



HAL
open science

Rayleigh–Taylor instability in impact cratering experiments

Victor Lherm, Renaud Deguen, Thierry Alboussière, Maylis Landeau

► **To cite this version:**

Victor Lherm, Renaud Deguen, Thierry Alboussière, Maylis Landeau. Rayleigh–Taylor instability in impact cratering experiments. *Journal of Fluid Mechanics*, 2022, 937, pp.A20. 10.1017/jfm.2022.111 . hal-03623435

HAL Id: hal-03623435

<https://hal.science/hal-03623435v1>

Submitted on 16 May 2022

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

Banner appropriate to article type will appear here in typeset article

1 **Rayleigh-Taylor instability in impact cratering** 2 **experiments**

3 **V. Lherm^{1,2}†, R. Deguen³, T. Alboussière² and M. Landeau⁴**

4 ¹Department of Earth and Environmental Sciences, University of Rochester, 227 Hutchison Hall,
5 Rochester, NY 14627, USA

6 ²Univ Lyon, ENSL, UCBL, UJM, CNRS, LGL-TPE, F-69007 Lyon, France

7 ³Univ. Grenoble Alpes, Univ. Savoie Mont Blanc, CNRS, IRD, Univ. Gustave Eiffel, ISTerre, 38000
8 Grenoble, France

9 ⁴Université de Paris, Institut de Physique du Globe de Paris, CNRS, 75005 Paris, France

10 (Received xx; revised xx; accepted xx)

11 When a liquid drop strikes a deep pool of a target liquid, an impact crater opens while the
12 liquid of the drop decelerates and spreads on the surface of the crater. When the density of the
13 drop is larger than the target liquid, we observe mushroom-shaped instabilities growing at the
14 interface between the two liquids. We interpret this instability as a spherical Rayleigh-Taylor
15 instability due to the deceleration of the interface, which exceeds the ambient gravity. We
16 investigate experimentally the effect of the density contrast and the impact Froude number,
17 which measures the importance of the impactor kinetic energy to gravitational energy, on
18 the instability and the resulting mixing layer. Using backlighting and planar laser-induced
19 fluorescence methods, we obtain the position of the air-liquid interface, an estimate of
20 the instability wavelength and the thickness of the mixing layer. We derive a model for
21 the evolution of the crater radius from an energy conservation. We then show that the
22 observed dynamics of the mixing layer results from a competition between the geometrical
23 expansion of the crater, which tends to thin the layer and entrainment related to the instability,
24 which increases the layer thickness. The mixing caused by this instability has geophysical
25 implications for the impacts that formed terrestrial planets. Extrapolating our scalings to
26 planets, we estimate the mass of silicates that equilibrates with the metallic core of the
27 impacting bodies.

28 **MSC Codes** (*Optional*) Please enter your MSC Codes here

29 **1. Introduction**

30 Rayleigh-Taylor (RT) instability occurs at the interface between two fluids of different
31 densities, when the density and pressure gradients are in opposite directions. In a gravitational
32 field, an interface separating a dense fluid supported by a lighter one is unstable (Rayleigh
33 1899). In this situation, which we refer as static, the average position of the interface does not
34 vary with time. If the interface is accelerated in the direction from the lighter to the denser
35 fluid, the configuration is also unstable (Taylor 1950). In this situation, which we refer as

† Email address for correspondence: victor.lherm@ens-lyon.fr

36 dynamic, the average position of the interface varies over time. In both cases, infinitesimal
37 perturbations at the interface will grow over time, leading to the interpenetration of the fluids
38 and to the reduction of their combined potential energy.

39 The RT instability was first investigated at planar interfaces using theoretical, numerical and
40 experimental methods, both in the early-time linear (*e.g.* Emmons *et al.* 1960; Chandrasekhar
41 1961; Tryggvason 1988) and the subsequent non-linear regimes (*e.g.* Linden *et al.* 1994;
42 Dalziel *et al.* 1999; Dimonte 1999). However, various phenomena such as inertial confinement
43 fusion experiments (*e.g.* Lindl 1998; Thomas & Kares 2012), supernovae explosions (*e.g.*
44 Arnett *et al.* 1989; Schmidt 2006), detonation of explosive charges (*e.g.* Balakrishnan &
45 Menon 2011) and collapsing bubbles (*e.g.* Prosperetti 1977; Lin *et al.* 2002) involve RT
46 instabilities at spherical interfaces. The spherical configuration was initially investigated in
47 static and dynamic cases, regarding the early-time linear stability of spherical interfaces
48 between two inviscid fluids (Bell 1951; Plesset 1954; Mikaelian 1990). Viscosity effects
49 responsible for energy dissipation at small-scale were also investigated in both cases
50 (Chandrasekhar 1955; Prosperetti 1977; Mikaelian 2016). Turbulent mixing related to the
51 late-time non-linear RT instability dynamics was also investigated for spherical interfaces
52 (Youngs & Williams 2008; Thomas & Kares 2012; Lombardini *et al.* 2014).

53 The RT instability also enters the dynamics of drop impacts. Above a given kinetic energy,
54 the drop impact is followed by the formation of a liquid crown above the crater (*e.g.* Rein
55 1993). The interface between the liquid of the crown and the air rapidly decelerates, which
56 leads to the formation of fluid fingers in part interpreted as an RT instability (*e.g.* Allen 1975;
57 Krechetnikov & Homsy 2009; Agbaglah *et al.* 2013).

58 In our experiment, we investigate another instability that develops under the water surface
59 when a denser-than-water drop impacts a deep liquid pool of water. Figure 1 shows a snapshot
60 from one such experiment, at a time when the impact of the drop has produced a sizeable
61 crater. We see on this picture that the liquid from the drop, which covers the crater floor, has
62 developed mushroom-shaped structures penetrating radially into the water pool. We interpret
63 these structures as a spherical RT instability associated with the deceleration of the crater
64 floor. The dynamics of the RT instability depends crucially on the acceleration history of the
65 interface (Mikaelian 1990; Dimonte & Schneider 2000).

66 In the case of a drop impact, this acceleration is dictated by the dynamics of the crater,
67 which depends on the impact parameters, *i.e.* drop radius, impact velocity, ambient gravity
68 and physical properties of the fluids such as surface tension, density and viscosity. Depending
69 on these impact parameters, various impact regimes such as floating, bouncing, coalescence
70 and splashing may occur (*e.g.* Rein 1993). In this work, we focus on the splashing regime.
71 Since the pioneering experiments of Worthington (1895), the splashing regime has been
72 extensively investigated, regarding in particular the scaling of the maximum crater size (Engel
73 1966; Leng 2001), the time evolution of the transient crater (Engel 1967; Morton *et al.* 2000;
74 Bisighini *et al.* 2010), the evolution and fragmentation of the fluid crown (Okawa *et al.* 2006;
75 Zhang *et al.* 2010), the formation and fragmentation of the central peak (Fedorchenko &
76 Wang 2004; Ray *et al.* 2015) and underwater acoustic properties of the impact (Prosperetti
77 & Oguz 1993). Furthermore, the effects of immiscibility (Lhuissier *et al.* 2013; Jain *et al.*
78 2019), viscoplasticity (Jalaal *et al.* 2019), impact angles (Okawa *et al.* 2006; Gielen *et al.*
79 2017) and thickness of the target layer (Berberović *et al.* 2009) on impact dynamics have
80 been examined. Based on these experiments, several models of the crater size evolution and
81 its related acceleration history were developed, using energy conservation (*e.g.* Engel 1966,
82 1967; Leng 2001) or momentum conservation in an irrotational flow (*e.g.* Bisighini *et al.*
83 2010).

84 Besides providing an example of an RT instability at a spherical interface, this drop impact
85 instability and the mixing related to it have geophysical implications. Terrestrial planets such

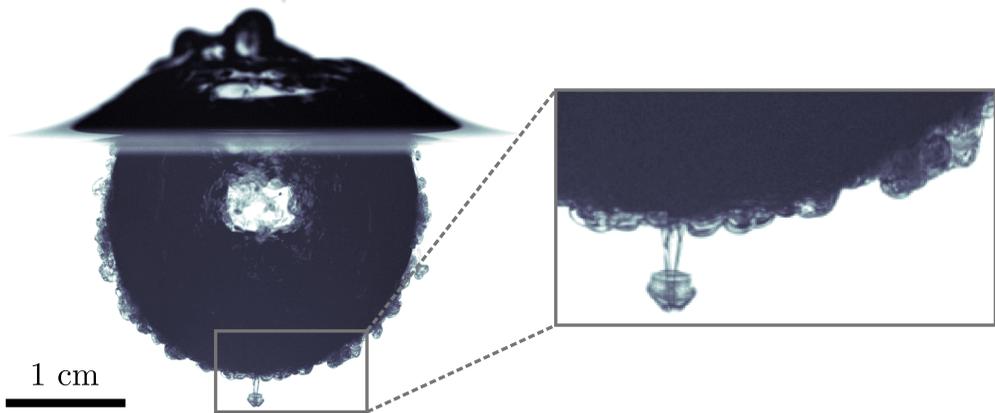


Figure 1: Crater produced by the vertical impact of a liquid drop onto a less dense liquid pool. A spherical RT instability develops around the crater when it decelerates, which results in mushroom-shaped plumes growing radially outward.

86 as the Earth formed 4.5×10^9 years ago by the successive accretion of increasingly massive
 87 bodies composed mainly of silicates and iron (Chambers 2010), the last giant impact being
 88 probably responsible for the formation of the Moon (Canup 2012; Cuk & Stewart 2012).
 89 During this process, planetary materials are heated by the kinetic energy released during the
 90 impacts, the reduction of gravitational potential energy as the metal of the impactors migrates
 91 toward the core and the decay of radioactive isotopes (Rubie *et al.* 2015). This energy supply
 92 contributes to the production of deep magma oceans (Solomatov 2015). In addition, accretion
 93 models show that most of the Earth mass was accreted from differentiated bodies (Kleine
 94 *et al.* 2002; Scherstén *et al.* 2006), *i.e.* with a separate core and mantle (in contrast, it should
 95 be noted that we use drops made of a single liquid in our experiments). Both the impacting
 96 body and the planetary surface are melted by the shock waves produced by the impact,
 97 releasing the liquid metal core of the impactor into a fully molten magma ocean (Tonks
 98 & Melosh 1993). This results in a situation where the metal core of an impactor strikes a
 99 less dense silicate magma ocean. A spherical RT instability can therefore develop during
 100 crater opening, producing mixing that contributes to the thermal and chemical equilibration
 101 between the metal core of the impactors and the silicates of the magma ocean.

102 The current dynamics of the Earth is partly inherited from its concomitant accretion and
 103 differentiation. Heat partitioning and chemical fractionation between the mantle and the core
 104 depend on the physical processes involved during differentiation (Stevenson 1990; Wood *et al.*
 105 2006), which includes in particular equilibration and dispersion occurring during planetary
 106 impacts (Canup 2004; Cuk & Stewart 2012; Kendall & Melosh 2016; Nakajima *et al.* 2021;
 107 Landeau *et al.* 2021). Heat partitioning sets the initial temperature contrast between the
 108 mantle and the core. It crucially determines the early thermal and magnetic evolution of the
 109 planet, in particular the formation and evolution of magma oceans (Labrosse *et al.* 2007; Sun
 110 *et al.* 2018) and the existence of an early dynamo (Williams & Nimmo 2004; Monteux *et al.*
 111 2011; Badro *et al.* 2018). Chemical fractionation has also major geodynamical implications,
 112 such as the nature and abundance of radioactive and light elements in the core (Corgne *et al.*
 113 2007; Siebert *et al.* 2012; Badro *et al.* 2015; Fischer *et al.* 2015). Geochemical data such as
 114 isotopic ratios and partitioning coefficients between metal and silicates provide constraints
 115 on the timing of accretion and physical conditions of core formation in terrestrial planets
 116 (Li & Agee 1996; Kleine *et al.* 2002; Righter 2011; Siebert *et al.* 2011). However, their

117 interpretation depends on the degree of chemical equilibration between the metal of the
 118 impactors' core and the magma ocean (Rudge *et al.* 2010; Rubie *et al.* 2011). Consequently,
 119 an estimate of the equilibration produced by the spherical RT instability during the impact
 120 is required in order to properly interpret geochemical data. In this paper, we examine the
 121 spherical RT instability produced during an impact using fluid dynamics experiments.

