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Chaotic Bouncing of a Droplet on a Soap Film

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We examine the complex dynamics arising when a water droplet bounces on a horizontal soap film suspended on a vertically oscillating circular frame. A variety of simple and complex periodic bouncing states are observed, in addition to multiperiodicity and period-doubling transitions to chaos. The system is simply and accurately modeled by a single ordinary differential equation, the numerical solution of which captures all the essential features of the observed behavior. Iterative maps and bifurcation diagrams indicate that the system exhibits all the features of a classic low-dimensional chaotic oscillator.

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Couder et al. [1] have recently shown that oil droplets, when placed on a vertically vibrated oil bath, may bounce indefinitely rather than coalescing. The dynamics of the bouncing droplets are extraordinarily rich. Feedback between the droplet and its wave field may lead to self-propulsion [2] and diffraction of these walking droplets as they pass through a slit [3]; moreover, multiple droplets may lock into complex orbital motions [2] or lattices [4]. We here demonstrate that a droplet on a vertically vibrated soap film may similarly avoid coalescence, and that the bouncing droplet represents a textbook example of a chaotic oscillator, with many features common to the bouncing of an inelastic ball on a solid substrate.

Drops of uniform size ($R = 0.08$ cm) bounce on a horizontal circular soap film of radius $A = 1.6$ cm vibrated with vertical displacement $B \cos(2\pi ft)$ (Fig. 1). The droplet and soap film consist of a glycerol-water-soap mixture (80% water, 20% glycerol, <1% soap) with density $\rho = 1.05$ g cm$^{-3}$, viscosity $\nu = 2$ cS, and surface tension $\sigma = 22$ dyne cm$^{-1}$. Drops of uniform size ($R = 0.8$ mm) and mass ($m = 2.25$ mg) were extruded from an insulin syringe (needle diameter 0.35 mm). For the bouncing states, the characteristic drop impact speeds ($U \sim 4$–32 cm s$^{-1}$) are much less than the characteristic wave speed on the film (a film thickness of 4 $\mu$m indicates a wave speed of $\sqrt{2\sigma/\rho h} \sim 330$ cm s$^{-1}$). The influence of capillary waves is thus assumed to be negligible, and the film described as quasistatic: it deforms instantaneously in response to the forcing imposed by the droplet. The Weber number $We = \rho U^2 R / \sigma$ lies between 0.06 and 3.9. During impact, the droplet remains roughly spherical: maximum center line distortions of 13% were observed (at $We = 3.9$), so the surface energy of drop distortion is less than 3% that associated with soap film distortion. Beneath the droplet, the soap film lies tangent to the droplet, and so roughly assumes the form of a spherical cap of radius $R$. Beyond the droplet, the pressure is atmospheric on either side of the soap film, which thus assumes the form of a catenoid as confirmed experimentally [5]. The spherical cap and catenoid match at a point $M$ corresponding to an angle $\alpha$ (Fig. 1). The vertical deflection of the soap film $Z$ and the resulting vertical force $F$ on the droplet may be expressed in terms of $\alpha$:

$$ Z = 1 - \cos \alpha + \sin^2 \alpha \ln \left[ \tan \frac{\alpha}{2} \frac{R}{Z} \left( 1 + \sqrt{1 - \beta^2} \right) \right] \frac{Z}{R}, \tag{1} $$

$$ F = 4 \pi \sigma R \sin^2 \alpha, $$

where $\beta = (R \sin^2 \alpha) / \alpha$. The force-displacement relation $F(Z)$ is shown in Fig. 2(a). In the range $0 < Z / R < 3$, the film responds as a linear spring, $F = k Z$, where the effective spring constant $k = (8 \pi / 3) \sigma$. The force then saturates, achieving a maximum at $\alpha = \pi / 2$, and decreases thereafter. This quasistatic description of the film was used successfully by the authors [5] to deduce a criterion for breakthrough of a droplet striking a stationary film [6]. When the droplet strikes a static soap film at a speed $U$, the drop is in apparent contact with the film for a time $\tau_c$. Dissipation during rebound results in the kinetic energy at takeoff being less than that at impact. Figure 2(b) illustrates the dependence of the dimensionless contact time $\tau_c = \tau_c / \tau_o$ and dissipated energy $\Delta(V^2 / 2)$ on the dimen-

