

Controlling droplets at the tips of fibers

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Many complex wetting behaviors of fibrous materials are rooted in the behaviors of individual droplets attached to pairs of fibers. Here, we study the splitting of a droplet held between the tips of two cylindrical fibers. We discover a sharp transition between two post-rupture states, navigated by changing the angle between the rods, in agreement with our bifurcation analysis. Depinning of the bridge contact line can lead to a much larger asymmetry between the volume of liquid left on each rod. This second scenario enables the near-complete transfer of an aqueous glycerol droplet between two identical vinylpolysiloxane fibers. We leverage this response in a device that uses a ruck to pass a droplet along a train of fibers, a proof-of-concept for the geometric control of droplets on deformable, architected surfaces.

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I. INTRODUCTION

Many materials in nature and industry are composed of collections of woven or aggregated fibers, or are decorated with hairs on their surface. When liquid invades the interstitial space between the fibers, diverse responses can result from a competition between elasticity and wetting, including clumping, matting, and self-assembly [1,2]. Yet, in a dense array of hairs, the channels between the hairs can prevent the liquid interface from penetrating into the structure [3]. This is particularly relevant in many biological settings, such as the hairy pads of insects [4,5], snails [6], geckos [7], and frogs [8], allowing them to adhere to or walk on wet or mucus-covered surfaces [9]. There, the interaction of a liquid interface with the *tips* of the micro-hairs is prevailing, if not more crucial, which raises the question: how does a liquid that sits atop a hairy surface behave as the material underneath is curled, twisted, or stretched?

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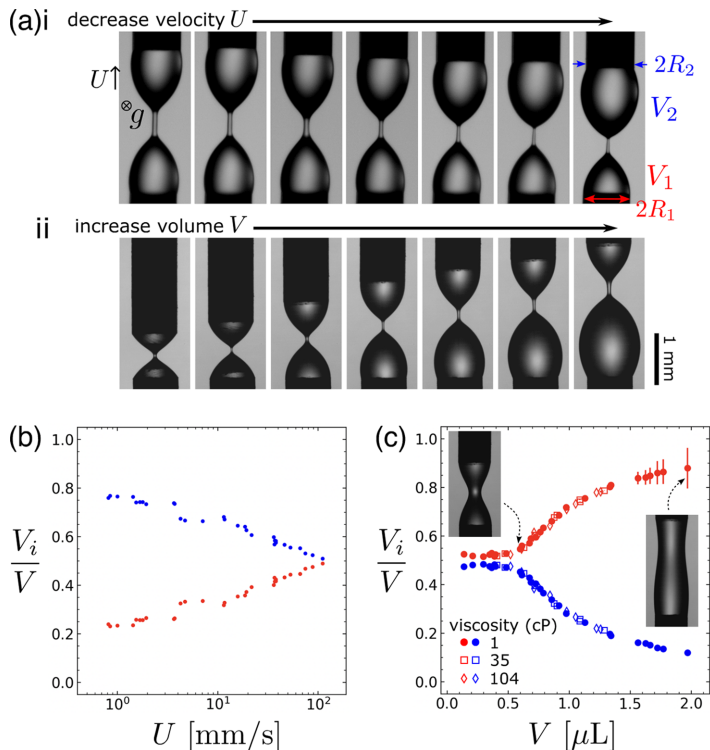


FIG. 1. Symmetry breaking in splitting a liquid bridge between two identical rods. (a) Snapshots of a liquid bridge pinching off from two separating rods into droplets of volumes V_1 and V_2 , for a range of (i) separation velocities U and (ii) total volumes V . Symmetry is gradually broken when U decreases or V increases. (b) Fractional volumes of the separate droplets in the post-rupture state versus U . (c) Fractional volumes of the separate droplets in the post-rupture state versus V . Inset: corresponding last stable states.

Before confronting multifiber, multidroplet systems, one may start by unraveling the fate of a single liquid bridge as it is pulled apart by the tips of two circular rods. The study of liquid bridges with fixed contact lines dates back to Plateau [10] and has received continuous attention throughout the centuries [11–14]. Surprisingly, a simple variation on this setup, in which one tilts one of the bridge ends and then pulls them apart, has been neglected. We show that the outcome of the liquid bridge rupture may be dramatically altered by small variations in the angle of the rods [Fig. 2(a)], in agreement with a bifurcation analysis.

Building on these results, we demonstrate how the system geometry and a depinning contact line can be coupled to enable a moderately nonwetting polymer (receding contact angle $\sim 60^\circ$) to achieve a near-complete transfer (99.995%) of a Newtonian liquid droplet from one fiber to another, a feat that is typically only achieved by functionalizing either the liquid [15] or the solid surface [16–19]. We harness this *geometric control* to pass a droplet down a train of fibers, highlighting the potential for soft metamaterials to control the transport of liquids in novel ways.