122 The layout of this paper is as follows. In section 2, we introduce the set of non-dimensional
 123 numbers used in this study and present the experimental procedure. In section 3, we describe
 124 the phenomenology of the cratering process and of the RT instability. In section 4, we obtain
 125 a model for the crater radius evolution from energy conservation. In section 5, we use the
 126 acceleration history of the cratering process to model the evolution of the thickness of the
 127 mixing layer resulting from the RT instability. In section 6, we finally apply this model to the
 128 differentiation of terrestrial planets.

129 2. Impact cratering experiments

130 2.1. Non-dimensional numbers

131 The impact dynamics of a liquid drop released above a deep liquid pool with a different
 132 density and viscosity depends on its impact velocity U_i and radius R_i , the densities ρ_1 and
 133 ρ_2 of the drop and the pool, the dynamic viscosities μ_1 and μ_2 of the drop and the pool, the
 134 surface tension at the air-liquid interface σ and the acceleration of gravity g . Since these
 135 eight parameters contain three fundamental units, the Vaschy–Buckingham theorem dictates
 136 that the impact dynamics depends on a set of five independent dimensionless numbers. We
 137 choose the following set:

$$138 \quad Fr = \frac{U_i^2}{gR_i}, \quad We = \frac{\rho_1 U_i^2 R_i}{\sigma}, \quad Re = \frac{\rho_2 U_i R_i}{\mu_2}, \quad \rho_1/\rho_2, \quad \mu_1/\mu_2. \quad (2.1)$$

139 The Froude number Fr is a measure of the relative importance of impactor inertia and gravity
 140 forces. It can also be interpreted as the ratio of the kinetic energy $\rho_1 R_i^3 U_i^2$ of the impactor to
 141 its gravitational potential energy $\rho_1 g R_i^4$ just before impact. The Weber number We compares
 142 the impactor inertia and interfacial tension at the air-liquid interface. The Reynolds number
 143 Re is the ratio between inertial and viscous forces. The ratios ρ_1/ρ_2 and μ_1/μ_2 compare,
 144 respectively, the density and the dynamic viscosity of the drop and the pool. Although the
 145 Atwood number $A = (\rho_1 - \rho_2)/(\rho_1 + \rho_2)$ is used widely in the context of RT instabilities,
 146 we use instead the density ratio because it appears more naturally in the equations governing
 147 the crater size evolution (see section 4). Since surface tension depends on salt concentration,
 148 a surface tension ratio between the drop and the pool is also involved. However, we ignore
 149 this parameter because the Weber number is much larger than unity and the surface tension
 150 of the drop varies by less than 20% compared with the pool. We also use a modified Froude
 151 number and a Bond number,

$$152 \quad Fr^* = \frac{\rho_1 U_i^2}{\rho_2 g R_i}, \quad Bo = \frac{\rho_2 g R_i^2}{\sigma}, \quad (2.2)$$

153 which are, respectively, the ratio of the kinetic energy of the impactor $\rho_1 R_i^3 U_i^2$ to the change
 154 of potential gravitational energy $\rho_2 g R_i^4$ associated with the opening of a crater of size R_i
 155 and the ratio of buoyancy forces to interfacial tension at the air-liquid interface. The Bond,
 156 Weber and modified Froude numbers are related through $Bo = We/Fr^*$. In the following,
 157 time, lengths, velocities and accelerations are made dimensionless using the drop radius and
 158 the impact velocity, *i.e.* using, respectively, R_i/U_i , R_i , U_i and U_i^2/R_i . These dimensionless

159 quantities are denoted with a superscript star. For example, we use a dimensionless time
 160 $t^* = t/(R_i/U_i)$.

161 Planetary impacts are generally classified between a strength regime and a gravity regime,
 162 depending on the resistance to deformation of the target material (*e.g.* Melosh 1989). When
 163 the yield stress of the solid planetary surface, or the viscous stress of the magma ocean, are
 164 negligible in comparison with the impact-induced stress, planetary impacts are in the gravity
 165 regime. This typically happens when the impactor is larger than a few kilometres (Holsapple
 166 1993). In this case, the Reynolds number, the Weber number and the Bond number are
 167 expected to be larger than 10^{10} , due to kilometric impactors with impact velocities of the
 168 order of 10 km.s^{-1} . This leads to extremely turbulent impact conditions where surface tension
 169 effects are negligible. Assuming that the impact velocity is close to the escape velocity of the
 170 target planet, *i.e.* $U_i \approx \sqrt{gR_t}$ with R_t the radius of the target, the modified Froude number
 171 scales as

$$172 \quad Fr^* \approx \frac{\rho_1 R_t}{\rho_2 R_i}. \quad (2.3)$$

173 Fr^* is then of the order of 1 for impactors comparable in size with the target, but increases
 174 by several orders of magnitude for small colliding bodies, *e.g.* the Froude number is about
 175 10^4 for a 1 km radius body impacting an Earth-sized planet. From the perspective of the
 176 impactor, small impacts are thus dynamically more extreme than giant impacts. Since the
 177 metal core of the impactor is mainly composed of iron, the density ratio and the viscosity
 178 ratio are, respectively, expected to be of the order of $\rho_1/\rho_2 \approx 2$ and $\mu_1/\mu_2 \approx 0.1$ (Solomatov
 179 2015).

180 Table 1 compares the value of the dimensionless parameters in our drop impact experiments
 181 and in planetary impacts. Since experimental Reynolds numbers ($Re \gtrsim 2500$) and Weber
 182 numbers ($We \gtrsim 60$) are large, viscosity and surface tension are mostly negligible during
 183 crater opening (*e.g.* Olevson 1969; Macklin & Metaxas 1976). Although Re and We are much
 184 larger during planetary impacts than in our experiments, this means that the cratering process
 185 is governed by inertia and buoyancy forces, in both our experiments and planetary impacts.
 186 We thus focus on the gravity regime (*e.g.* Melosh 1989), where the cratering dynamics
 187 depends mainly on two dimensionless parameters, the Froude number Fr and the density
 188 ratio ρ_1/ρ_2 .

189 In order to characterise the crater geometry and the RT instability following the impact,
 190 we vary the drop radius, drop density and impact velocity. In our experiments, we obtain
 191 Froude numbers and modified Froude numbers larger than unity, in the range $Fr \approx$
 192 $60 - 1200$ and $Fr^* \approx 60 - 2100$, respectively. We focus on five density ratios $\rho_1/\rho_2 \approx$
 193 $\{1.0, 1.2, 1.4, 1.6, 1.8\}$, in comparison with a reference case without density contrast. We
 194 have also conducted a few experiments at $\rho_1/\rho_2 \approx 0.8$ using ethanol in the drop. The
 195 upper limit of our experimental density ratios is close to the density ratio expected during a
 196 planetary impact.

197 *2.2. Experimental set-up*

198 *2.2.1. Drop production, fluids and cameras*

199 In these experiments, we release a liquid drop in the air above a deep liquid pool contained
 200 in a $16 \times 16 \times 30$ cm glass tank (figure 2). The pool level is set at the top of the tank. The
 201 aim is to minimise the thickness of the meniscus on the sides of the tank in order to obtain
 202 an image of the crater all the way to the surface.

203 We generate the drop using a needle supplied with fluid by a syringe driver at a slow
 204 and steady pace. When the weight of the drop exceeds the surface tension forces, the drop
 205 comes off. We use a metallic needle with an inner diameter of 1.6 mm and a nylon plastic

Dimensionless number	Experiments	Planetary impacts
Fr	60 – 1200	$1 - 10^5$
Fr^*	60 – 2100	$1 - 10^5$
We	60 – 1300	$\geq 10^{14}$
Bo	0.7 – 1	$\geq 10^{10}$
Re	2500 – 13500	$\geq 10^{11}$
ρ_1/ρ_2	1 – 1.8	2
A	0 – 0.29	0.33
μ_1/μ_2	0.9 – 1.2	0.1

Table 1: Typical values of the main dimensionless parameters (equations 2.1 and 2.2) in the experiments and typical planetary impacts. For planetary impacts, dimensionless numbers are estimated with a density of 4000 kg.m^{-3} for molten silicates and of 8000 kg.m^{-3} for molten metal, a dynamic viscosity of 0.1 Pa.s for molten silicates and of 0.01 Pa.s for molten metal (Solomatov 2015). Surface tension between air and molten silicates and between air and molten metal, are taken to be 0.3 J.m^{-2} (Taniguchi 1988) and 1.8 J.m^{-2} (Wille *et al.* 2002), respectively. We assume the impact velocity to be one to three times the escape velocity (Agnor *et al.* 1999; Agnor & Asphaug 2004). We assume the impactor-to-target radius to be in the range $10^{-4} - 1$.

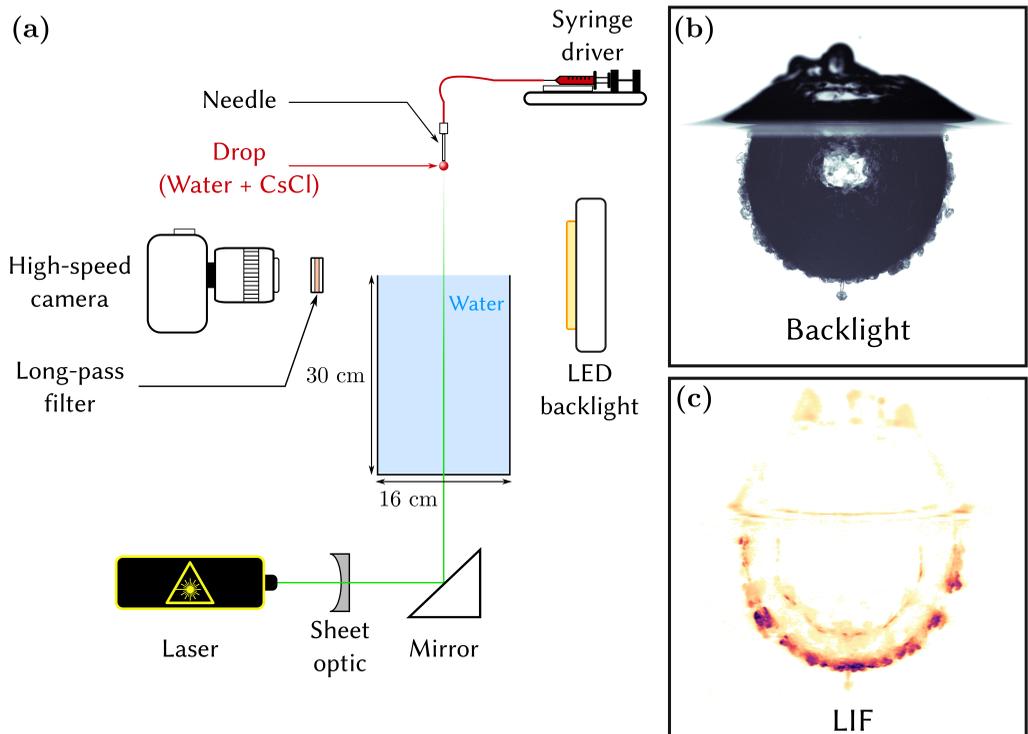


Figure 2: (a) Schematic view of the experimental set-up, including backlight and LIF configuration set-up. (b) Snapshot obtained using the backlight configuration. (c) Snapshot obtained using the LIF configuration.

206 needle with an inner diameter of 4.7 mm, generating two series of drop size with typical
 207 radius in the range 1.7 – 2.0 mm and 2.3 – 2.7 mm, respectively. We estimate the drop size,
 208 which depends on the drop density, for each experiment based on a calibration using mass
 209 measurements of dozens of drops, independent density measurement and assuming the drop
 210 is spherical. We validate this method using high-speed pictures of the drop prior to impact
 211 where we can directly measure the drop radius.

212 Typical impact velocities are in the range 1 – 5 m.s⁻¹. We calculate the impact velocity for
 213 each experiment using a calibrated free-fall model for the drop, including a quadratic drag.
 214 We also validate this method using high-speed pictures of the drop prior to impact where we
 215 can directly measure the drop velocity.