![FIG. 1. Experimental system: a droplet of radius $R = 0.08$ cm bounces on a soap film of radius $A = 1.6$ cm vibrated with vertical displacement $B \cos(2\pi ft)$. The soap film assumes the form of a spherical cap beneath the droplet, and a catenoid beyond the matching point $M$.](image)
The film forcing provides energy to the droplet. If this energy precisely balances that lost through dissipation during impact, the droplet executes a periodic bouncing motion. We denote by \((m, n)\) a periodic state in which the droplet bounces \(n\) times and the soap film oscillates \(m\) times during a single period. A myriad of simple \((n = 1)\) and complex \((n > 1)\) periodic states were observed experimentally, including \((1, 1), (2, 1), (3, 1), (3, 3),\) and \((2, 2)\) (Fig. 3). Mode transitions characterized by either aperiodic transients or period-doubling cascades were observed as the forcing parameters were varied (see supplementary material in Ref. [8]). Multi-periodicity is apparent in Figs. 3(a)–3(c): multiple periodic solutions \((1, 1), (2, 1),\) and \((3, 1)\) arise for precisely the same forcing parameters \((\omega, \Gamma) = (0.6, 1.1),\) but different initial conditions. With \(\omega = 0.6\) fixed, complex periodic states were apparent at higher \(\Gamma\) [Figs. 3(d) and 3(e)], and ultimately chaos emerges for \(\Gamma \geq 1.1\) [Fig. 3(f)]. As seen in Figs. 4(b) and 4(c), the computed periodic solutions of Eq. (2) are remarkably close to those observed experimentally; in particular, the landing and takeoff phases are in good agreement. As a caveat, we note that complex modes \((m,\)
defined as

\[ \text{from one impact to the next, for various initial conditions represented by vertical lines denoted by } \frac{V}{C_{0}}(1, 1), (2, 1), \text{ and } (3, 1) \text{ at } \frac{V}{C_{0}} = 0.6. \]

The landing and takeoff phases measured experimentally are \(\frac{V}{C_{0}} \approx 1.82\), respectively. For each set of initial conditions \((V_i, \phi_i)\), contours of the net energy acquired by the drop during impact, \(V_{i+1}^2 - V_i^2\), are computed. The shaded area corresponds to initial conditions for which energy is gained during impact; in the white area, energy is lost. Simple modes \((m, 1)\) are represented by single points that necessarily lie on the boundary of the shaded area: no energy is gained or lost during impact, so the bouncing is perfectly periodic over a single forcing cycle. Complex modes \((m, n > 1)\) are represented by closed curves that cross the zero-energy boundary. For example, in the \((3, 3)\) mode \([\text{Figs. 3(d) and 5(a)}]\), energy is transferred to the droplet during the first two impacts, increasing the leap height; however, during the third bounce, energy is lost and the initial conditions are recovered. At \(\Gamma = 1.82\), the motion is chaotic, and a strange attractor emerges on the Poincaré section \([\text{Fig. 5(b)}]\). (See supplementary material in Ref. [8] for a Smale map diagnostic.)

