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Leidenfrost levitated liquid tori

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Abstract – Levitating a liquid over a vapor film was limited to droplets. Here we show that on curved substrates a larger quantity of fluid can be suspended. This opens a new possibility for exploring free-liquid-surface phenomena without any contact with a solid. In one of the simplest possible situations, a large fluid torus is levitated over a circular trough. A poloidal flow inside the ring generates a wave on its inner side, making it polygonal. This wave is described by a solitonic model which balances surface tension and the pressure depletion due to the distortion of the poloidal flow.

The situations in which any fluid can be steadily kept isolated from any contact with a dense medium (solid or liquid) are scarce and often obtained in extreme situations like microgravity for liquids or electromagnetic confinement for toroidal plasmas. A third (and more accessible) possibility is provided by the Leidenfrost effect [1] where a liquid, deposited on a very hot plate, is kept levitating over the thin vapor film created beneath the drop by evaporation. The shape of these drops results from the competition between surface tension and gravity and thus scaled with the capillary length $l_c = \sqrt{\gamma/\rho g}$ (where γ is the surface tension, ρ the fluid density and g gravity). For water drops of a diameter of one centimeter or less, the Laplace pressure is large enough to obtain stable shapes. In this limit the underlying vapor escapes easily by the side of the near-contact area [2,3]. If a larger quantity of liquid is deposited, the effect of gravity causes the drop to take a pancake shape of thickness of $2l_c$. On a planar substrate, the region of near contact becomes too large, and all the vapor generated underneath cannot entirely escape around the edges. Some of it accumulates and forms bubbles that move through the fluid layer and disturbs it violently [4]. Due to these perturbations, even though it provides a free-surface situation, the use of the Leidenfrost effect has been limited to small drops, such as their deformation modes [5], or self-propulsion on a ratchet substrate [6].

The aim of the present article is twofold. We first show that by the use of non-planar substrates it is possible to levitate quantities of liquid orders of magnitude larger than the traditional Leidenfrost effect, without any formation of bubbles at the surface. As an example we show in a second part that the waves propagating on one type of levitated liquid rings are generated by a secondary instability of the toroidal vortex it contains. We then illustrate that obtaining a fluid bounded only by free surfaces is a tool for the investigation of the coupling between hydrodynamics instabilities and surface waves.

In order to circumvent the size limitation, the idea is to use non-planar substrates such that on their whole surface one of the two main curvatures is large. As the levitated fluid spreads into the troughs, there is everywhere a path for the vapour to evacuate easily towards the edges. The vapour evacuation is thus organized by the imposed geometry. A large number of possible configurations of troughs meets these requirements. They can be circular or form various grid patterns. The general topology of the levitated liquid is a torus with multiple holes. In this letter, we limit ourselves to a simple geometry where the substrate has only one circular trough, which allows us to stabilize a large single-hole liquid torus (see fig. 1).

The substrates we used are 20 mm thick brass plates with circular troughs of several profiles and different radii R_G (see fig. 1). The brass substrate is placed on a

