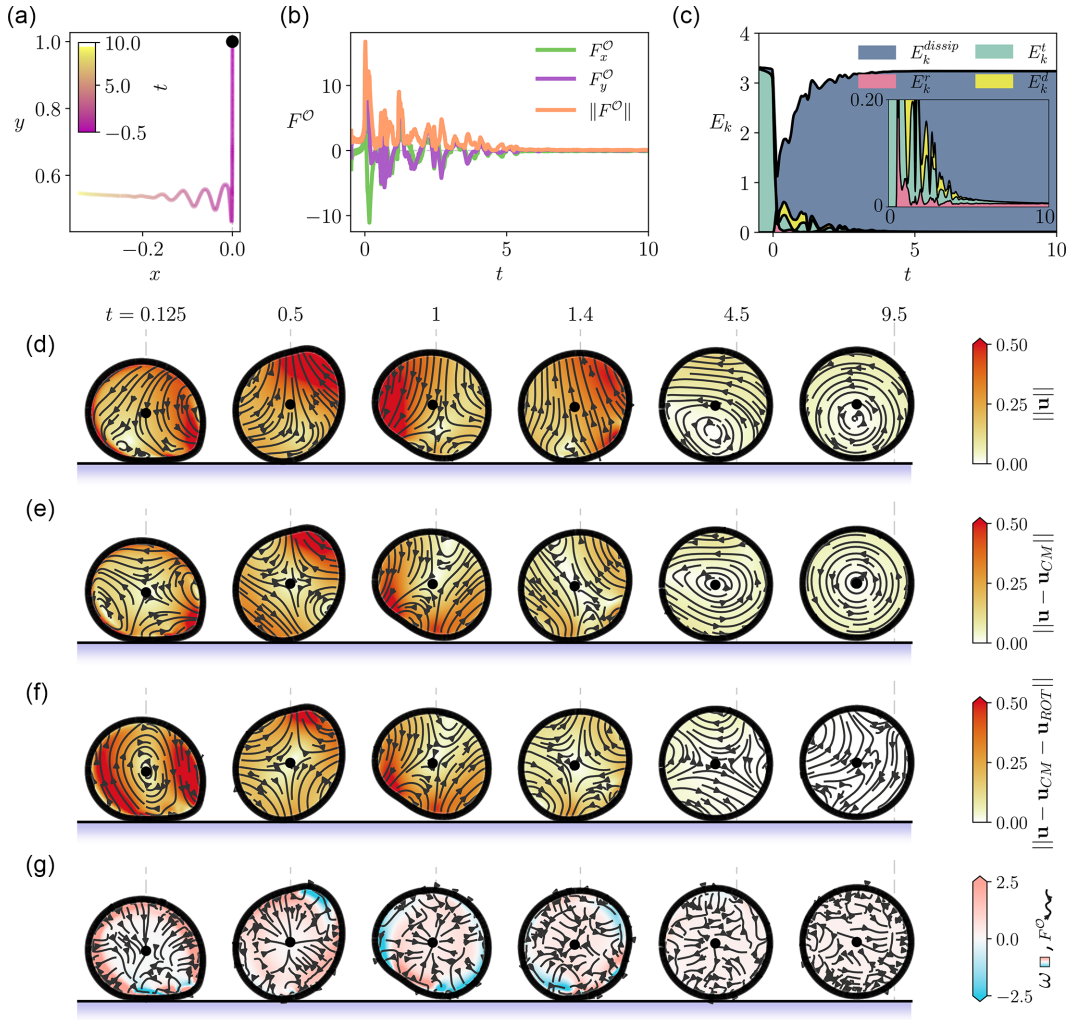


FIG. 10. Dynamics of a roller droplet with high odd viscosity ($\text{Re}^\mathcal{E} = 50$, $\text{Re}^\mathcal{O} = 1$, $\text{We} = 0.2$). (a) Center-of-mass trajectory, marked by time. (b) Components of the average odd force over time. (c) Kinetic energy budget over time, where superscripts r , t , and d denote rotation, translation, and deformation, respectively. We additionally show the dissipated energy, highlighting the strong dissipation that suppresses droplet rebound. The colored regions are stacked cumulatively, such that the thickness of each band corresponds to the contribution of the associated component. The inset provides a zoom of the y -axis, making it clearer that a residual rotational velocity field persists at late times. (d) (Lab-frame) velocity fields at selected time instants. (e) Velocity fields with the pure translation component removed, revealing the late-time solid-body rotation induced by odd viscosity. (f) Deformation component of the velocity field. (g) Vorticity fields (background colors) and the odd force (streamlines). See Video V for a dynamic visualization of this simulation [43].