A bifurcation diagram \([\text{Fig. 6(a)}]\) represents the solution of Eq. (2) as a function of \(\Gamma\) for \(\omega = 1.1\). For \(\Gamma < 0.17\), no periodic bouncing states are possible: the droplet resides at rest on the soap film. The first periodic solution \((2, 1)\) appears at \(\Gamma = 0.17\). As \(\Gamma\) is progressively increased, simple modes \((m, 1)\) and complex modes \((m, n)\) appear in turn. At branch points, modes \((m, n)\) execute a period-doubling transition to a \((2m, 2n)\) mode. For example, the mode \((2, 1)\) gives rise to a period-doubling cascade that terminates at \(\Gamma \approx 1.67\) \([\text{Fig. 6(b)}]\). Thereafter, the complex periodic modes degenerate into a strange attractor that may coexist with stable periodic orbits emerging from other branches \([\text{shaded area in Fig. 6(a)}]\). For \(\Gamma > 1.91\), no such periodic branches persist, and the chaotic attractor is the only attracting set.

\[ V_{i+1} = f(V_i, \phi_i), \quad \phi_{i+1} = g(V_i, \phi_i). \]

Poincaré sections are represented in Figs. 5(a) and 5(b) for \(\Gamma = 0.82\) and \(\Gamma = 1.82\), respectively. For each set of initial conditions \((V_i, \phi_i)\), contours of the net energy acquired by the drop during impact, \(V_{i+1}^2 - V_i^2\), are computed. The shaded area corresponds to initial conditions for which energy is gained during impact; in the white area, energy is lost. Simple modes \((m, 1)\) are represented by single points that necessarily lie on the boundary of the shaded area: no energy is gained or lost during impact, so the bouncing is perfectly periodic over a single forcing cycle. Complex modes \((m, n > 1)\) are represented by closed curves that cross the zero-energy boundary. For example, in the \((3, 3)\) mode \([\text{Figs. 3(d) and 5(a)}]\), energy is transferred to the droplet during the first two impacts, increasing the leap height; however, during the third bounce, energy is lost and the initial conditions are recovered. At \(\Gamma = 1.82\), the motion is chaotic, and a strange attractor emerges on the Poincaré section \([\text{Fig. 5(b)}]\). (See supplementary material in Ref. [8] for a Smale map diagnostic.)

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During the period-doubling cascade of (2, 1), bifurcations occur at \( \Gamma = 1.361, 1.631, 1.6679, 1.67244, 1.67319, 1.673349, 1.673380, \ldots \). Defining \( \delta_i = (\Gamma_{i+1} - \Gamma_i)/(\Gamma_{i+2} - \Gamma_{i+1}) \) yields the first terms of the \( \delta \) suite as 7.3, 8.1, 6.0, 4.7, 5.1, \ldots. This suite slowly converges to a value larger than the Feigenbaum constant 4.6692, which is universal for one-dimension quadratic iterative maps [9].

We recall that for area-preserving two-dimension maps, Tabor [10] deduces a universal constant of 8.7211. Our system shares many features with the elastic ball bouncing on a vertically vibrated rigid substrate [11], including multiperiodicity and period-doubling transitions to chaos. However, qualitative differences exist between the two systems owing to the differences in the collisional dynamics. In the bouncing ball problem, the collision is instantaneous and characterized entirely by the coefficient of restitution; in our system, the collision is of finite duration and the coefficient of restitution depends on the impact speed. Consequently, in our system there are no sticking solutions as arise at weak forcing for the bouncing ball, and the multiperiodicity is considerably enhanced. It is hoped that our study will inform studies of droplets bouncing on a fluid bath [1–3,12,13]. For example, for a droplet bouncing on a soap film, the most unstable bouncing state may be (2, 1) rather than the (1, 1) observed for the bouncing elastic ball. Finally, we note that in terms of ease of both experimental study and theoretical description, this system is perhaps the simplest fluid chaotic oscillator yet explored.

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**FIG. 6.** (a) Bifurcation diagram for \( \omega = 1.1 \): impact speed \( V \) as a function of the forcing acceleration \( \Gamma \). Note that a mode \((m, n)\) necessarily has \( n \) branches. Chaotic solutions first appear at \( \Gamma = 1.67 \). (b) Period-doubling cascade of the (2, 1) branch. The shaded region is expanded in the inset.