II. RESULTS

A. Experiment

In the experiment, we form a liquid bridge of de-ionized water (surface tension $\gamma = 72$ mN/m, density $\rho = 1000$ kg/m³) between the circular faces of two nearly identical solid pins (Vermont Gage A series, with radii $R_1 = 460$ μ m, $R_2 = 470$ μ m). We then draw the pins apart at speed U .

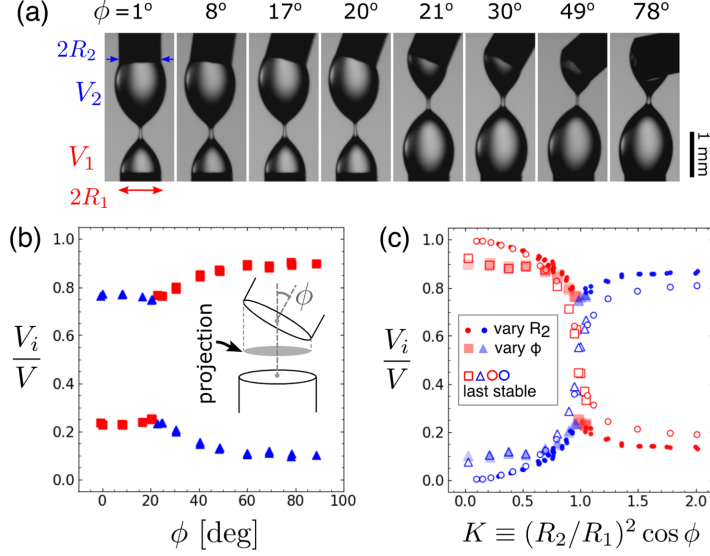


FIG. 2. Splitting a droplet between two angled rods. (a) Snapshots of a liquid bridge pinching off for a range of inclination angles ϕ . A sharp transition occurs near $\phi \approx 20^\circ$. (b) Fractional volumes of the separate droplets in the post-rupture state versus ϕ . (c) Fractional volumes versus the projected area ratio K for the same data shown in (b) (transparent), as well as for a pair of rods at $\phi = 0$ with varying R_2 (dots). Open markers: last stable states for both $\phi = 0$ (circles) and $\phi > 0$ (squares/triangles).

When the gap reaches a threshold size, the liquid bridge destabilizes and the liquid separates into two parts. We record this pinch-off process using a high-speed camera (Phantom Miro Lab 340). Our experiments are performed at a small Bond number $Bo = \rho g R_1^2 / \gamma \sim 10^{-2}$, where the droplets are negligibly deformed by gravity. To further ensure the minimal influence of gravity, we orient our setup horizontally [gravity points into the paper in Fig. 1(a)].

Figure 1(a)(i) shows repeated experiments at decreasing separation velocities U with snapshots taken approximately 0.1 ms before the liquid bridge breaks. At a high separation speed of $U \sim 100$ mm/s, the liquid bridge narrows at its midpoint and pinches off into two nearly symmetric droplets. This is intuitive, given the setup's symmetry of the boundary conditions. However, the symmetry is gradually broken when the process is slowed down. At a low separation speed ($U \sim 0.1$ mm/s, last image in the sequence), there is a significant difference in the sizes of the top and bottom droplets after they break apart. To quantify this effect, we measure the volume V_1 of the droplet remaining on the bottom rod, and that on the top rod V_2 versus U . Figure 1(b) shows that the volume left on each fiber varies approximately logarithmically with the velocity U .

A similar broken symmetry is observed by increasing the total volume of the liquid. The image sequence of Fig. 1(a)(ii) shows the uneven growth of the top and bottom droplet when the total volume of the liquid bridge increases. Figure 1(c) also shows that, starting from an even split, the volume distribution is first insensitive to the added volume when the total volume $V \lesssim 0.5 \mu\text{L}$. Then, for $V \gtrsim 0.5 \mu\text{L}$, the portion of the liquid left on each fiber becomes more and more disparate upon further increases in the initial volume V .

To investigate a possible role of the viscosity, we vary the droplet viscosity over two orders of magnitude by using mixtures of glycerol and water. Surprisingly, even though liquid pinch-off is a dynamic process, the amount of liquid left on each rod in the final state is observed to be independent of the droplet viscosity [different symbols in Fig. 1(c)]. We emphasize that the fluid velocities generated near the waist during the pinch-off are many orders of magnitude faster than the relatively slow velocity of the needles U .

Using high-speed imaging, we are able to trace the profile of the dynamic liquid bridge back to its marginally-stable state, where it is on the verge of destabilizing (which we call “last stable” hereafter). By examining the corresponding last stable states for each experiment, we find that the transition at $V \approx 0.5 \mu\text{L}$ corresponds to the point where the liquid interface meets the face of the rod at an approximately 90° angle at the last stable state [inset, Fig. 1(c)]. Below this volume, the liquid bridge splits nearly evenly between the rods.