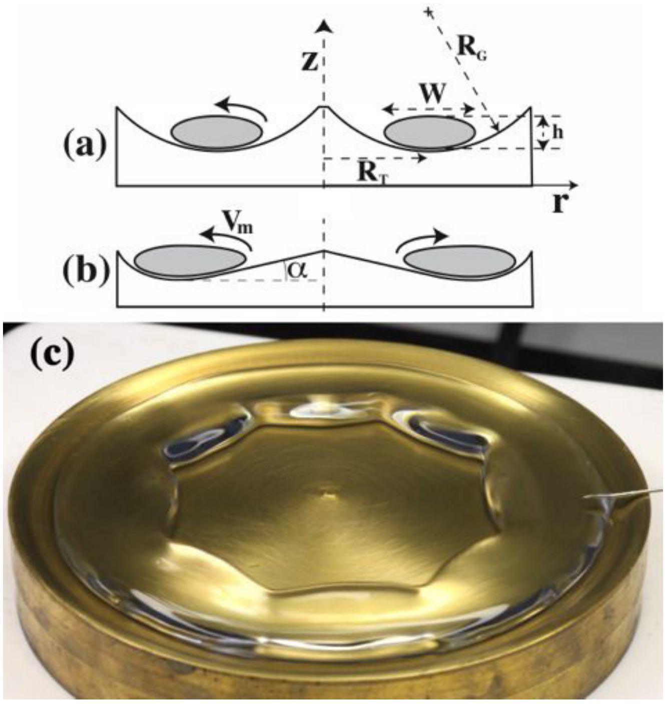

Fig. 1: (Colour on-line) (a) Sketch of the toroidal configuration. (b) An asymmetric variant with a trough having a gentle slope inward and a steep slope outward. (c) Photograph of a polygonal pattern observed in the configuration (b).

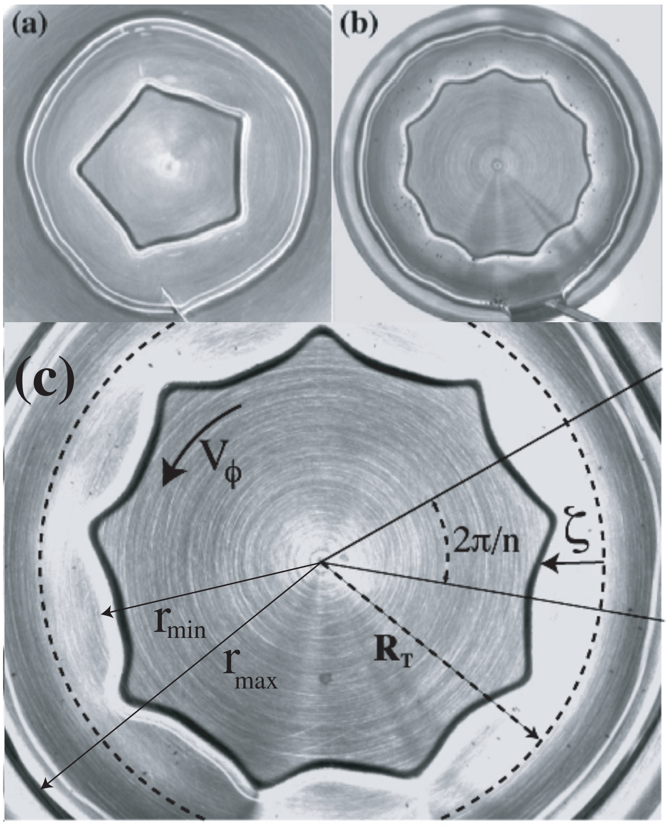

Fig. 2: Photographs of three polygonal modes observed on two different substrates: (a) for $R_T = 30$ mm, (b) and (c) for $R_T = 50$ mm, using two different volumes of water, which correspond to different number of sides.

horizontal 1000 W temperature-regulated hot plate heated at T between 200 °C and 400 °C. When distilled water is poured into the trough, a large crescent-shaped levitated drop forms. When more water is added the extremities of the crescent drop connect: a levitated liquid torus is formed. It thins down with time due to evaporation. In order to compensate this natural decay, we used a motorized syringe to create a weak constant flow Q of hot water into the levitating annulus. The torus thus eventually reaches a steady regime, where evaporation and injection rates are equal.

Only the surface of the transverse section is now limited, by the appearance of bubbles for the width and by the capillary length for the height. These limitations lead to a maximum area of the vertical section of few cm^2. However, no limit has been observed in the azimuthal direction. A liter of water could be set in levitation with a circular trough of radius 30 cm. The profile of the section of the liquid torus depends on the substrate geometry defined in fig. 1. If the trough is deep and narrow (*i.e.*, with R_G of the order of few l_c and $R_G < R_T$), a stable liquid annulus is obtained. For shallower troughs as shown in fig. 1(b) where the radius $R_G > R_T$ the levitated liquid torus is flattened by gravity into a doughnut shape and an interesting phenomenon is observed at the surface on which we will now focus.