peak in the average $F_x^\mathcal{O}$ is observed, as in the previous case (Fig. 8). The key difference is the larger maximum spreading reached by this droplet, which leads to extended spreading and retraction phases. During these longer phases, the positive $F_x^\mathcal{O}$ generated by the spreading and retraction is sufficient to accelerate the droplet to the right, overcoming the initial leftward velocity emerged at impact. Consequently, the droplet's horizontal translation is fully reversed to the right before it completely bounces off the surface. After bouncing, the subsequent dynamics are similar to those described for the case in Fig. 8.

Figure 10 shows the velocity fields for an impact resulting in a droplet that rolls over the surface. This case corresponds to $\text{Re}^\mathcal{E} = 50$, $\text{Re}^\mathcal{O} = 1$, and $\text{We} = 0.2$, representing a case with very high odd viscosity. The strong surface tension at this low Weber number, combined with the suppression of spreading due to high odd viscosity, results in minimal droplet

deformation after impact. In this simulation, in addition to the kinetic-energy decomposition given in Eq. (D6), we also show the cumulative energy dissipated by even viscosity, defined as

$$E^{\text{dissip}}(t) = 2 \text{We} \int_0^t \left(\int_\Omega \mu^\mathcal{E} \|\mathbf{D}\|^2 d\Omega \right) dt', \quad (\text{D7})$$

where $\mathbf{D} = \frac{1}{2}\dot{\boldsymbol{\gamma}}$ is the deformation-rate tensor. To ensure that the full initial kinetic energy is accounted for at late times, the dissipation is computed both inside and outside the droplet using the VOF transition function $\mu^\mathcal{E}$ from Eq. (A13). Immediately after impact, the flow generates a strongly localized vorticity layer near the droplet-ambient interface. This vorticity layer has predominantly the same sign along the interface, indicating the rapid emergence of coherent rotational motion. The localized vorticity layer, clearly visible in the first frame of panel (g), is spatially correlated with enhanced