The above results identify a strong asymmetry in both the last stable state and the final state, for the *quasistatic* splitting of a liquid bridge with large volume. What determines the uneven distribution? To investigate this question, Fig. 2(a) shows a series of experiments at sequentially larger relative angles ϕ , using the same set of rods and keeping the total volume V constant. Since the rotated rod has a slightly larger radius, $R_2 > R_1$, the top rod initially collects a larger fraction of the liquid, consistent with what was observed at $\phi = 0^\circ$ in Fig. 1. Strikingly, as ϕ increases to surpass a threshold of $\phi_{\text{thresh}} \approx 20^\circ$, the system goes through a sudden inversion so that the bottom rod retains more liquid. Figure 2(b) shows that there is a significant jump in the volume distribution (V_1 versus V_2) at an abrupt, discrete transition. At yet larger angles, the contact line partially moves to the side wall [last panel, Fig. 2(a)]. Despite this complication, the bottom rod still collects a larger fraction of liquid.

B. Projected area hypothesis

To rationalize this behavior, we start by proposing an intuitive picture of this transition. At $\phi = 0$, it is reasonable to expect that, regardless of the details of the dynamic pinch-off process, the larger rod will end up with a larger droplet [20]. For $\phi > 0$, the end of a static bridge should be narrower toward where it joins onto the tilted rod. One may thus expect that the *projected* area of the rod—which shrinks with increasing ϕ —could play an important role [see schematic in Fig. 2(b)]. This motivates us to consider the ratio K of the projected contact areas as the key control parameter:

$$K \equiv (R_2/R_1)^2 \cos \phi. \quad (1)$$

At $K = 1$, the projected area of the larger, tilted rod becomes equal to that of the smaller rod, which suggests a transition at $\cos^{-1}(R_1^2/R_2^2)$. As a first test of this intuition, plugging in values for R_1, R_2 for the case shown in Fig. 2(a) gives a transition at 18.9° , close to the measured transition at 20° .

This way of thinking rests on the following suggested equivalence: changing the projected area of the rod (by tilting it) has the same effect as changing the actual area of a *nontilted* rod. To test this idea, we perform further experiments where we split a drop using aligned rods ($\phi = 0$) while systematically varying the radius of the top rod R_2 . We plot the final volume proportions V_i/V as a function of the area ratio for the aligned rods in Fig. 2(c) (dots). In the same plot, we again present the data for rods at angles from Fig. 2(b), now using the projected area ratio K as the new variable (transparent squares/triangles). The transitions for both protocols are found to occur at $K = 1$. This result supports our proposed picture, wherein the projected area of the tilted rod should be compared with the area of the untilted rod, to determine which rod receives a larger droplet in the final state.

C. Threshold angle

We now test our projected area hypothesis [Eq. (1)] across a wider range of rod sizes and droplet volumes. We perform a series of splitting and coalescing cycles while adjusting ϕ each time to hone in on the precise transition point, ϕ_{thresh} . The hallmark of sitting at ϕ_{thresh} is the ability to obtain two distinct end states with an imperceptible change in initial conditions. Figure 3(a)(i) shows an example of such a fine-tuned state; we can then extract the profiles and superimpose them on one another (purple versus green outlines) to highlight the two possible states when the relative rod angle is in close proximity to ϕ_{thresh} . By allowing the water to slowly evaporate, we can map out ϕ_{thresh} as a function of the liquid volume V , which we measure *in-situ* from the images of the split droplets.

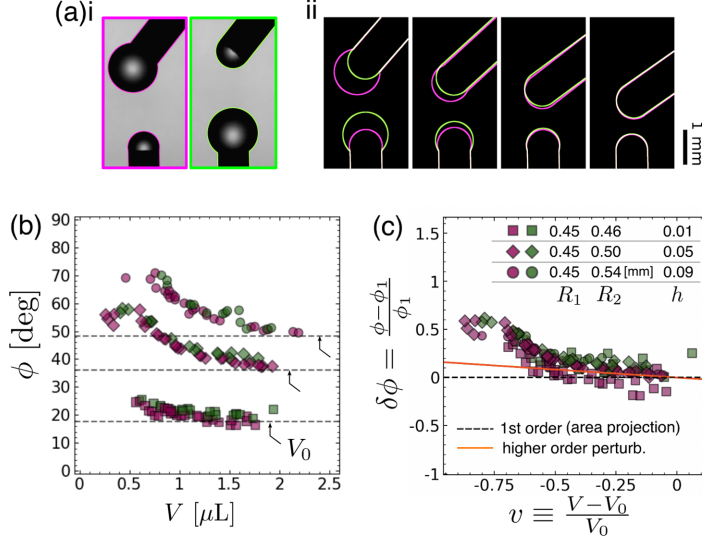


FIG. 3. Threshold angles for different droplet volumes and rod sizes. (a)(i): Two possible outcomes near the threshold angle. (ii): Superimposed outlines of the bifurcated states at threshold angles, as the total volume V is gradually reduced. (b) Measured ϕ near the bifurcations, for three sets of rod size ratios. Green: the bottom rod receives a larger droplet. Pink: the bottom rod receives a smaller droplet. Horizontal dashed lines: 1st-order perturbation theory Eq. (4). Arrows: corresponding Plateau limit $V = V_0$, where the perturbation theory is asymptotically accurate. (c) Data collapse in $\delta\phi$ - v space. Solid orange line: higher order perturbation theory.