We study steady regimes that are controlled by the temperature of the substrate and the injection rate. At low substrate temperature ($T < 220$ °C) and when the annulus has a small section the levitated liquid is at rest and a stable annulus is observed. As the temperature is increased there is a threshold depending both on temperature and torus dimension above which a poloidal flow appears, the fluid rises up on the inner side of the ring. These fluid motions are sketched in fig. 1(b). There are two possible origins for the existence of the toroidal flow. An estimate of the temperature difference from the bottom of the liquid to the top (4 K) suggests that the surface tension gradient is large enough to trigger a Marangoni convection [5]. Furthermore the escape motion of the vapor could also drag the liquid into motion [6,7]. Both processes lead to instability so that it is unclear why motionless regimes can exist. This is an interesting question that is beyond the scope of the present article. We limit ourself to an experimental description. The local velocities at the surface of the torus have been measured by following the motion of tracers of diameter $d = 50\,\mu$m as they are trapped by surface tension at the surface. The mean velocity V_m is of the order of 50 mm/s and does not depend on the torus width W. The striking observation is that when this liquid flow appears, a propagative wave grows on the inner side of the torus, forming a polygonal pattern rotating in the laboratory frame (see fig. 2). On all the shallow substrates of the type shown in figs. 1(a) and (b), the permanent

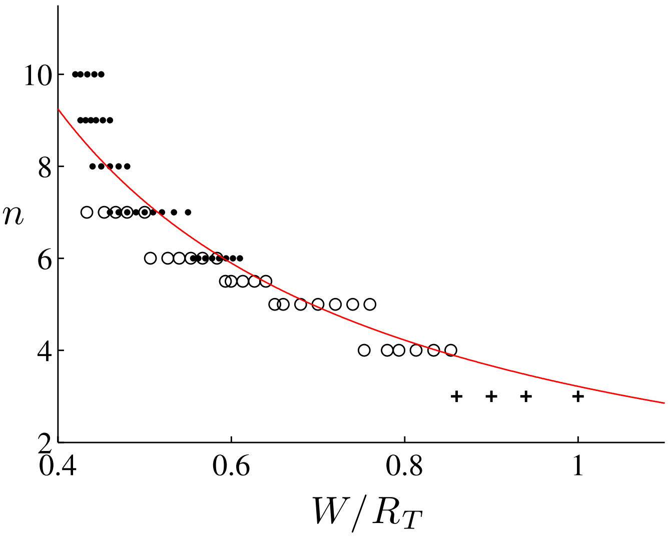

Fig. 3: (Colour on-line) Number of polygon sides as a function of the aspect ratio W/R_T of the torus for three different substrates: $R_T = 15$ mm (+), $R_T = 30$ mm (○) and $R_T = 50$ mm (·). The wavelength $\lambda = 2\pi R_T/n$ is approximately proportional to the torus thickness so that $n \sim R_T/W$ (solid line).

injection of water leads to stable and reproducible polygonal patterns (see figs. 1(c) and 2).

In order to study this instability quantitatively, we used specially designed substrates in which the trough has an asymmetrical profile (see fig. 1(b)) with a weak constant slope inwards and a steep edge outwards that fixes the external perimeter of the doughnut. The number n of sides increases continuously with the doughnut perimeter. Using three substrates of different perimeters $2\pi R_T$, we observed polygons having from three to twelve sides (see fig. 2). The wavelength $\lambda = 2\pi R_T/n$ is found to be of the order of the width W so that $n \propto R_T/W$ (see fig. 3). In the laboratory frame, these polygonal patterns rotate along the azimuthal direction θ with a linear velocity V_θ. The rotation has an equal probability to be clockwise or counterclockwise, but this degeneracy can be lifted. Since the friction of the liquid on the substrate is very weak, the whole liquid annulus can be set into an azimuthal rotation with a velocity U by means of a weak air jet impinging tangentially on the upper surface. The wave is then always counter-propagating so as to minimize its velocity in the laboratory frame. For all values of U, the velocity V_θ of the wave is constant in the liquid frame of reference. We find $V_\theta = 105 \pm 5$ mm/s for the asymmetrical trough shown in fig. 1(b). This is twice lower than the usual minimum of velocity for capillary gravity waves at an equivalent depth of 5 mm corresponding to the height of the levitated ring. This can be easily understood by considering that the interface does not move vertically but along the substrate slope forming an angle $\alpha = 5°$ with the horizontal (cf. fig. 1(b)). A good interpretation of the observed velocity is obtained by replacing in the usual dispersion relation of surface waves, g by a reduced value $g\sin\alpha$ and the fluid depth by the width W of the torus.