216 We use an aqueous solution of caesium chloride CsCl ($\rho_1 = 998 - 1800 \text{ kg.m}^{-3}$, $\mu_1 =$
 217 $0.9 \times 10^{-3} - 1.2 \times 10^{-3} \text{ Pa.s}$) in the drop and water ($\rho_2 = 998 \text{ kg.m}^{-3}$, $\mu_2 = 10^{-3} \text{ Pa.s}$) in the
 218 pool. In a few experiments, we also use pure ethanol ($\rho_1 = 790 \text{ kg.m}^{-3}$, $\mu_1 = 1.2 \times 10^{-3} \text{ Pa.s}$)
 219 in the drop. The density is measured for each experiment using an Anton Paar DMA 35 Basic
 220 densitometer. Since the typical measurement errors on density are less than 0.1%, we neglect
 221 errors on density. We obtain viscosities using data from Haynes (2016). We neglect errors
 222 on viscosity since the typical error is less than 0.01%. The surface tension at the air-water
 223 interface is $\sigma = 72.8 \pm 0.4 \text{ mJ.m}^{-2}$ (Haynes 2016).

224 We obtain errors on the velocity and radius of the impacting drop from the variability in
 225 mass measurements and from error propagation in the velocity model, respectively. We then
 226 propagate uncertainties on fluid properties and impact parameters to uncertainties on the
 227 dimensionless numbers (equations 2.1 and 2.2).

228 We use two imaging configurations, backlight and laser-induced fluorescence (LIF)
 229 configurations, most suited to determine the crater shape and characterise the mixing layer
 230 thickness, respectively. In both configurations, we position the camera at the same height
 231 as the water surface. We record images at 1400 Hz with a 2560×1600 pixels resolution
 232 ($34 \mu\text{m.pixel}^{-1}$) and a 12 bits dynamic range, using a high-speed Phantom VEO 640L camera
 233 and a Tokina AT-X M100 PRO D Macro lens.

234 2.2.2. Backlighting

235 In the backlight configuration (figure 3a), we measure the evolution of the mean crater radius.
 236 The crater is illuminated from behind by a light emitting diode (LED) backlight panel and it
 237 appears dark owing to refraction of light at the air-water interface. Image processing involves
 238 spatial calibration using a sight, background removal, intensity threshold, image binarisation
 239 and allows us to determine the crater boundary.

240 We then fit the experimental crater boundary radius $R_\theta^{\text{exp}}(\theta, t)$, which depends on the polar
 241 angle θ and time t , using a set of shifted Legendre polynomials \tilde{P}_l up to degree $l = 5$, as
 242 follows:

$$243 \quad R_\theta^{\text{exp}}(\theta, t) = \sum_{l=0}^5 a_l^{\text{exp}} \tilde{P}_l(\cos \theta), \quad (2.4)$$

244 where a_l^{exp} are the experimental fitted coefficients. The shifted Legendre polynomials are
 245 an affine transformation of the standard Legendre polynomials $\tilde{P}_l(x) = P_l(2x - 1)$ and are
 246 orthogonal on $[0, 1]$, *i.e.* on a half-space. We obtain the experimental mean crater radius
 247 from the $l = 0$ coefficient, *i.e.* $R^{\text{exp}} = a_0^{\text{exp}}$.

248 Uncertainties are dominated by the extrinsic variability between experiments in the same
 249 configuration. We repeat each experiment at least four times consecutively in order to estimate
 250 uncertainties on dimensionless parameters and target quantities. This allows us to include
 251 uncertainties resulting from reflections and refraction at the crater boundary. In comparison,

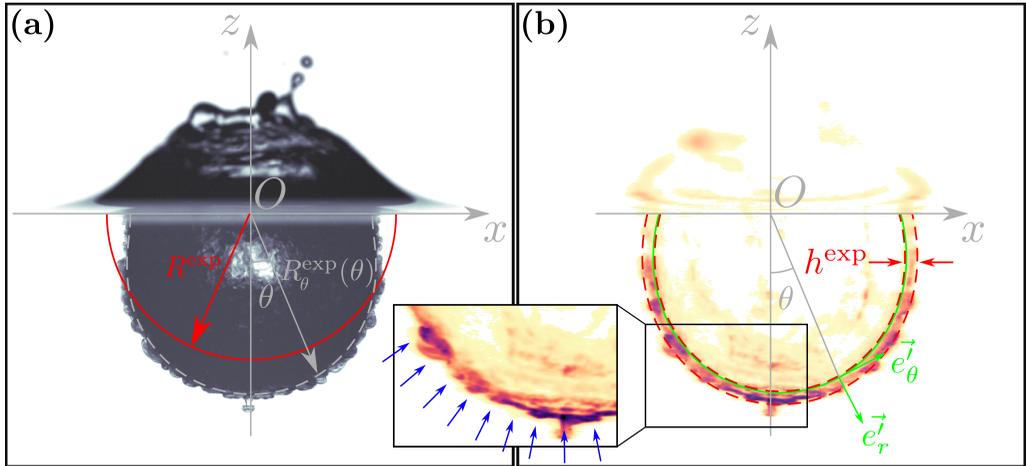


Figure 3: (a) Detection of the crater boundary. The dashed line corresponds to the fitted crater boundary using a set of shifted Legendre polynomials up to degree $l = 5$, the degree $l = 0$ gives the solid line corresponding to the mean crater radius. (b) Detection of the mixing layer. The solid line corresponds to the fitted crater boundary using a set of Legendre polynomials up to degree $l = 5$. The dashed lines correspond to the weighted average of the mixing layer boundaries, calculated using the second moment of the LIF intensity about the mean position of the layer. Blue arrows indicate the position of the plumes produced by the instability.

252 the intrinsic uncertainties related to the spatial resolution of the camera, spatial calibration
 253 and image processing, are small. Spatial resolution of the camera, $30 \text{ pixel} \cdot \text{mm}^{-1}$, is adequate
 254 given the size of the target, allowing this uncertainty to be neglected. We also neglect spatial
 255 calibration errors, typically around 0.2 pixel. Finally, given the camera resolution and dynamic
 256 range, we obtain a good contrast on the crater and the impacting drop, which allows us to
 257 neglect errors related to image processing.

258 2.2.3. LIF

259 In the LIF configuration (figure 3b), we measure the thickness of the mixing layer and the
 260 number of plumes produced by the RT instability. The fluorescent dye (Rhodamine 6G,
 261 $1.5 \text{ mg} \cdot \text{l}^{-1}$) contained in the fluid of the drop is excited by a vertical laser sheet (532 nm).
 262 The fluorescent dye then re-emits light between 500 nm and 700 nm. This emission signal is
 263 then recorded by the camera and isolated from the laser signal with a long-pass filter (> 540
 264 nm). We generate the laser sheet using a 10 W Nd:YAG continuous laser in combination
 265 with a divergent cylindrical lens and a telescope, producing a 1 mm thick sheet. We divert
 266 the laser sheet vertically using a 45° mirror beneath the tank. In order to isolate the mixing
 267 layer, we process images with spatial calibration using a sight, background removal and
 268 laser sheet corrections, removing sheet inhomogeneities. We then filter and remove artefacts
 269 due to reflections on the surface and on the air-water interface. In particular, we remove the
 270 internal reflection of the mixing layer (*e.g.* figure 2c). We eventually obtain a measured LIF
 271 intensity field from which we can identify the mixing layer. Scalar diffusion of the dense
 272 liquid is negligible since the diffusion length during crater opening (typically 50 ms) is
 273 approximately $10 \mu\text{m} = 0.3 \text{ pixel}$. During crater opening, diffusion of Rhodamine 6G is even
 274 more negligible.

275 As for the backlight configuration, we fit the position of the experimental crater boundary
 276 $R_\theta^{\text{exp}}(\theta, t)$, which corresponds to the inner boundary of the mixing layer, using a set of shifted
 277 Legendre polynomials up to degree $l = 5$. We then define a local frame of reference $(\vec{e}_r', \vec{e}_\theta')$,

278 where \mathbf{e}'_r is normal to the fitted crater boundary and \mathbf{e}'_θ is tangent to it. For each polar position
 279 θ about the crater boundary, we calculate the local mean position of the mixing layer $\langle r' \rangle$,
 280 using the position of the pixels in the local frame of reference (r', θ) and the corresponding
 281 LIF intensity field I :

$$282 \quad \langle r' \rangle(\theta) = \frac{\int r' I(r', \theta) dr'}{\int I(r', \theta) dr'}. \quad (2.5)$$

283 We then calculate the local standard deviation $\sigma_{r'}$ about the local mean position of the mixing
 284 layer,

$$285 \quad \sigma_{r'}(\theta) = \sqrt{\frac{\int [r' - \langle r' \rangle(\theta)]^2 I(r', \theta) dr'}{\int I(r', \theta) dr'}}. \quad (2.6)$$

286 We eventually obtain the experimental mixing layer thickness h^{exp} with

$$287 \quad h^{\text{exp}} = \frac{\int_{-\pi/2}^{\pi/2} 2\sigma_{r'}(R_\theta^{\text{exp}})^2 |\sin \theta| w(\theta) d\theta}{\int_{-\pi/2}^{\pi/2} (R_\theta^{\text{exp}})^2 |\sin \theta| w(\theta) d\theta}, \quad (2.7)$$

288 using a weighted average where $w = 1/[1 + \exp\{k(|\theta| - \theta_0)\}]$ is a symmetric logistic weight
 289 function whose steepness is $k = 30$ and whose sigmoid midpoint is $\theta_0 = \pi/3$. The logistic
 290 function allows us to give more weight to the bottom of the crater, between $\theta = 0$ and $\theta = \theta_0$
 291 and less to the top of the crater, close to $\theta = \pm\pi/2$. The use of such a weight function is
 292 motivated by the polar dependency of the LIF signal quality. The crater is indeed illuminated
 293 from below, so that the laser sheet undergoes absorption (due to the fluorescent dye) and
 294 refraction (due to small scale variations of the index of refraction) as it goes through the
 295 mixing layer. Imaging of the mixing layer may also be perturbed directly by the air-water
 296 interface, causing reflection of the laser sheet. Consequently, close to the surface, *i.e.* at
 297 $\theta = \pm\pi/2$, imaging of the mixing layer may undergo significant perturbations, leading to a
 298 poor estimate of its extent. All these effects are amplified when the crater slightly deviates
 299 from a hemispherical shape. When the drop is denser than the pool, the crater is stretched
 300 downward, leading to an ellipsoidal crater centred below the surface of the pool. At low
 301 Weber and Froude numbers, the crater is also deformed by the propagation of a capillary
 302 wave which is superimposed on the density effects. The path of the laser sheet through the
 303 mixing layer thus increases and is more likely to cross the air-water interface.

304 We also count manually the number of plumes produced by the RT instability for each
 305 experiment at the same dimensionless time $t^* \sim 10$ (figure 3b, inset). This time was chosen
 306 from visual inspections, which suggest that at this time the plumes have not yet interacted
 307 significantly with each other, which is relevant since this number of plumes is to be compared
 308 with a scaling of the early-time instability wavelength (appendix A).

309 As in the backlight configuration, uncertainties are dominated by the extrinsic variability
 310 between experiments in the same configuration and hence, we repeat each experiment
 311 at least four times consecutively. We neglect intrinsic uncertainties related to the spatial
 312 resolution of the camera, the spatial calibration and image processing, since as in the backlight
 313 configuration they are still comparatively small.

314 3. Phenomenology

315 In our experiments, the development of the RT instability is governed by the crater
 316 evolution following the impact and particularly by its acceleration history. Using both
 317 backlight and LIF configurations, we characterise the air-water interface evolution and the

318 mixing layer evolution. From these measurements, we describe the observed RT instability.
 319 We base our description on two typical experiments, with and without density contrast,
 320 with similar evolution of the crater radius, in the backlight (figure 4a,c) and the LIF
 321 configurations (figure 4b,d). Experiments without density contrast are also available as
 322 animations in supplementary movie 1 (backlight configuration) and supplementary movie 2
 323 (LIF configuration), while experiments with density contrast are available in supplementary
 324 movie 3 (backlight configuration) and supplementary movie 4 (LIF configuration) – available
 325 at <https://doi.org/10.1017/jfm.2022.111>.