As the volume decreases, it takes a larger inclination for the transition to occur, while the volume disparity in the bimodal end states decreases, as can be seen in the increasing overlap of the outlines [Fig. 3(a)(ii)]. We collect these data in Fig. 3(b), showing that ϕ_{thresh} decreases with increasing V for three values of R_2/R_1 that we tested.

D. Role of dynamics

The intricate process of a liquid bridge dynamically pinching into separate droplets occurs across multiple scales [21]. One may wonder: does this dynamics significantly redistribute the liquid across the bridge waist, or is the splitting fraction largely set already at the static shape determined by the size and relative angle of the two rods? To associate a well-defined fraction of volume to each rod at the last stable state, we slice the liquid bridge with an imaginary plane at the waist of the bridge [see Supplemental Material (SM) [22]]. We show additional data in Fig. 2(c), where the volumes are measured from the last stable bridge, V_i/V (open circles for $\phi = 0$; open squares/triangles for $\phi > 0$). These data show a close agreement with the droplets eventually separated from the liquid bridge. These results show that only near $K = 1$ do the pinch-off dynamics play a significant role in redistributing the liquid.

E. Bifurcation theory

We have seen so far that the outcome of the droplet-splitting process is captured reasonably well by the last stable shape before the onset of the dynamic pinch-off instability. Motivated by this observation, we obtain an approximation of the threshold angle ϕ_{thresh} from the last stable shape of the liquid bridge.

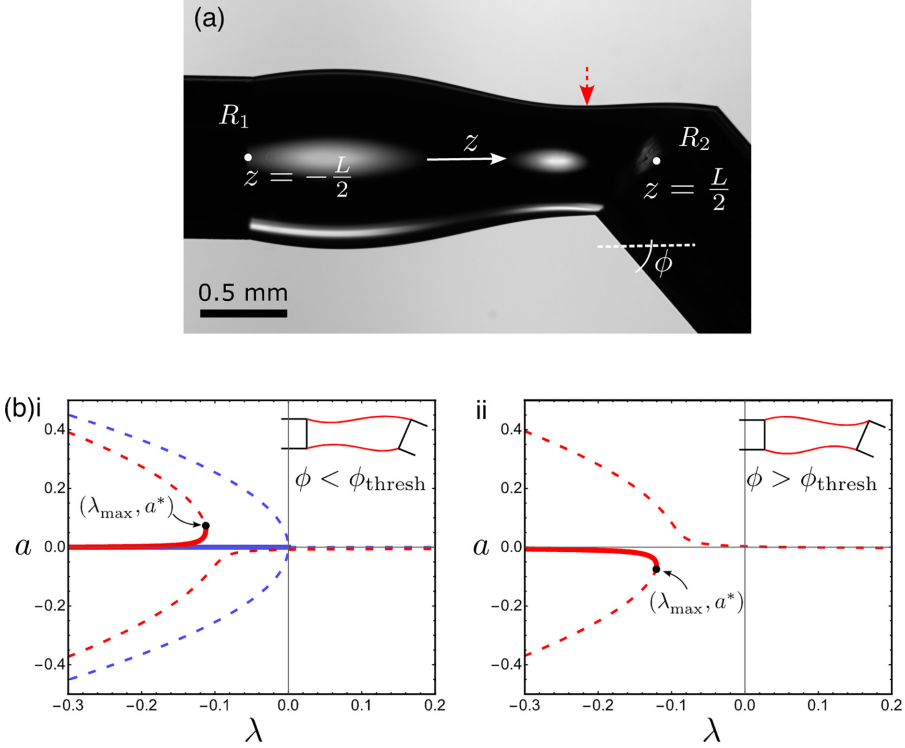


FIG. 4. Bifurcation analysis of tilted bridges. (a) Photograph of a liquid bridge spanning two rod faces with radii R_1 and R_2 , separated by L and tilted by ϕ . Red arrow: position of the bridge waist. (b) The amplitude a of the antisymmetric deformation as a function of the length parameter λ for $\phi = 0$ (blue) and $\phi > 0$ (red). The tilt angle ϕ controls the two unfoldings [(i) and (ii), red] of a perfect subcritical pitchfork bifurcation [(i), blue] with the change occurring across ϕ_{thresh} . The location of the waist changes drastically for the two unfoldings (inset schematic). Stable (unstable) branches are depicted as solid (dashed) curves, with a switch occurring at turning points $(\lambda_{\text{max}}, a^*)$ (black circles). The angled bridge bifurcations are plotted with $h = 0.053$ and $v = -0.1$.