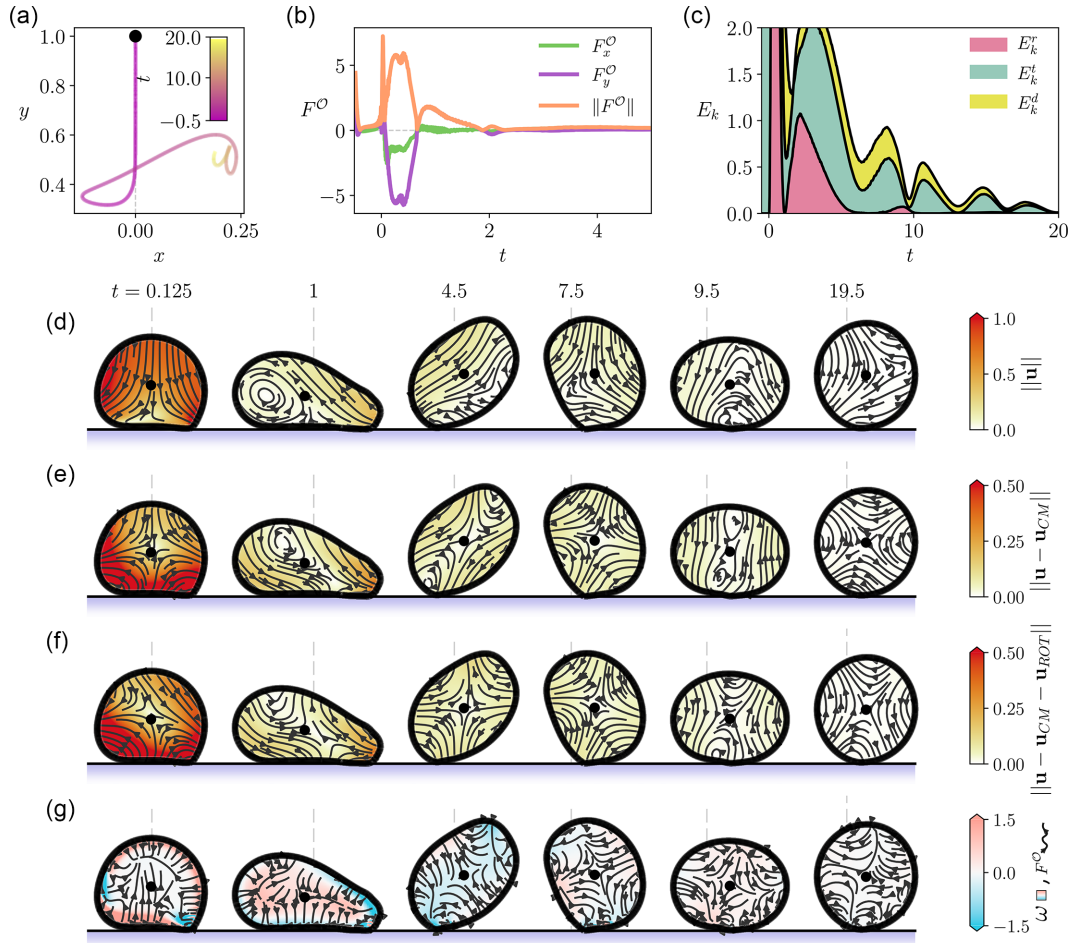


FIG. 11. Dynamics of a roller droplet with moderate odd viscosity ($\text{Re}^\varepsilon = 50$, $\text{Re}^O = 4$, $\text{We} = 7$). (a) Center-of-mass trajectory, marked by time. (b) Components of the average odd force over time. (c) Kinetic energy budget over time, where superscripts r , t , and d denote rotation, translation, and deformation, respectively. The colored regions are stacked cumulatively, such that the thickness of each band corresponds to the contribution of the associated component. (d) (Lab-frame) velocity fields at selected time instants. (e) Velocity fields with the pure translation component removed. (f) Deformation component of the velocity field. (g) Vorticity fields (background colors) and the odd force (streamlines). See Video VI for a dynamic visualization of this simulation [43].

even-viscous dissipation. This enhanced dissipation, visible in the energy budget of panel (c), leads to a rapid loss of kinetic energy during the impact stage. A fraction of the initial kinetic energy is also temporarily converted into surface energy through the small deformation experienced by the droplet immediately after impact, before being restored and dissipated as the droplet recovers its circular shape. As a consequence of the strong early dissipation, the droplet does not undergo sufficient spreading and retraction to sustain rebound. Nonetheless, part of the remaining kinetic energy is transferred into a persistent solid-body rotational mode, as confirmed by the inset of the energy decomposition in panel (c). Over time, nearly the entire velocity field is converted into this rotational component, so that the droplet effectively rolls over the surface. Due to the superhydrophobic, nonwetting properties of the surface, the coupling between the rotational motion and the substrate remains weak, resulting in only minimal net translation.

Figure 11 shows the velocity fields for an impact resulting in a droplet that slides over the surface and eventually stops.

This case corresponds to $\text{Re}^\varepsilon = 50$, $\text{Re}^O = 4$, and $\text{We} = 7$. Initially, the behavior is similar to that observed in previous figures: The droplet accelerates to the left during impact and then to the right during the spreading and retraction phases. The rightward acceleration is strong enough to visibly cause the droplet to slide to the right while still in contact with the surface. As shown by the kinetic energy distribution in panel (c), the leftward and rightward accelerations balance each other, leading to a rapid vanishing of the droplet's translational velocity. Although the kinetic energy continues to oscillate due to capillary-driven surface oscillations, it eventually dissipates through even viscosity, and the droplet comes to a near halt, reaching almost zero kinetic energy.