326

3.1. Crater geometry

327

328

329

330

331

332

333

334

335

336

337

338

339

340

341

342

343

344

345

346

347

The impact of the drop first causes the formation of an impact crater with a flat bottom due to the spreading of the drop (figure 4a, i). The liquid of the drop is deformed and accumulated on the crater floor (figure 4b, i) on a time scale $t^* \approx 2 - 3$, akin to previous results (*e.g.* Engel 1961, 1962). As the crater grows, the cavity becomes hemispherical, on a time scale $t^* \approx 10$, as a result of the overpressure produced at the contact point between the impacting drop and the surface (figure 4a, ii,iii). The liquid of the drop then spreads over the crater sides toward the surface, producing a layer with an approximately uniform thickness at the surface of the crater (figure 4b, iii,iv).

The impact also produces a cylindrical fluid crown (Fedorchenko & Wang 2004) (figure 4a, i-iv), along with a surface wave propagating radially outward from the crater on the horizontal surface (Leng 2001). As can be seen from the LIF intensity field (figure 4b, i-iv), the fluid of the drop mostly accumulates on the surface of the crater, leaving the crown mainly composed of fluid from the pool. As soon as the crown decelerates, the cylindrical sheet produces liquid ligaments around the crown rim, which eventually fragment into drops (figure 4a, ii,iii) (Krechtnikov & Homsy 2009; Zhang *et al.* 2010; Agbaglah *et al.* 2013).

When the crater reaches its maximum size, the crater is almost hemispherical (figure 4a, iv). When $\rho_1/\rho_2 > 1$, the crater cavity is entrained downward, which stretches vertically the crater (figure 4c, iv). After reaching its maximum size, the crater starts to collapse due to buoyancy forces (figure 4a, v). The resulting converging flow leads to the formation of a thick upward jet which, in view of the LIF intensity field (figure 4b, vi), appears to be made mostly of liquid from the drop.

348

3.2. RT instability

349

350

351

352

353

354

355

356

The presence of an instability at the interface between the drop and target fluids is particularly clear when comparing experiments without and with a density contrast between the two liquids (figure 4). We will also base our description on figure 5, which shows composite images obtained from sequences of LIF images from these two experiments. They are built by extracting from each LIF image a column of pixels centred on the crater vertical axis of symmetry, before juxtaposing them to show the LIF intensity as a function of vertical position and time. This provides a qualitative picture of the evolution of the radial dispersion of the liquid drop.

The early-time evolution of the drop liquid is similar in both experiments up to $t^* \sim 10$. The liquid from the drop spreads over the crater floor and forms a thin layer of approximately uniform thickness (figure 4 (all subpanels i,ii) and figure 5). In the experiment with no density contrast (figure 4a,b and figure 5a), the thickness of the layer keeps decreasing with time as the crater grows. This is a direct consequence of the crater expansion, the liquid of the drop being redistributed over an increasingly large surface area.

363

364

The behaviour of the layer is markedly different when the drop is denser than the pool. Figure 4c,d and figure 5b show that the thickness of the mixing layer initially decreases up

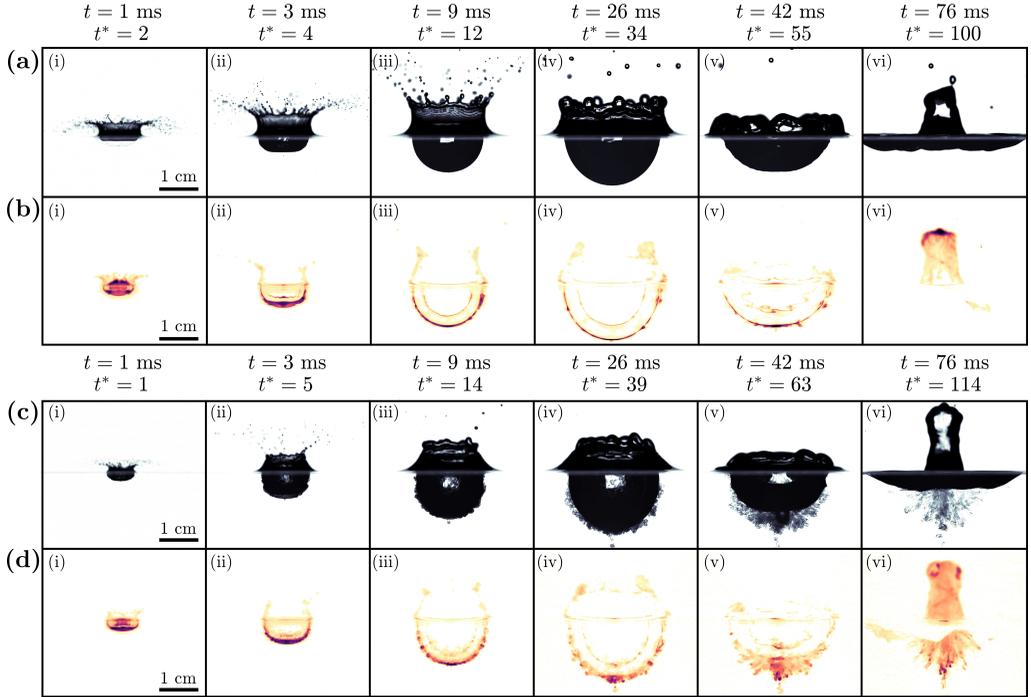


Figure 4: Liquid drop impact onto a deep liquid pool without density contrast ($\rho_1/\rho_2 = 1$) in the backlight ($Fr = 483$) (a) and the LIF ($Fr = 481$) (b) configuration. Liquid drop impact onto a deep liquid pool with density contrast ($\rho_1/\rho_2 = 1.8$) in the backlight ($Fr = 542$) (c) and the LIF ($Fr = 572$) (d) configuration. The liquid drop impact sequences are also available as animations in supplementary movies 1 (a), 2 (b), 3 (c) and 4 (d).

365 to $t^* \approx 10$, but it then increases over time while small scale corrugations of the interface
 366 evolve into radially growing plumes. The fine scale structure of the layer as seen from the
 367 LIF intensity field (figure 4d, iii,iv) is reminiscent of the structure of mixing layers observed
 368 in RT experiments in planar configurations (*e.g.* Dalziel *et al.* 1999).

369 These observations can be rationalised as follows. During the crater opening phase, the
 370 rate at which the crater grows gradually decreases with time, which results in a deceleration
 371 of the boundary between the mixing layer and the surrounding liquid. This situation is known
 372 to be unstable with respect to the RT instability when the liquid from the drop is heavier than
 373 the liquid from the pool (Taylor 1950). Measurements of the crater acceleration as a function
 374 of time (figure 7c of §4.2) shows that the deceleration of the crater boundary can be more
 375 than 10 times larger than the acceleration of gravity, which explains why the dense liquid
 376 plumes grow radially outward rather than in the vertical direction. The growth of the RT
 377 instability tends to increase the thickness of the mixing layer. However, the expansion of the
 378 crater spreads the dyed liquid over an increasing large surface area, which tends to make the
 379 layer thinner. The competition between these two effects explains why the thickness of the
 380 mixing layer first decreases as the crater expands and then increases when the RT instability
 381 dominates.

382 The instability first goes through a stage where the perturbations of the interface are small in
 383 comparison with the radius of the crater and the wavelength of the instability (figure 4d, i,ii).
 384 The growth of the RT instability competes with the geometrical expansion of the crater, which
 385 stretches the instabilities, hence reducing the amplitude of the perturbations and increasing

12

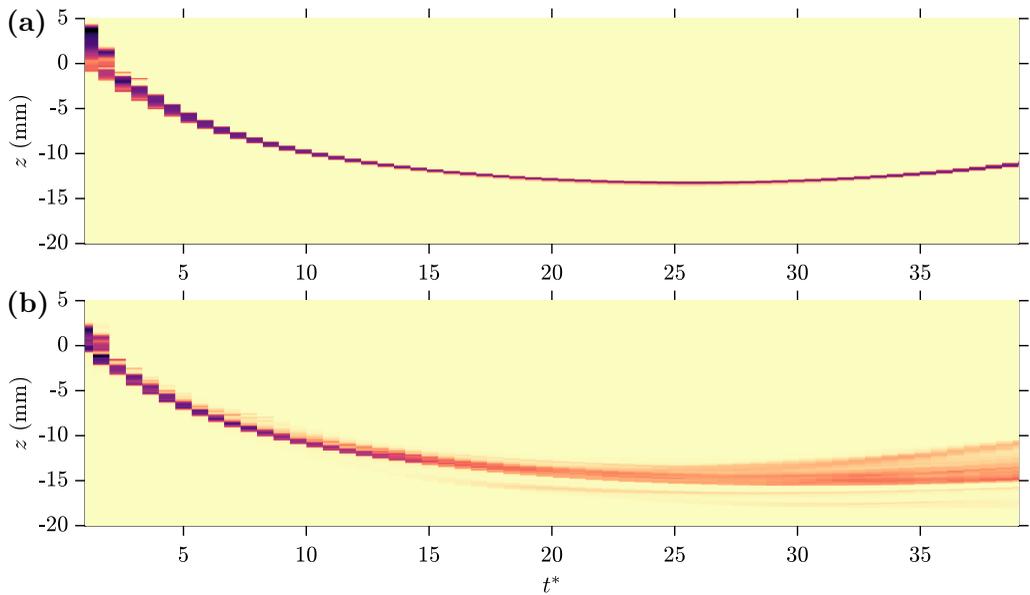


Figure 5: Composite images showing the LIF intensity as a function of vertical position and time, on the crater axis of symmetry ($\theta = 0$). The two panels have been built from experiments without (a) ($\rho_1/\rho_2 = 1.0$) and with (b) ($\rho_1/\rho_2 = 1.8$) a density contrast between the drop and the pool. The surface of the pool prior to impact is at $z = 0$.

386 their wavelength. This initial stage is dominated by the fast vigorous crater expansion, which
 387 contributes to decrease the mixing layer thickness by spreading the volume of the layer over
 388 an increasingly large surface area. Considering the resolution of our experiments, we are not
 389 able to observe the small initial perturbations that could exist in spite of the vigorous crater
 390 expansion. We observe that the amplitude of the perturbations eventually reaches the same
 391 size as the instability wavelength at $t^* \simeq 10$.

392 The two fluids involved being miscible, *i.e.* surface tension is zero, all wavenumbers
 393 are expected to be unstable with respect to the RT instability (Chandrasekhar 1955). In
 394 addition, owing to larger velocity gradients at large wavenumbers, viscosity is responsible
 395 for the energy dissipation of short wavelengths. The growth rate of the instability then
 396 decreases as the inverse of the wavenumber (Chandrasekhar 1961). Consequently, a mode
 397 of maximum instability depending on the acceleration history and impact parameters is
 398 expected to develop. This mode likely determines the typical number of plumes and the
 399 corresponding wavelength. The growth rate of the instability and the wavelength selection
 400 are also possibly influenced by the thin layer configuration, with a stabilising effect, if the
 401 mixing layer thickness is much smaller than the typical instability wavelength (Keller &
 402 Kolodner 1954; Villermaux 2020).

403 At longer times ($t^* \gtrsim 10$), geometrical effects produced by crater expansion loose intensity
 404 and become comparable with the instability. This coincides with a stage where the instability
 405 is strongly influenced by three-dimensional effects, leading to the formation of plumes below
 406 the hemispherical surface of the crater (figure 4d, iii). As the instability grows towards a
 407 more turbulent layer, the mode of maximum instability is modified by non-linear interactions.
 408 Plumes then start interacting with each other, producing a mixing layer (figure 4d, iv).