Under cylindrical coordinates (z, θ) , force equilibrium at low Bond numbers requires that the profile $F(z, \theta)$ of a static liquid bridge satisfy:

$$\mathcal{M}[F(z, \theta)] = \text{const.}, \quad (2)$$

where $\mathcal{M}/2$ is the mean curvature of the profile [11]. Since the difference between the radii of the two rods is small, we solve Eq. (2) perturbatively around a reference cylinder of radius $R_0 = (R_1 + R_2)/2$ and length L , with its two face centers located at $z = \pm L/2$ [Fig. 4(a)]. At the cusp of collapse, the Plateau limit, where the length-to-diameter ratio $L/(2R_0)$ equals π , the reference cylinder has a volume of $V_0 = 2\pi^2 R_0^3$. To lowest order, the cylinder loses stability along a mode $\sin(2\pi z/L)$ that is antisymmetric about the midplane $z = 0$ [23]. This mode characterizes the dominant symmetry-breaking deformation of the reference cylinder towards a shape $F(z, \theta) = R_0(1 + a \sin(2\pi z/L))$ upon varying the angle of tilt ϕ , volume V , and rod radii away from their values in the reference cylinder. We define dimensionless parameters:

$$\lambda \equiv L/(2R_0) - \pi, \quad h \equiv (R_2 - R_1)/(R_2 + R_1), \quad v \equiv (V - V_0)/V_0$$

to represent these deviations in the length, rod radii, and volume of the reference cylinder, respectively.

Bifurcation theory enables us to express the amplitude a of the antisymmetric deformation as a series expansion in the parameters λ , h , v , and ϕ . As detailed in the SM [22], we present an analytical scheme utilizing the Lyapunov-Schmidt decomposition [24,25] to solve Eq. (2) [12,13,23,26], while tilting the circular top support with all contact lines pinned to the edges of the supports. The procedure yields the following bifurcation equation as a power series in $(\lambda, h, \phi, v; a)$:

$$G(\lambda, h, \phi, v; a) \equiv \frac{2h}{\sqrt{2+2\pi^2}} + \frac{3a\lambda - \pi av}{\sqrt{2+2\pi^2}} - \frac{8\pi h\lambda}{(2+2\pi^2)^{3/2}} - \frac{\phi^2}{4\sqrt{2+2\pi^2}} + \frac{3\pi a^3}{2\sqrt{2+2\pi^2}} - \frac{15\pi^2 a^2 h}{8\sqrt{2}(1+\pi^2)^{3/2}} + \dots = 0. \quad (3)$$

Equation (3) has been truncated for brevity and the full equation used in our analysis can be found in the SM [22].

Plotting the solutions of Eq. (3) for a as a function of λ elucidates the mechanism for the sharp inversion of the volume distribution upon an angle change, presented in Fig. 2. First, for a cylindrical bridge, setting $(h = \phi = v = 0)$ produces a perfect subcritical pitchfork bifurcation [23] [blue, Fig. 4(b)(i)], and the stability is lost at the Plateau limit $\lambda = 0$. Nonzero values of h , ϕ , v , however, lead to two possible unfoldings of the perfect pitchfork bifurcation [red, Fig. 4(b)], depending on if the tilt of the top support surpasses a threshold ϕ_{thresh} . As the two rods are pulled apart, the stability of the liquid bridge is lost at turning points $(\lambda_{\text{max}}, a^*)$ [black points, Fig. 4(b)]. More significantly, the amplitude a^* of the bridge profile switches its sign at $\phi = \phi_{\text{thresh}}$, giving rise to a *finite jump* of the waist position [red arrow, Fig. 4(a)] of the liquid bridge and therefore to a drastic change of the partitioned volumes.

To obtain the threshold angle ϕ_{thresh} , we effect a perturbation expansion:

$$\phi_{\text{thresh}} = \phi_1 \epsilon + \phi_2 \epsilon^{1+p_2} + \dots, \quad \lambda_{\text{max}} = \lambda_1 \epsilon + \lambda_2 \epsilon^{1+q_2} + \dots$$

for positive integers p_2 and q_2 , where the small quantity $\epsilon \sim 10^{-1}$ tracks the strength of the deviations set by our experiments: $v \sim \epsilon$ and $h \sim \epsilon^2$. At the turning points, $\partial a / \partial \lambda \rightarrow \infty$. We enforce the equivalent condition $\partial G / \partial a = 0$ so that the bridge is at the cusp of collapse. The Newton-Puiseux algorithm [27] provides a systematic procedure for obtaining the roots of the bifurcation equation (3) in powers of ϵ and, to leading order, we obtain $a^* = \epsilon^{2/3} (\frac{1}{12\pi})^{1/3} (8h - \phi_1^2)^{1/3}$ (see SM [22]). Therefore, up to order $\epsilon^{2/3}$, the amplitude a^* switches its sign at $\phi_1 = \sqrt{8h}$, so that to the first order of ϵ ,