APPENDIX E: MINIMAL ODD MASS-SPRING MODEL

To provide a minimal phenomenological interpretation of the parity-breaking rebound dynamics, we consider a simple odd mass-spring (OMS) model [see Fig. 12(a)]. Such reduced inertial-capillary oscillator analogies have previously

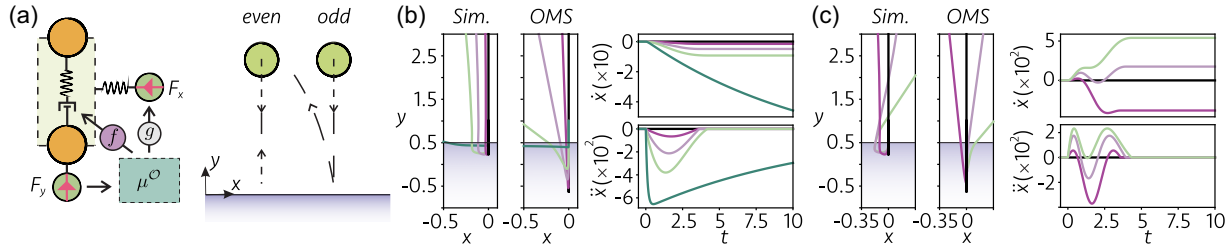


FIG. 12. Predicting dynamics with a minimal model. (a) An OMS system provides a simple model of odd droplet impact: odd viscosity leads to left-right symmetry breaking and dissipation through flow modulation. (b) Results of simulations are shown for $\text{Re}^O = \{+\infty, 100, 25, 10, 1\}$. The simple OMS model captures the left bounce and the transition to overdamped, fully suppressed droplets; nonetheless, differences remain apparent. The droplet in the OMS model experiences horizontal acceleration ($\ddot{x}$) while interacting with the wall, leading to a horizontal velocity ($\dot{x}$) that remains constant after bouncing (here, constants are $c_1 = 9.3$, $c_2 = 0.63$, $k = 0.78$, and $c_3 = 0$). (c) Introducing a nonlinear function g allows capturing both leftward and rightward bounces, as the total horizontal acceleration can change sign. Here, the simulations are for $\text{Re}^O = \{+\infty, 25, 8, 6\}$ and model constants are $c_1 = 6.04$, $c_2 = 0.63$, $k = 0.78$, and $c_3 = 3$.

been used to rationalize droplet impact and rebound dynamics [69–71]. Here, the purpose of the OMS model is not to quantitatively reproduce the full hydrodynamic evolution, but rather to illustrate how parity-breaking couplings can generically convert compressional motion into transverse dynamics and suppress rebound.

A conventional mass-spring system impacts and rebounds vertically from a wall [69]. In contrast, the odd mass-spring model incorporates nonlinear coupling between compression and transverse motion due to odd viscosity. For $t < 0$, the droplet falls vertically at a constant speed U . For $0 < t < t_b$, where t_b is the bouncing time, the droplet interacts with the wall ($y < 0.5$). During this period, horizontal forces are exerted, breaking the line symmetry of the trajectory. Additionally, odd viscosity modulates the flow field, leading to energy dissipation and suppression of bouncing. This interaction with the wall can be formulated as

$$m\ddot{y} + f(\mu^O)\dot{y} + k(y - 0.5) = 0, \quad (\text{E1})$$

$$m\ddot{x} + g(\mu^O)y = 0, \quad (\text{E2})$$

with initial conditions $x = \dot{x} = 0$, $y = 0.5$, and $\dot{y} = -U$. Equations (E1) and (E2) represent a standard damped harmonic oscillator in the y direction, coupled with a parity-breaking forcing in the x direction. We assume for $t > t_b$, the droplet continues its trajectory with zero acceleration, i.e., $\ddot{x} = \ddot{y} = 0$. The choice of functions f and g determines the system's dynamics. A simple choice is $f = c_1 |\mu^O|$ and $g = -c_2 \mu^O (1 - c_3 y^2)$, where c_i ($i = 1, 2, 3$) are constants. For small vertical deformation, we have $g \approx -c_2 \mu^O$, and the mass accelerates horizontally to the left. If the vertical deformation is large, the horizontal force may switch sign, causing the mass to accelerate to the right. Additionally, if the magnitude of dissipation in the y direction, due to flow modulation caused by odd viscosity, is sufficiently large, bouncing is completely suppressed. In Figs. 12(b) and 12(c), we compare the results of this simplified model with simulations. Despite its simplicity, the model captures the general trajectory dynamics. Nevertheless, notable differences arise, as it neglects even viscous dissipation and, more importantly, does not account for surface deformation or heterogeneous flow fields.

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