409 The importance of the density ratio is apparent in figure 6, which shows the LIF intensity
 410 when the crater reaches its maximum size, as a function of the Froude number and density
 411 ratio. For density ratios smaller than unity (first column), no instability is observed. For

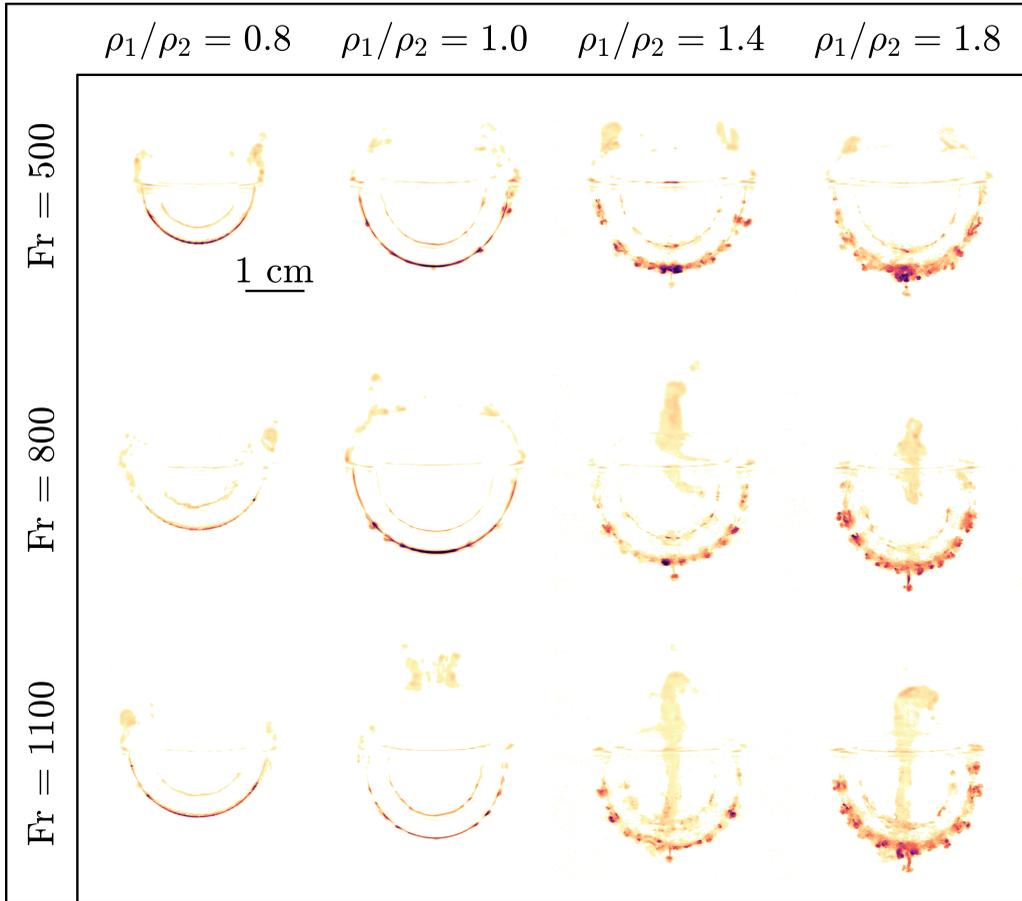


Figure 6: Snapshots of the LIF intensity field when the crater reaches its maximum size, as a function of the Froude number and the density ratio.

412 density ratios about unity (second column), small corrugations are observed between the
 413 dyed liquid and the ambient, possibly resulting from a large-scale shear instability in the
 414 layer. During crater opening, the air-water interface is not exactly hemispherical and the
 415 velocity field is not exactly radial (Bisighini *et al.* 2010). This creates a velocity shear across
 416 the interface, which may lead to the development of a Kelvin-Helmholtz instability. For
 417 density ratios larger than unity (third and fourth columns), the thickness of the dyed liquid
 418 layer is significantly larger than in other cases, as a consequence of the development of the
 419 RT instability. For a given Froude number, the mixing layer thickness obtained when the
 420 crater reaches its maximum size increases with the density ratio. For a given density ratio,
 421 the mixing layer thickness does not change significantly with the Froude number (figure 6).

422 A point of terminology may be in order here. The configuration of our experiments falls
 423 somewhat in-between the canonical Rayleigh-Taylor configuration, in which the acceleration
 424 is constant and incompressible Richtmyer-Meshkov instability, in which the acceleration
 425 changes impulsively (Richtmyer 1960; Meshkov 1969; Jacobs & Sheeley 1996). Since in our
 426 experiments the acceleration varies continuously during the crater opening phase, we have
 427 chosen to refer to it as a Rayleigh-Taylor instability, as has been done before in the literature
 428 (e.g. Mikaelian 1990, 2016). However, the magnitude of the acceleration decreases quite
 429 fast, by almost two orders of magnitude, (figure 7c of §4.2), giving it a somewhat impulsive

430 nature. This suggests that the development of the instability may share some similarities with
 431 the Richtmyer-Meshkov instability. .

432 4. Evolution of the crater size

433 Experiments in the backlight configuration provide the time evolution and the maximum of
 434 the mean crater radius, a required step in the understanding of the RT instability. We derive
 435 an energy model for the evolution of the crater radius, velocity and acceleration and compare
 436 it with experiments, from which we obtain scaling laws for the maximum crater radius and
 437 the crater opening time.

438 4.1. Energy conservation model

439 We use an energy conservation model (Engel 1966, 1967; Leng 2001) accounting for the
 440 density difference between the impacting drop and the targeted pool. We assume the crater
 441 shape to be hemispherical and the flow around the crater to be incompressible and irrotational.
 442 Since the crater opening dynamics is mainly driven by impactor inertia ($Re \gtrsim 2500$) and
 443 gravity forces, viscous dissipation is not included in the model. The formation of the crown
 444 and the surface wave during the impact, in particular their potential, kinetic and surface
 445 energies, are not explicitly included in the model.

446 On the basis of these assumptions, the sum of the crater potential energy E_p , the crater
 447 surface energy E_σ and the crater kinetic energy E_k , at any instant of time is equal to the
 448 impacting drop kinetic energy E_i just before the impact. The potential energy of the crater is

$$449 \quad E_p = \iiint \rho_2 g z dV = \int_0^R \rho_2 g \pi (R^2 - z^2) z dz = \frac{1}{4} \pi \rho_2 g R^4, \quad (4.1)$$

450 where R is the mean crater radius (section 2) and z is the depth. The crater surface
 451 energy corresponds to the formation of new surface due to crater opening. It is equal to
 452 the difference of surface energy between the initially planar surface area of the pool πR^2 and
 453 the hemispherical surface area of the cavity $2\pi R^2$, *i.e.*

$$454 \quad E_\sigma = \sigma (2\pi R^2 - \pi R^2) = \sigma \pi R^2. \quad (4.2)$$

455 The crater kinetic energy corresponds to the kinetic energy of the pool fluid below the initial
 456 surface and is related to the flow velocity potential. A radial velocity potential of the form
 457 $\Phi = -A/r$, the solution of the Laplace equation $\nabla^2 \Phi = 0$, is able to satisfy the boundary
 458 conditions at $r = R$. At the crater boundary, the radial velocity is $u_r(r = R) = (\partial \Phi / \partial r)_{r=R} =$
 459 \dot{R} , giving $A = \dot{R} R^2$ and

$$460 \quad \Phi = -\frac{\dot{R} R^2}{r}. \quad (4.3)$$

461 The radial velocity, the tangential velocity and the resulting velocity are, respectively

$$462 \quad \begin{cases} u_r = \frac{\dot{R} R^2}{r^2}, \\ u_\theta = 0, \\ \|\mathbf{u}\| = \sqrt{u_r^2 + u_\theta^2} = \frac{\dot{R} R^2}{r^2}. \end{cases} \quad (4.4)$$

463 The crater kinetic energy is then

$$464 \quad E_k = \int \frac{1}{2} \rho_2 \|\mathbf{u}\|^2 dV = \int_R^{+\infty} \pi \rho_2 \dot{R}^2 R^4 \frac{1}{r^2} dr = \pi \rho_2 \dot{R}^2 R^3. \quad (4.5)$$

465 The impacting drop kinetic energy is

$$466 \quad E_i = \frac{2}{3} \pi \rho_1 R_i^3 U_i^2. \quad (4.6)$$

467 Energy conservation between E_p , E_σ , E_k and E_i gives

$$468 \quad \frac{1}{4} \rho_2 g R^4 + \sigma R^2 + \xi \rho_2 R^3 \dot{R}^2 = \frac{2}{3} \phi \rho_1 R_i^3 U_i^2, \quad (4.7)$$

469 where ϕ and ξ are fitted parameters. The coefficient ϕ corresponds to a correction parameter
470 accounting for the terms not included in the model, *i.e.* viscous dissipation and crown energy
471 terms. The coefficient ξ is a correction parameter accounting for the difference between the
472 deliberately simplified velocity field used in the model and the true flow.

473 Normalising the crater radius and opening velocity by the impacting drop radius R_i and
474 velocity U_i , respectively, energy conservation becomes

$$475 \quad \frac{1}{4} \frac{1}{Fr^*} R^{*4} + \frac{1}{Fr^* Bo} R^{*2} + \xi \left(\frac{\rho_1}{\rho_2} \right)^{-1} R^{*3} \dot{R}^{*2} = \frac{2}{3} \phi, \quad (4.8)$$

476 where $R^* = R/R_i$ and $\dot{R}^* = \dot{R}/U_i$ are dimensionless.

477 For each experiment, we calculate ϕ as follows. Assuming that the velocity field vanishes
478 simultaneously in the pool (Prosperetti & Oguz 1993), the crater kinetic energy vanishes
479 when the crater reaches its maximum size, which, taking $R^* = R_{max}^{*exp}$ and $\dot{R}^* = 0$ in equation
480 4.8, gives

$$481 \quad \phi = \frac{3}{2} \frac{1}{Fr^*} \left(R_{max}^{*exp} \right)^2 \left[\frac{1}{4} \left(R_{max}^{*exp} \right)^2 + \frac{1}{Bo} \right]. \quad (4.9)$$

482 Knowing ϕ , we then fit the time evolution of the mean crater radius to the experiments with
483 equation 4.8 using a least square method, the kinetic energy correction parameter ξ being
484 a fit parameter. Fitting ϕ and ξ for each experiment shows that both parameters depend on
485 the Froude number (see appendix B for details). Knowing ϕ and ξ , we solve the ordinary
486 differential equation 4.8 using the boundary condition $R^*(1) = 1$. This assumes that the
487 crater radius is initially the same as the drop radius, at $t = R_i/U_i$.

488 4.2. Time evolution

489 Figure 7 shows the evolution of the mean crater radius (figure 7a), the mean crater velocity
490 (figure 7b) and the mean crater acceleration (figure 7c). This figure compares the fitted energy
491 model with experimental data, in two reference cases, with and without density difference
492 between the impacting drop and the pool. In both cases, the fitted mean crater radius, opening
493 velocity and acceleration are in close agreement with the experimental data. In the $\rho_1/\rho_2 = 1$
494 case, $\phi = 0.40$ and $\xi = 0.35$. In the $\rho_1/\rho_2 = 1.8$ case, $\phi = 0.39$ and $\xi = 0.34$. Experimental
495 data are consistent with the qualitative observations of section 3. The crater first opens, before
496 it reaches its maximum size and eventually starts to collapse.

497 At early times, the crater potential and surface energies are negligible in comparison with
498 the crater kinetic energy. This implies that the kinetic energy of the impactor is converted
499 exclusively into kinetic energy of the flow around the crater. We thus neglect the two first
500 terms on the left-hand side of equation 4.8, which in this limit gives

$$501 \quad \begin{cases} R^* = [Q(t^* - 1) + 1]^{2/5} \\ \dot{R}^* = \frac{2}{5} Q [Q(t^* - 1) + 1]^{-3/5} \\ \ddot{R}^* = -\frac{6}{25} Q^2 [Q(t^* - 1) + 1]^{-8/5} \end{cases}, \quad (4.10)$$

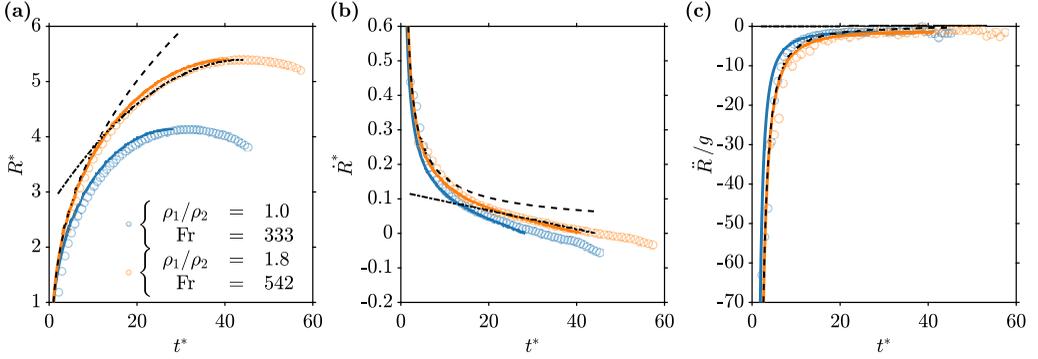


Figure 7: Time evolution of the normalised mean crater radius R^* (a), the normalised mean crater velocity \dot{R}^* (b) and the mean crater acceleration \ddot{R}^* normalised by the acceleration of gravity g (c), for two impact parameters. Circles and solid lines correspond, respectively, to experimental data and fitted energy model (equation 4.8). Dashed lines and dash-dotted lines correspond, respectively, to early-time power-law solution (equation 4.10) and late-time quadratic solution (equations 4.12) with $\rho_1/\rho_2 = 1.8$ and $Fr = 542$.