$$\phi_{\text{thresh}} = \sqrt{8h}. \quad (4)$$

When $R_2/R_1 \approx 1$, it follows from Eq. (4) that $(R_2/R_1)^2 \cos \phi_1 \approx 1 - 4(R_2/R_1 - 1)^2/3 + \mathcal{O}((R_2/R_1 - 1)^3)$ so that $K \approx 1$ from Eq. (1), validating our hypothesis that the area projection is the leading cause for the volume distribution transition. In Fig. 3(b), we plot Eq. (4) (dashed lines), and show that despite neglecting higher order terms and the dynamics of pinch-off, our bifurcation analysis successfully predicts different transition angles for different pairs of rods. As expected, the predictions from the bifurcation analysis become asymptotically accurate for volumes approaching the Plateau limits of $V_0 \approx 1.9 \mu\text{L}$, $V_0 \approx 2.1 \mu\text{L}$, and $V_0 \approx 2.4 \mu\text{L}$, with R_2 increasing in this order [arrows, Fig. 3(b)]. However, the theoretical threshold angle is volume-independent at this order which deviates from our data at small volumes.

To capture the volume dependence of the threshold angle, we continue the perturbation expansion to higher orders. The next order correction is given by $p_2 = 1$ and $\phi_2 = -\sqrt{2h}v/3$, so that to order ϵ^2 ,

$$\phi_{\text{thresh}} = \sqrt{8h} - \frac{\sqrt{2h}}{3} v. \quad (5)$$

Taking into account the correction term with volume $V = 1.23 \mu\text{L}$, we obtain $\phi_{\text{thresh}} = 20.06^\circ$, in excellent agreement with the measured transition [Fig. 2(a)]. Further, Eqs. (4) and (5) suggest

that the deviation of the threshold angle from its lowest-order value, $\delta\phi \equiv (\phi_{\text{thresh}} - \phi_1)/\phi_1$, is proportional to the volume deviation v . We show in Fig. 3(c) that in the $\delta\phi$ - v space, our data indeed collapse, demarcating a narrow boundary between the two distributions (pink versus green). In the same space we have also plotted the first order [Eq. (4)] and the second order [Eq. (5)] solutions as the dashed and the solid lines, respectively. Remarkably, a perturbation expansion around a straight cylinder gives a quantitative explanation for the observed transition boundary of a liquid bridge down to a volume of 40% of a Plateau cylinder. For even smaller volumes $v \rightarrow -1^+$, higher order corrections play a non-negligible role, and the current expansion is insufficient. A detailed account of our bifurcation analysis and perturbation procedure can be found in the SM [22].

F. Complete droplet transfer

If the fluid can be coaxied into contacting a smaller area of the tilted rod, it may have the effect of mimicking a smaller rod that would collect a smaller share of the liquid. To test this idea, we perform droplet-splitting experiments using vinylpolysiloxane (VPS) fibers (Zhermack 8 shore A) of radius $540\ \mu\text{m}$, which we make by pouring mixed precursor fluids into custom molds and letting them cure, following Ref. [28]. We produce a clean, well-controlled fiber face by creating an excess thin film of VPS (much larger than the fiber face), which is peeled off to reveal a flat and smooth fiber face (see SM [22]). Since the residue droplets now can be much smaller, we use a water/glycerol mixture (viscosity $\eta = 104\ \text{cP}$) to reduce evaporation. Figure 5(a) shows the droplet-splitting process now is complex: the liquid contact line may depin and evolve simultaneously with the shape of the liquid bridge as the rods are pulled apart. The receding contact angle θ_r is measured to be approximately 70° , shown in Fig. 5(b). A substantial variation of the contact angle with the contact-line displacement s reveals a significant stick-slip motion. Before the contact line sweeps across the entire surface, the liquid bridge becomes unstable and pinches off from the bridge waist.

Despite this complexity, the end result gives a simple trend. Figure 5(c) shows the end states of a series of experiments splitting a liquid bridge of $V = 1.45\ \mu\text{L}$ between two VPS fibers at a sequentially larger angle ϕ , where the residue volume remaining in the top fiber decreases drastically at large angles. We measure the final droplet volumes on each fiber and show in Fig. 5(d) that they are relatively flat up to $\phi \approx 25^\circ$, when the contact line remains pinned. For $\phi \gtrsim 25^\circ$, the contact line evolves before pinch-off, leading to a sharp decline in the residue V_2 by four orders of magnitude. At $\phi \approx 70^\circ$, over 99.995% of the liquid transfers to the bottom fiber. As ϕ increases further, the liquid bridge starts to creep onto the side wall, leading to a slightly increased post-rupture residue. Here, we focus on the regime $\phi \lesssim 70^\circ$, where this complication does not occur.