The shape of the polygons, the experimental geometry, and the liquid flow inside the torus are reminiscent of other faceting instabilities observed in partially filled rotating tanks [8,9], and in circular hydraulic jumps [10,11]. In all these systems, a circular front limiting the boundaries of a toroidal vortex becomes polygonal. This constitutes one of the most intriguing phenomena observed at the frontier between vortex dynamics and free-surface flows. An open question is whether the mechanism of instability is the same in all these systems [12,13]. Our experimental set-up offers the opportunity to study this question in ideal conditions, *i.e.*, where the vortex is clearly delimited as it coincides with the position of the free surface.

When zooming in one wavelength, an obvious feature appears: the strong asymmetry between the regions where the fluid protrudes inwards (the sides of the polygons) and the constrictions (the corners). The position of the inner side interface is defined by $\zeta(\theta)$ the radial distance to the azimuthal axis R_T of the torus using the conventions given in fig. 2. For a given mode the amplitude of the wave increases with the imposed temperature. We measured systematically the wave profile $A(\theta) = \zeta - \zeta_{max}$ for various amplitudes $A_0 = |\zeta_{min} - \zeta_{max}|$ obtained for several temperatures T between 220 °C and 400 °C. They are shown in fig. 4(a) for a pattern $n = 6$. In order to characterize these shapes, we measured the angular width $\Delta\theta$ at mid-height as a function of the amplitude A_0 for all patterns (see fig. 4(b)). Thus the patterns of various orders can be compared. Figure 4(c) is a plot of ζ as a function of the normalized angle $n\theta/2\pi$ for patterns with $n = 6, 7, 8$ where the wave amplitude is the same. The wave profile remains identical, the apparent differences observed in fig. 2 being only due to the mean curvature imposed at large scale by the geometry. This suggests that we can describe the wave profile by a universal equation with only one parameter fixing both its amplitude and its width. Here we propose a simple heuristic model to obtain such an equation.

A characteristic of a torus is that its inner side has a perimeter shorter than the outer side. For this reason the toroidal geometry induces by flow conservation a velocity gradient between the two vertical interfaces at $r = r_{min}$ and $r = r_{max}$. This gradient has been measured by using the same tracers as previously. We found $\delta(V) = V(r_{min}) - V(r_{max}) \sim 60$ mm/s. Applying the Bernoulli relation to the streamline at the free surface of the torus, we obtain a pressure difference ΔP_v:

$$\Delta P_v = \frac{1}{2}\rho\delta(V^2) = \frac{1}{2}\rho\left(V^2(r_{min}) - V^2(r_{max})\right) \sim 2\,\text{Pa}. \tag{1}$$

On the other hand, the horizontal curvature of the inner perimeter $1/r_{min}$ creates a capillary pressure jump $\Delta P_\gamma = \gamma/r_{min} \sim 2$ Pa which tends to maintain this perimeter circular. Note that these two differences of pressure at the