502 where $Q = \left(\frac{25}{6} \frac{\phi}{\xi} \frac{\rho_1}{\rho_2}\right)^{1/2}$. This consistently verifies the imposed boundary condition $R^*(1) = 1$
 503 and depends on the density ratio ρ_1/ρ_2 and on the correction parameters ϕ and ξ . It is in
 504 agreement with experimental data at early times (figure 7, dashed lines) and similar scaling
 505 laws from previous works (Leng 2001; Bisighini *et al.* 2010).

506 At late times, the crater velocity becomes very small. If surface tension can be neglected,
 507 taking the time derivative of equation 4.8 and then making the assumption $\dot{R}^* = 0$ gives

$$508 \quad \ddot{R}^* = -\frac{1}{2} \frac{1}{Fr\xi}. \quad (4.11)$$

509 Using $R^*(t_{max}^*) = R_{max}^*$ and $\dot{R}^*(t_{max}^*) = 0$ as boundary conditions, we obtain a quadratic
 510 solution

$$511 \quad \begin{cases} \ddot{R}^* = -\frac{1}{2} \frac{1}{Fr\xi} \\ \dot{R}^* = -\frac{1}{2} \frac{1}{Fr\xi} (t^* - t_{max}^*) \\ R^* = R_{max}^* - \frac{1}{4} \frac{1}{Fr\xi} (t^* - t_{max}^*)^2 \end{cases} . \quad (4.12)$$

512 This is in good agreement with experimental data at late times using the experimentally
 513 determined values of R_{max}^* and t_{max}^* (figure 7, dash-dotted lines). Using scaling laws for
 514 R_{max}^* and t_{max}^* (to be obtained in section 4.3), the late-time quadratic evolution of the mean
 515 crater radius can be fully predicted as function of ϕ and ξ .

516 4.3. Maximum crater radius and opening time

517 We first consider the limit of no surface tension ($Bo \rightarrow +\infty$), which amounts to neglecting
 518 the second term on the left-hand side of equation 4.8. The maximum size of the crater is then
 519 obtained by taking $\dot{R}^* = 0$ in equation 4.8, before solving for R^* to obtain

$$520 \quad R_{max}^*|_{Bo \rightarrow +\infty} = \left(\frac{8}{3}\right)^{1/4} \phi^{1/4} Fr^{*1/4}. \quad (4.13)$$

521 Figure 8a shows the maximum crater size in our experiments as a function of a least square
 522 best-fit power law scaling in the form

$$523 \quad R_{max}^{*lsq} = c_1 Fr^{*c_2}. \quad (4.14)$$

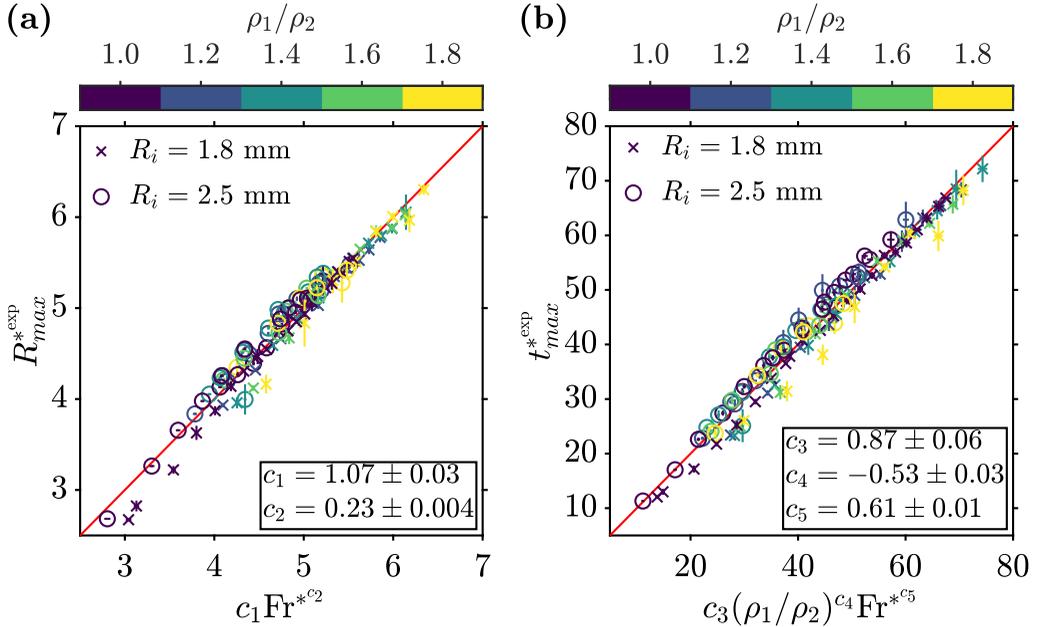


Figure 8: (a) Experimental normalised maximum crater radius $R_{max}^{*\text{exp}}$, as a function of the least squares best-fit power-law scaling $R_{max}^{*\text{lsq}}$ (equation 4.14). (b) Experimental normalised crater opening time $t_{max}^{*\text{exp}}$, as a function of the least squares best-fit power-law scaling $t_{max}^{*\text{lsq}}$ (equation 4.17). Colours scale as the density ratio ρ_1/ρ_2 . Circles and crosses correspond, respectively, to large and small drop size series.

524 The exponent $c_2 = 0.23 \pm 0.004$ for Fr^* is close to the theoretical $1/4$ prediction of equation
 525 4.13 and is in agreement with previous works on liquids (Prosperetti & Oguz 1993; Leng
 526 2001; Bisighini *et al.* 2010) and granular materials (Walsh *et al.* 2003; Takita & Sumita
 527 2013). The prefactor $c_1 = 1.07 \pm 0.03$ is close to the value predicted by the model (equation
 528 4.13): since $\phi = 0.38 \pm 0.04$ in our experiments, the predicted model prefactor is indeed
 529 equal to 1.0 ± 0.03 . The prefactor c_1 is also consistent with those obtained in previous works
 530 (e.g. $c_1 = 1.1$ in Leng (2001)).

531 We now turn to estimating the crater opening time, defined as the time t_{max}^* at which the
 532 maximum crater size is reached. Having neglected surface tension, we integrate equation 4.8
 533 between $t^* = 0$ and $t^* = t_{max}^*|_{Bo \rightarrow +\infty}$ to obtain

$$534 \quad t_{max}^*|_{Bo \rightarrow +\infty} = \xi^{1/2} \left(\frac{\rho_1}{\rho_2} \right)^{-1/2} \int_0^{R_{max}^*|_{Bo \rightarrow +\infty}} \frac{R^{*3/2}}{\left(\frac{2}{3}\phi - \frac{1}{4} \frac{R^{*4}}{Fr^*} \right)^{1/2}} dR^*. \quad (4.15)$$

535 Using equation 4.13 for $R_{max}^*|_{Bo \rightarrow +\infty}$ and integrating, we obtain

$$536 \quad t_{max}^*|_{Bo \rightarrow +\infty} = \frac{1}{2} \left(\frac{8}{3} \right)^{1/8} B \left(\frac{1}{2}, \frac{5}{8} \right) \left(\frac{\rho_1}{\rho_2} \right)^{-1/2} \phi^{1/8} \xi^{1/2} Fr^{*5/8}, \quad (4.16)$$

537 where B is the beta function.

538 Figure 8b shows the opening time in our experiments as a function of a least square best-fit
 539 power law scaling in the form

$$540 \quad t_{max}^{*\text{lsq}} = c_3 (\rho_1/\rho_2)^{c_4} Fr^{*c_5}. \quad (4.17)$$

541 The exponent $c_4 = -0.53 \pm 0.03$ for ρ_1/ρ_2 agrees with the theoretical $-1/2$ prediction of
 542 equation 4.16. The exponent $c_5 = 0.61 \pm 0.01$ for Fr^* is also close to the $5/8 = 0.625$
 543 prediction of equation 4.16 and agrees with previous works at $\rho_1/\rho_2 = 1$ (Leng 2001;
 544 Bisighini *et al.* 2010). The prefactor $c_3 = 0.87 \pm 0.06$ is close to the value predicted by
 545 equation 4.16: Since $\phi = 0.38 \pm 0.04$ and $\xi = 0.34 \pm 0.03$ in our experiments, the prefactor
 546 predicted by the model is indeed equal to 0.79 ± 0.04 . The prefactor c_3 is also consistent
 547 with prefactors obtained in previous works (*e.g.* $c_3 = 0.59$ in Leng (2001)).

548 In the above paragraphs, we develop a leading-order model that neglects surface tension.
 549 We now consider the effect of surface tension on the maximum crater size and crater opening
 550 time. With only the kinetic energy term set to zero in equation 4.8, we obtain the maximum
 551 crater radius including surface tension

$$552 \quad R_{max}^* = R_{max|Bo \rightarrow +\infty}^* \left[\sqrt{1 + \frac{3}{2}(\sqrt{Fr^* \phi} Bo)^{-2}} - \frac{\sqrt{6}}{2}(\sqrt{Fr^* \phi} Bo)^{-1} \right]^{1/2}. \quad (4.18)$$

553 This scaling depends on the dimensionless parameter $(\sqrt{Fr^* \phi} Bo)^{-1}$, which brings in
 554 the effect of surface tension on the cratering dynamics. When $Bo \rightarrow +\infty$, we have
 555 $(\sqrt{Fr^* \phi} Bo)^{-1} \rightarrow 0$ and we retrieve the scaling without surface tension (equation 4.13).

556 Figure 9a shows the maximum crater radius normalised by the maximum crater radius
 557 scaling without surface tension $R_{max}^{*exp}/R_{max|Bo \rightarrow +\infty}^*$, as a function of $(\sqrt{Fr^* \phi} Bo)^{-1}$. This
 558 corresponds to the ratio between the experimental data and the scaling law without surface
 559 tension (equation 4.13). As expected, the scaling without surface tension overestimates the
 560 experimental maximum crater radius because it neglects surface energy. This overestimate
 561 decreases with $(\sqrt{Fr^* \phi} Bo)^{-1}$, *i.e.* when surface tension effects become negligible in
 562 comparison with impactor inertia and gravity forces. The difference between experimental
 563 data and the scaling law without surface tension is properly accounted for by equation 4.18.