G. Mesoscale simulations

To further explore the effect of the contact angle on the post-rupture volumes, we use mesoscale simulations using the Many-body Dissipative Particle Dynamics (MDPD) method, since it is much more convenient to tune the contact angle precisely [29,30] in these simulations by adjusting the liquid-solid interactions. This method is well suited for processes dominated by liquid viscosity and surface tension [31,32], ranging from droplet motions on rigid solid substrates [29], dynamic capillary wetting [33], self-cleaning of hydrophobic rough surfaces [34], and droplet wrapping of an elastic rod [30]. The simulations follow the same setup as in experiments with a liquid bridge ($1.45\ \mu\text{L}$) placed between two rods at different inclined angles, whose radii are close to $500\ \mu\text{m}$. The liquid has an effective surface tension of $21.2\ \text{mN/m}$ and kinetic viscosity of $2.9 \times 10^{-5}\ \text{m}^2/\text{s}$. The rods are discretized into particle clouds and the bottom rod is assigned a uniform velocity to mimic the rigid-body motion. To compare with our experiments given the measured contact angles [Fig. 5(b)], we ran our simulation at two contact angles, $\theta = 70^\circ$ and $\theta = 80^\circ$, and the resulting volume ratios are shown in Fig. 5(d) (solid and dashed lines). Note that the MDPD method yields an effective contact angle with negligible hysteresis, which is difficult to achieve in experiments.

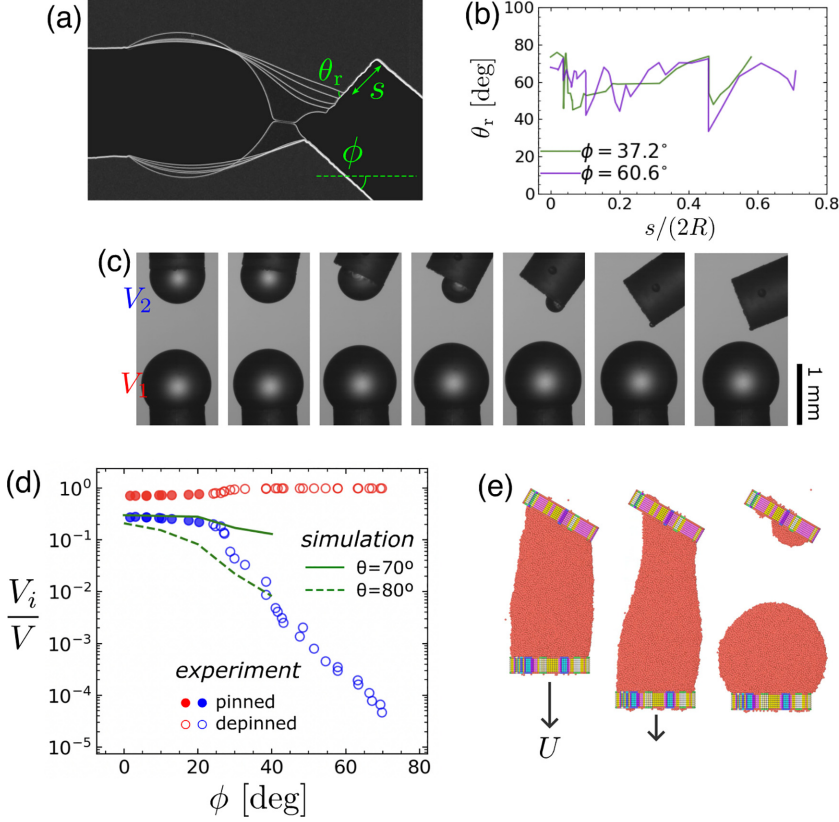


FIG. 5. Depinning the contact line allows a larger range of splitting ratios. (a) Image overlays of the profiles of a liquid bridge between the tips of two VPS polymer fibers at an inclination $\phi = 49.6^\circ$. (b) The receding contact angle θ_r as the contact line slides across fiber face, for two values of ϕ . (c) Post-rupture states between two VPS fibers with a range of ϕ . (d) Post-rupture volumes versus ϕ . The smaller droplet reduces its size by four orders of magnitude as ϕ increases, facilitating a near-complete transfer. Solid markers: pinned contact line at low ϕ . Open markers: depinned contact line at large ϕ . Solid and dashed lines: MDPD simulations for two effective contact angles θ . (e) Snapshots of the MDPD simulation with $\theta = 80^\circ$, exhibiting contact-line depinning.

Despite this simplification, the simulations capture the experimentally observed trend of V_2/V as a function of ϕ , as well as the depinning process of the contact line [Fig. 5(e)].