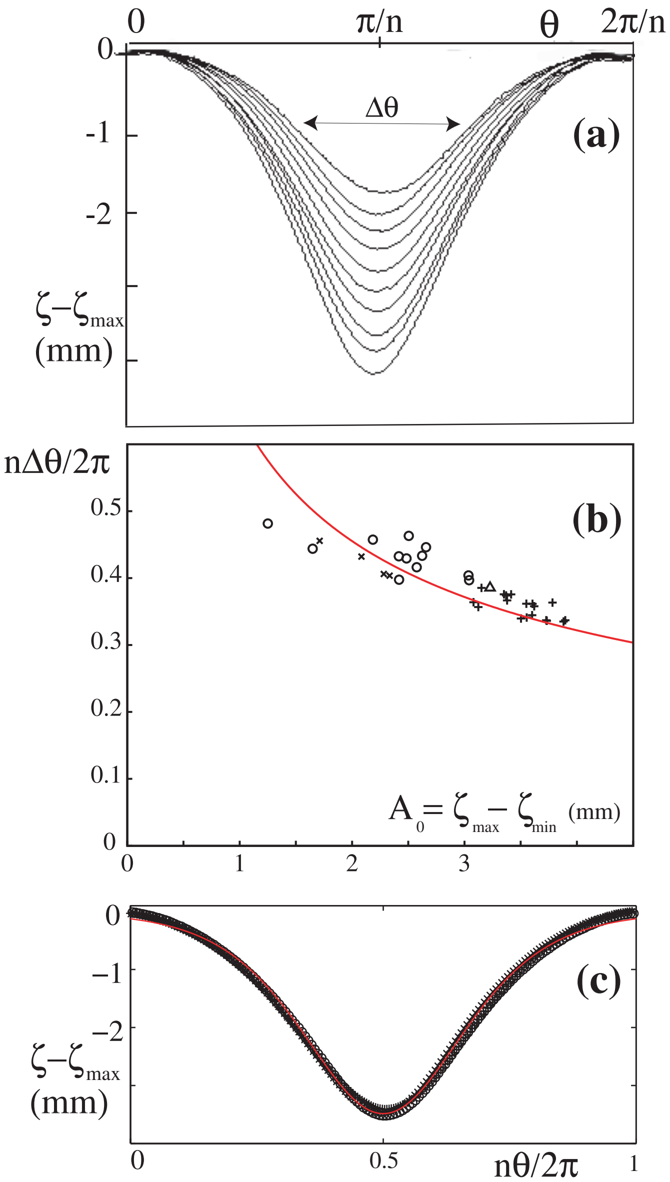

Fig. 4: (Colour on-line) (a) Various wave profiles for $n = 6$ obtained with substrate temperatures between 220 °C and 400 °C. $\Delta\theta$ is defined as the full width at half-maximum. (b) Evolution of $n\Delta\theta/2\pi$ as a function of the wave amplitude (for $n = 5$ (△), 6 (□), 7 (×), 8 (○), 9 (+)). The solid line is a fit with $n\Delta\theta/2\pi \propto A_0^{-1/2}$. (c) The superposition of normalized wave profiles for three different patterns ($n = 6$ (□), 7 (×), 8 (○)). The solid line is a fit using eq. (4).

interface are much smaller than both the hydrostatic pressure and the Laplace pressure coming from the vertical curvature of the interface. These main terms lead to the propagation of the waves when a deformation appears. However, we have observed the spontaneous growth of waves only when a poloidal flow exists inside the torus. Thus, we assume that an increase of Bernoulli pressure is induced by the reorganization of the flow near the corner, and pulls the interface towards the liquid phase. We assume the simplest possible expression for this Bernoulli extra pressure, *i.e.*, $\delta(V^2)/V^2 \sim -\beta_1 A/h + \beta_2 A^2$, where β_1 and β_2 are constants. The interface deformation can be limited by two different forces: the capillary pressure which limits the curvature of the inner torus perimeter, and the hydrostatic pressure which limits the wave amplitude. This latter can be written $\rho g \sin\alpha(\zeta - \zeta_{max})$ as the interface moves along the slope of the substrate. Roughly, we can admit that the wave profile of the internal perimeter is ruled by the following pressure equilibrium:

$$\frac{\rho}{2}\delta(V^2) = \frac{\gamma}{r_{min}}\frac{\partial^2\zeta}{\partial\theta^2} - \rho g \sin\alpha(\zeta - \zeta_{max}). \tag{2}$$