564 We finally obtain an expression for the crater opening time by integrating equation 4.8
 565 between $t^* = 0$ and $t^* = t_{max}^*$:

$$566 \quad t_{max}^* = \xi^{1/2} \left(\frac{\rho_1}{\rho_2} \right)^{-1/2} \int_0^{R_{max}^*} \frac{R^{*3/2}}{\left(\frac{2}{3}\phi - \frac{R^{*2}}{Fr^* Bo} - \frac{1}{4} \frac{R^{*4}}{Fr^*} \right)^{1/2}} dR^*. \quad (4.19)$$

567 Writing $\mathcal{R} = R^*/R_{max}^*$, this gives

$$568 \quad t_{max}^* = t_{max|Bo \rightarrow +\infty}^* \frac{4}{B \left(\frac{1}{2}, \frac{5}{8} \right)} \int_0^1 \frac{f(x)^{5/4} \mathcal{R}^{3/2}}{\left[1 - \mathcal{R}^4 + 2xf(x)(\mathcal{R}^4 - \mathcal{R}^2) \right]^{1/2}} d\mathcal{R}, \quad (4.20)$$

569 where $f(x) = \sqrt{1+x^2} - x$ and $x = (2Fr^* \phi Bo^2/3)^{-1/2}$. Using a first-order Taylor expansion
 570 in $(\sqrt{Fr^* \phi} Bo)^{-1}$, we obtain

$$571 \quad t_{max}^* = t_{max|Bo \rightarrow +\infty}^* \left[1 - \frac{\sqrt{6} B \left(\frac{1}{2}, \frac{1}{8} \right)}{8 B \left(\frac{1}{2}, \frac{5}{8} \right)} (\sqrt{Fr^* \phi} Bo)^{-1} + O \left((\sqrt{Fr^* \phi} Bo)^{-2} \right) \right]. \quad (4.21)$$

572 Figure 9b shows the crater opening time normalised by the crater opening time scaling
 573 without surface tension $t_{max}^{*exp}/t_{max|Bo \rightarrow +\infty}^*$, as a function of $(\sqrt{Fr^* \phi} Bo)^{-1}$. This corresponds
 574 to the ratio between experimental data and the scaling law without surface tension (equation
 575 4.16). Although this scaling law without surface tension is close to the experimental opening
 576 time, it increasingly overestimates experiments as $(\sqrt{Fr^* \phi} Bo)^{-1}$ increases, *i.e.* when surface

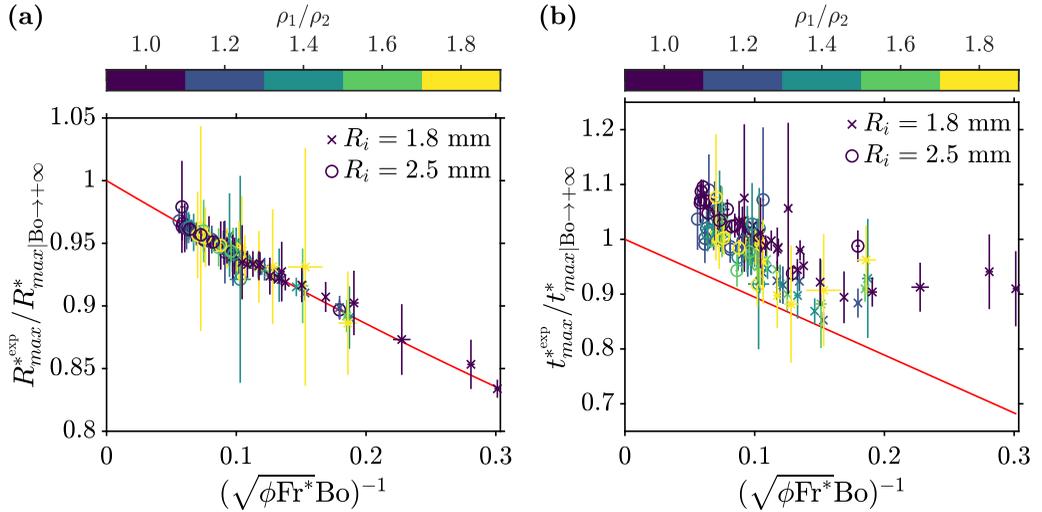


Figure 9: (a) Experimental maximum crater radius normalised by the maximum crater radius scaling without surface tension $R_{max}^{*exp}/R_{max|Bo \rightarrow +\infty}^*$ (equation 4.13), as a function of $1/(\sqrt{\phi Fr^* Bo})$. The solid line corresponds to the surface tension correction of equation 4.18. (b) Experimental crater opening time normalised by the crater opening time scaling without surface tension $t_{max}^{*exp}/t_{max|Bo \rightarrow +\infty}^*$ (equation 4.16), as a function of $1/(\sqrt{\phi Fr^* Bo})$. The solid line corresponds to the surface tension correction of equation 4.21. Colours scale as the density ratio ρ_1/ρ_2 . Circles and crosses correspond, respectively, to large and small drop size series.

577 tension effects become significant. When surface energy is significant, we indeed expect the
 578 crater to open at a slower rate. The difference between experimental data and the scaling
 579 law without surface tension is corrected in part by equation 4.21, which is in reasonable
 580 agreement with the experimental data (the relative error is $\pm 10\%$). The residual differences
 581 between experimental and predicted crater opening time may come from our assumption
 582 that the correction parameters ϕ and ξ are independent of time. This also explains why
 583 residual differences of the maximum crater radius are smaller. Since we obtain the maximum
 584 crater radius assuming that the kinetic energy term of equation 4.8 vanishes, we predict the
 585 maximum crater radius without time integration.

586 These scaling laws, with and without surface tension, are now discussed by normalising
 587 the time evolution of the crater radius for all experiments (figure 10). Figure 10a shows the
 588 crater radius normalised by scaling laws neglecting surface tension (equations 4.13 and 4.16).
 589 With this normalisation, we expect the experimental crater radius to collapse on the case
 590 without surface tension (figure 10a, red line). However, we obtain a residual dependency on
 591 the dimensionless parameter related to surface tension $(\sqrt{Fr^* \phi Bo})^{-1}$. Figure 10b shows the
 592 crater radius normalised by scaling laws including surface tension (equations 4.18 and 4.21).
 593 We find that the experimental crater radius collapses better when accounting for surface
 594 tension effects (figure 10b).

595 Although the maximum crater radius collapses in figure 10b, there is a residual dispersion
 596 on the crater radius evolution at early times when surface tension effects are significant.
 597 This is explained by renormalising the crater radius as $\tilde{R} = R^*/R_{max|Bo \rightarrow +\infty}^*$ and time as

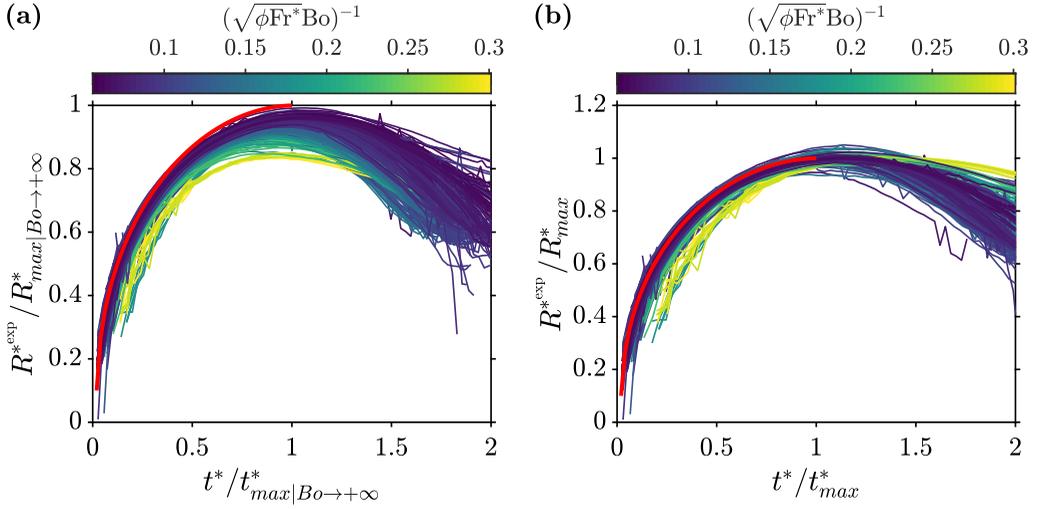


Figure 10: (a) Experimental mean crater radius normalised by the maximum crater radius scaling $R^{*\text{exp}} / R_{\text{max}|Bo \rightarrow +\infty}^*$, as a function of time normalised by the crater opening time scaling $t^* / t_{\text{max}|Bo \rightarrow +\infty}^*$, with scaling laws neglecting surface tension. (b) Experimental mean crater radius normalised by the maximum crater radius scaling $R^{*\text{exp}} / R_{\text{max}}^*$, as a function of time normalised by the crater opening time scaling t^* / t_{max}^* , with scaling laws including surface tension. The thick solid line corresponds to the solution of equation 4.8 when $1/(\sqrt{\phi Fr^* Bo}) = 0$, *i.e.* without surface tension. Colours scale as $1/(\sqrt{\phi Fr^* Bo})$.

598 $\tilde{t} = t^* / t_{\text{max}|Bo \rightarrow +\infty}^*$ in equation 4.8. This gives

$$599 \quad \tilde{R}^4 + \frac{\sqrt{6}}{\sqrt{Fr^* \phi Bo}} \tilde{R}^2 + \frac{16}{B \left(\frac{1}{2}, \frac{5}{8}\right)^2} \tilde{R}^3 \left(\frac{d\tilde{R}}{d\tilde{t}}\right)^2 = 1. \quad (4.22)$$

600 The surface tension term depends on $(\sqrt{Fr^* \phi Bo})^{-1}$, which results in a residual dispersion
 601 of experimental curves at early times. This remains correct when we normalise equation 4.8
 602 with scaling laws including surface tension, although the surface tension term depends on
 603 $(\sqrt{Fr^* \phi Bo})^{-1}$ in a different way. We also obtain a residual dispersion of the crater radius
 604 at late times (figure 10b). Since the energy model (equation 4.8) applies only during the
 605 opening of the crater, we expect the experimental curves to collapse only when $t^* < t_{\text{max}}^{*\text{exp}}$.

606 5. Evolution of the RT instability

607 Experiments in the LIF configuration allow us to describe quantitatively the evolution of the
 608 thickness of the mixing layer. Using the energy conservation model (section 4), we derive
 609 a model for the mixing layer thickness evolution, from which we obtain scaling laws for
 610 the maximum mixing layer thickness. We also obtain a theoretical scaling for the early-time
 611 instability wavelength, which we compare with experiments (appendix A).

612 5.1. Buoyancy–drag model

613 In addition to the energy conservation model assumptions (section 4), we assume the mixing
 614 layer to be uniform with a constant thickness around the crater boundary. We first consider
 615 the situation where no instability develops. In this situation, the volume of the layer remains
 616 constant. As the crater radius increases, the drop liquid spreads over an increasingly large

617 surface area and the mean layer thickness h thus decreases with time (figure 5a). Denoting
 618 by $\bar{u}(r, t)$ the laterally averaged velocity field associated with the opening of the crater, the
 619 time derivative of h is then given by

$$620 \quad \dot{h} = \bar{u}(R + h) - \bar{u}(R). \quad (5.1)$$

621 Since $\bar{u} = \dot{R}(R/r)^2$ corresponds to the radial potential flow of equation 4.4, this gives

$$622 \quad \dot{h} = \dot{R} \left[\frac{R^2}{(R + h)^2} - 1 \right], \quad (5.2)$$

623 where the right-hand side will be referred to as a geometrical thinning term. This equation
 624 originates from the mass conservation of the layer in the absence of instability.

625 We now consider the effect of the RT instability (figures 5b and 11), which we model
 626 as an entrainment process. Assuming that the ambient liquid is gradually incorporated into
 627 the mixing layer at a rate $u'(t)$ (a volumetric flux), the time derivative of h is then the sum
 628 of the geometrical thinning term (equation 5.2) and entrainment rate u' (figure 11c). The
 629 velocity u' and length scale h also correspond to the velocity and integral length scale of a
 630 mixing-length turbulent model describing the mixing layer. After non-dimensionalization,
 631 this writes as

$$632 \quad \dot{h}^* = \dot{R}^* \left[\frac{R^{*2}}{(R^* + h^*)^2} - 1 \right] + u'^*. \quad (5.3)$$

633 Now u' can also be seen as the velocity the tip of the RT plumes would have in the absence
 634 of geometrical thinning. With this interpretation in mind, we describe the evolution of u'
 635 using a buoyancy–drag model of the mixing layer (Dimonte 2000; Oron *et al.* 2001; Zhou
 636 2017). We consider that the RT plumes with a density $\bar{\rho} = \rho_2 + \Delta\rho$ penetrate into the less
 637 dense surrounding liquid with a density ρ_2 . The equation of motion then reads as

$$638 \quad \bar{\rho} \frac{du'^*}{dt^*} = -\beta\Delta\rho\ddot{R}^* - C\rho_2 \frac{u'^{*2}}{h^*}, \quad (5.4)$$

639 where β and C are the RT buoyancy prefactor and the drag coefficient, respectively. This
 640 equation corresponds to a balance between the fluid inertia on the left-hand side, buoyancy
 641 in the first term on the right-hand side and inertial drag in the second term on the right-hand
 642 side. The acceleration of the crater boundary \ddot{R} being significantly larger than g (figure 7),
 643 we neglect Earth's gravity in the buoyancy term.