H. Pinned versus sliding contact line

We note that, so far, there are two possible scenarios for the contact line: (i) it remains pinned (for water on metal rods), (ii) it depins from the far edge to slide across the fiber face (for water/glycerol on VPS). The contact-line configuration generally depends on multiple scales of the triple-phase interface, involving solid curvature, surface chemistry, physical defects, and the history of its motion—a debated subject not fully understood to date [35–38]. However, different from a liquid spanning extended objects [39–43], in our experiments, the geometrically sharp corners of the rods’ faces help facilitate a range of contact angles through the Gibbs condition [44,45]. This allows us to treat the observed contact angles as the *outcome* of force balance on the bridge, rather than a boundary condition. We can thus target either pinning or depinning by picking appropriate rods (e.g., metal or VPS, respectively). Accordingly, without the need to design the magnitude of contact

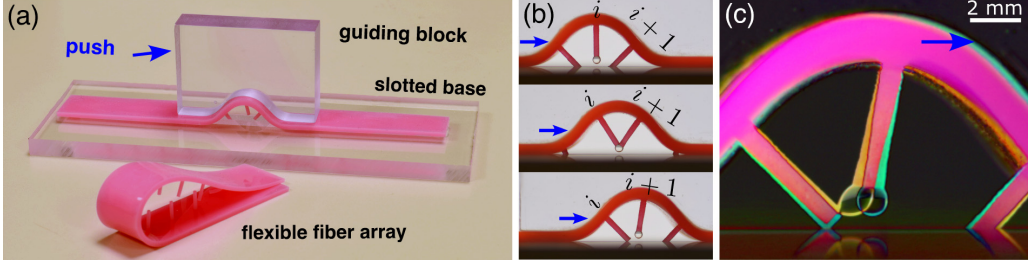


FIG. 6. Using material deformations to move a droplet from fiber tip to fiber tip. (a) A VPS elastomer is molded into a flexible fiber array. A guiding block imposes a ruck of the flexible substrate with designed angles between adjacent fibers. (b) Demonstration of a sustained droplet transfer, when the guiding block is continuously pushed in one direction (blue arrows). A drop is nearly completely transferred from site i to $i + 1$. (c) Subtraction of two consecutive frames at the moment of liquid bridge pinch-off.

angles, we have lumped the effects of the contact angles and their hysteresis into contact-line pinning (Figs. 1–4) or depinning (Fig. 5), which we have harnessed to the effect of droplet transfer.

I. “Caterpillarity”

The vanishing residue droplet on the VPS fiber face at large angles [Fig. 5(b)] suggests that a single droplet can be passed along a series of fibers, so long as the motions and angles of the fibers are well-controlled. To demonstrate this concept, we mold a VPS elastomer to form a hairy surface that is elastically deformable. The cured sample is laid down on a base plate with a slot where the fibers hang freely. We form a ruck in the sample and use a block to guide the ruck along the length of the sample, as shown by Fig. 6(a). The shape of the bump in the guiding block, as well as the length and the spacing of the fibers, are designed to bring pairs of adjacent fibers toward each other, one by one, along the underside of the sample.

To operate the device, we attach a droplet of glycerol/water mixture ($\eta = 104$ cP) to the end of one fiber i [Fig. 6(b), first panel]. The propagating bump rotates this fiber counterclockwise, bringing the droplet into contact with the neighboring fiber $i + 1$ to its right (second panel). As the bump passes through, fiber i rotates outward again clockwise with a large angle, leaving the droplet almost entirely on the next fiber (last panel). Figure 6(c) illustrates the pinch-off by showing the difference between two subsequent images. Crucially, the orientation and motion of the rods create the conditions for near-complete transfer from Fig. 5(b). In this way, the drop is transferred by one fiber in the traveling direction of the wave, so that it can repeat this process again. We show the entire process in the Movie in the SM [22].

III. CONCLUSION

It has been established since Savart [46] that a *free* liquid cylinder breaks due to an innate dynamic instability, insensitive to nozzle conditions [47,48]. In contrast, our angled two-rod system provides a route for external control over the end state of fragmentation. By simply turning the relative angle of the rods, a pair of binary post-rupture states can be accessed precisely. For the case of a pinned contact line, our experiments and analysis show that the transition between the post-rupture states is essentially independent of the liquid properties (density, viscosity, and surface tension), and is thus applicable for any liquid to the first approximation.

Delving deeper, a tilt in our opposing rods is linked to the creation of distinct turning points in an otherwise perfect pitchfork bifurcation. The robustness and sensitivity of our results are ultimately rooted in a *connectivity change* in the bifurcation diagram (Fig. 4). Unlike other biasing effects from the axial flow, gravity, or density [23,49], the angle change—combined with contact line motions—

provides an easily tunable, *geometric* handle that should apply at microscopic scales, where body forces fall to the wayside as surface forces become dominant.

Soft materials designed with a repeated unit cell can exhibit remarkable mechanical responses not inherent to their native materials [50]. Yet, the potential of designing mechanical metamaterials that can manipulate fluids are avenues that remain relatively unexplored [51]. Here we attain a high degree of control over a liquid droplet with an otherwise-unremarkable polymer surface, by harnessing the material geometry to transfer liquid from one pillar to another. Our device may serve as a proof-of-concept for the broader idea of using patterned solids to drive liquids and vice versa.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [53].

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Correction: A data availability statement and corresponding reference have been added.