Defining an amplitude as $A = \zeta - \zeta_{max}$ and using the previous expression suggested for $\delta(V^2)$ yields

$$\frac{\partial^2 A}{\partial\theta^2} = -\beta_1(We - We_c)A + \beta_2 A^2, \tag{3}$$

in which $We = \rho r_{min} V^2/\gamma$ is the Weber number, and We_c is the critical value which depends on the reduced gravity term $g\sin\alpha$, and $\beta_2 A^2$ is a non-linear term of second order expected because of the lack of symmetry between $-A$ and $+A$. Finally, the gravity term is responsible for the threshold, and the wave profile is a local equilibrium between a dynamical difference of pressure created by the flow and a Laplace pressure. This amplitude equation is analogous to the Korteweg-de Vries one, describing a large family of non-linear waves [14]. Above a strong enough forcing ($We > We_c$), eq. (3) admits localized solutions:

$$A = -A_0 \,\mathrm{sech}^2 \frac{2\pi\theta}{n\Delta\theta}. \tag{4}$$

As one can verify in fig. 4(c), this shape describes the observed profile for our corners, and the scaling $n\Delta\theta/(2\pi) \sim A_0^{-1/2}$ implied by eq. (3) holds reasonably well with all the collected data.

A remarkable fact in this experiment is that the amplitude of these solitons is negative. This has never been observed at this scale, corresponding usually to gravity waves. The sole observation of a dark soliton in hydrodynamic waves is due to Falcon *et al.* for a wavelength ten times smaller than ours, in a capillary regime where it has been shown that the amplitude can become negative [15].

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REFERENCES

[1] Leidenfrost J. G., *De Aquae Communis Nonnullis Qualitatibus Tractatus* (Duisburg) 1756.

[2] Snoeijer J. H., Brunet P. and Eggers J., *Phys. Rev. E*, **79** (2009) 036307.

[3] Duchemin L., Lister J. R. and Lange U., *J. Fluid Mech.*, **533** (2005) 161.

[4] Biance A. L., Clanet C. and Quéré D., *Phys. Fluids*, **15** (2003) 1632.

[5] Tokugawa N. and Takaki R., *J. Phys. Soc. Jpn.*, **63** (1994) 1758.

[6] Linke H., Aleman B. J., Melling L. D., Taormina M. J., Francis M. J., Dow-Hygelund C. C., Narayanan V., Taylor R. P. and Stout A., *Phys. Rev. Lett.*, **96** (2006) 154502.

[7] Snezhko A., Jacob E. B. and Aranson I. S., *New J. Phys.*, **10** (2008) 043034.

[8] Vatistas G. H., *J. Fluid Mech.*, **217** (1990) 241.

[9] Jansson T. R. N., Haspang M. P., Jensen K. H., Hersen P. and Bohr T., *Phys. Rev. Lett.*, **96** (2006) 174502.

[10] Ellegaard C., Hansen A. E., Haaning A., Marcussen A., Bohr T., Hansen J. and Watanabe S., *Nature*, **392** (1998) 767.

[11] Bush J. W. M., Aristoff J. M. and Hosoi A. E., *J. Fluid Mech.*, **558** (2006) 33.

[12] Ellegaard C., Hansen A. E., Haaning A., Marcussen A., Bohr T., Hansen J. and Watanabe S., *Nonlinearity*, **12** (1999) 1.

[13] Bergman R., Tophoj L., Homan T., Hersen P., Andersen A. and Bohr T., *J. Fluid Mech.*, **679** (2011) 415.

[14] Korteweg D. J. and Vries G. D., *Philos. Mag. J. Sci.*, **39** (1895) 422.

[15] Falcon E., Laroche C. and Fauve S., *Phys. Rev. Lett.*, **89** (2002) 204501.