644 Using mass conservation in the uniform, hemispherical and thin mixing layer, the
 645 dimensionless density difference is

$$646 \quad \frac{\Delta\rho}{\bar{\rho}} = \frac{1}{1 + \frac{3}{2}R^{*2}h^* \frac{\rho_2}{\Delta\rho_0}}, \quad (5.5)$$

647 where $\Delta\rho_0$ is the initial density difference between the impacting drop and the pool.
 648 Inserting the density excess (equation 5.5) into the equation of motion (equation 5.4) made
 649 dimensionless, we obtain

$$650 \quad \frac{du'^*}{dt^*} = -\beta \frac{\dot{R}^*}{1 + \frac{3}{2}R^{*2}h^* \frac{\rho_2}{\Delta\rho_0}} - C \frac{1}{1 + \frac{2}{3} \frac{1}{R^{*2}h^*} \frac{\Delta\rho_0}{\rho_2}} \frac{u'^{*2}}{h^*}. \quad (5.6)$$

651 Together with the crater radius evolution (equation 4.8), equations 5.3 and 5.6 are coupled
 652 ordinary differential equations. We solve this initial value problem numerically using fixed
 653 initial conditions at $t^* = 1$. As in section 4, we choose $R^*(1) = 1$, which means that the crater
 654 radius is initially the same as the drop radius. We choose the initial mixing layer thickness

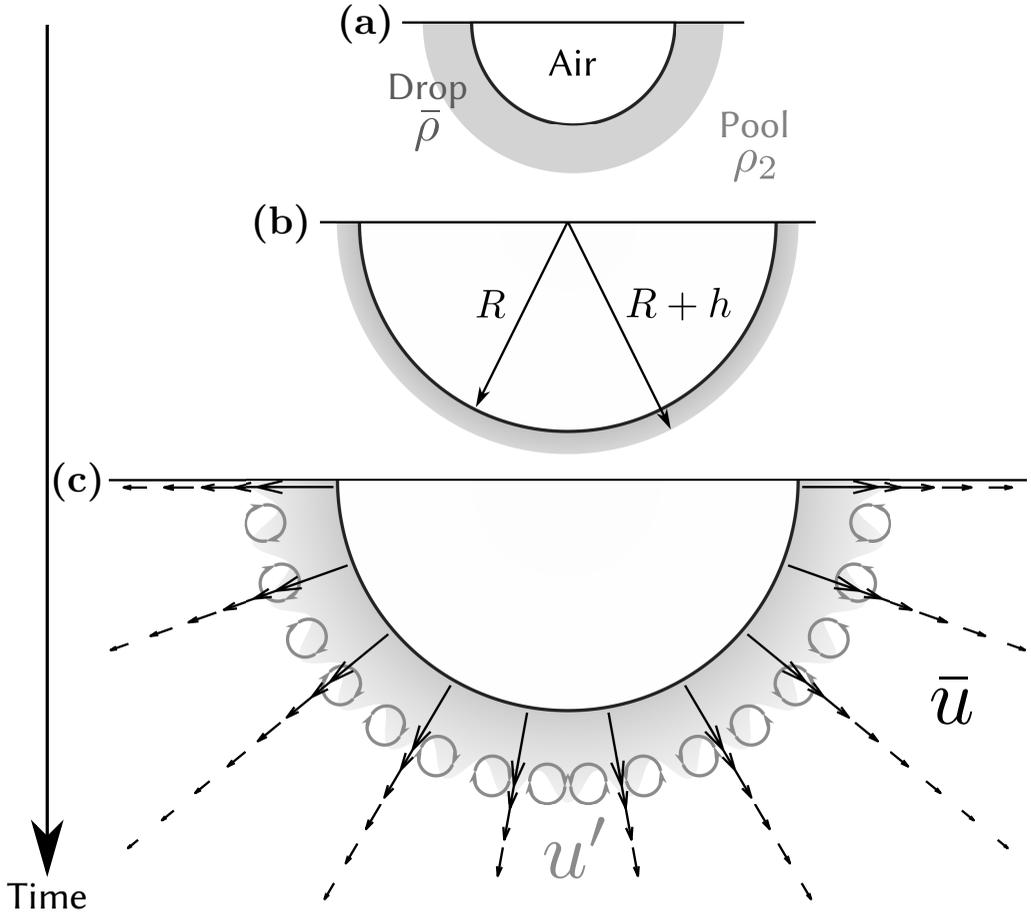


Figure 11: Mixing layer evolution. After having spread quickly on the crater boundary to become a thick layer (a), the liquid layer of the drop gradually gets thinner as the crater grows (b). At some point, crater expansion becomes weak enough, allowing for the RT instability to develop (c). We decompose the velocity field into a velocity component \bar{u} produced by crater opening and velocity fluctuations u' produced by the RT instability.

655 $h^*(1)$ by fitting the experiments without density contrast ($\rho_1/\rho_2 = 1$) with the system of
 656 differential equations without entrainment (equations 4.8 and 5.2). Parameter $h^*(1)$ is a fitting
 657 parameter, as well as the energy partitioning coefficient ϕ and the kinetic energy correction
 658 coefficient ξ (section 4). We experimentally obtain $h^*(1) = h_0^* = 0.62 \pm 0.15$. This value
 659 is larger than the theoretical $h^*(1) = 3^{1/3} - 1 \approx 0.44$ obtained when the liquid of the drop
 660 is distributed uniformly around a hemispherical crater with a radius $R^*(1) = 1$. This may
 661 be explained by the non-hemispherical cavity at the beginning of crater opening, the initial
 662 accumulation of the fluid of the drop on the crater floor and possible initial interpenetration
 663 between the drop and the ambient fluid. We also choose $u'^*(1) = 0$, assuming that the
 664 amplitude of the velocity fluctuations in the layer are initially small. For each experiment,
 665 using these fixed initial conditions, we use the experimentally measured crater radius $R^{*\text{exp}}$
 666 and mixing layer thickness $h^{*\text{exp}}$ to determine the best value for the fitting parameters of the
 667 system of differential equations. This includes the energy partitioning coefficient ϕ and the
 668 kinetic energy correction coefficient ξ (equation 4.8), as well as the buoyancy prefactor β
 669 and the drag coefficient C (equation 5.6). Fitting β and C for each experiment, we find that C

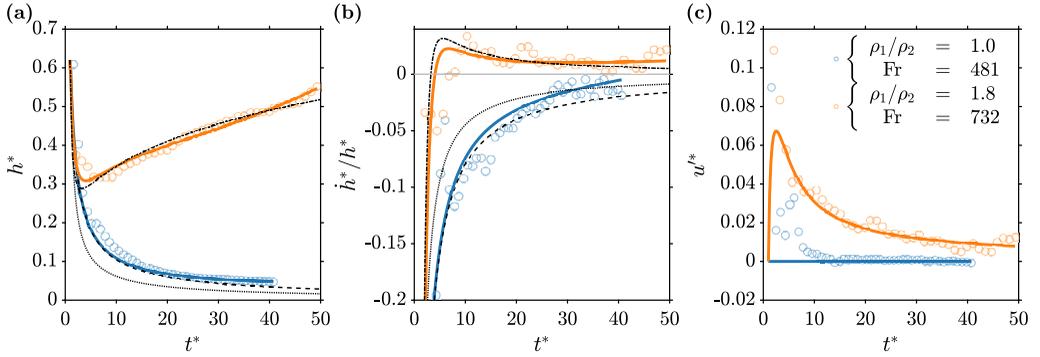


Figure 12: Time evolution of the normalised mixing layer thickness h^* (a), the normalised mixing layer growth rate \dot{h}^*/h^* (b) and the normalised inward flux due to entrainment u^* (estimated from equation 5.3) (c), for two impact parameters with and without initial density difference. Circles and solid lines correspond, respectively, to experimental data and fitted buoyancy–drag model (equations 4.8, 5.3 and 5.6). Dashed lines and dotted lines correspond, respectively, to the complete early-time power-law analytical solution (equation 5.7) and the approximate early-time power-law analytical solution (equation 5.9), calculated for the $\rho_1/\rho_2 = 1.0$ and $Fr = 481$ experiment. Dash-dotted lines correspond to the late-time power-law analytical solution (equation 5.13 with $h_0 = 0.62$, $\dot{h}_0 = -1.14$ and $C = 0.71$), calculated for the $\rho_1/\rho_2 = 1.8$ and $Fr = 732$ experiment.

670 is a decreasing function of the density ratio, with no resolvable effect of the Froude number.
 671 Also, β shows no resolvable trend with either the Froude number or the density ratio (see
 672 appendix C for details).

673

5.2. Time evolution

674 Figure 12 shows the evolution of the mixing layer thickness (figure 12a), growth rate (figure
 675 12b) and of the estimated inward flux due to entrainment (figure 12c). This figure compares
 676 the fitted mixing layer evolution model with experimental data, in two reference cases, with
 677 or without a difference of density between the impacting drop and the pool. Experimental
 678 data are consistent with the qualitative observations of section 3. The mixing layer starts to
 679 thin due to crater expansion, with a negative growth rate. Then, when the density of the drop
 680 is larger than the target liquid, it thickens owing to the RT instability, with a positive growth
 681 rate.

682 In figure 12c, we estimate u^* from equation 5.3 based on experimental measurements of R^*
 683 and h^* , *i.e.* u^* is the difference between the measured time derivative of h^* and the prediction
 684 of the model in the absence of entrainment. The model underestimates u^* for dimensionless
 685 times typically smaller than 10, for experiments with and without density difference. This
 686 underestimate may be explained by shear instabilities at the interface between the mixing
 687 layer and its surroundings (*e.g.* figure 6, $\rho_1/\rho_2 = 1.0$). These instabilities increase the growth
 688 rate of the layer and are neglected in the buoyancy–drag model, explaining why the model
 689 underestimates the inward flux u^* . Furthermore, the liquid of the drop initially accumulates
 690 at the crater floor. Since the weighted average of the mixing layer thickness h^* gives more
 691 weight to the bottom of the crater (equation 2.7), the initial mean thickness is then larger in
 692 the experiments than in the model, where the layer thickness is uniform. The liquid of the
 693 drop then flows on the crater sides. For a given volume of the mixing layer, the measured
 694 layer velocity \dot{h}^* is then larger than the early-time velocity predicted with a uniform mixing
 695 layer (equation 5.1). This also explains why the uniform model underestimates the inward
 696 flux u^* in comparison with experimental data.

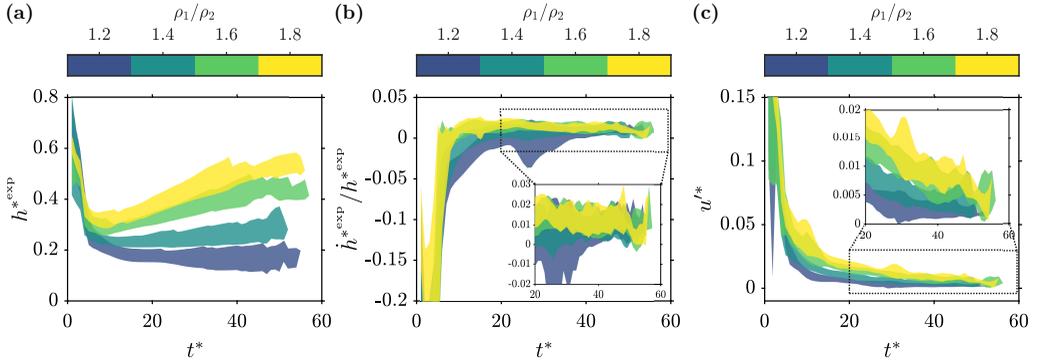


Figure 13: Time evolution of the experimental normalised mixing layer thickness $h^{*\text{exp}}$ (a), the experimental mixing layer growth rate $\dot{h}^{*\text{exp}}/h^{*\text{exp}}$ (b) and the normalised inward flux due to entrainment u'^* (c) (estimated from equation 5.3). Experiments are clustered by density ratio group, the extent of which is defined by the standard deviation of the experiments in that group. Colours scale as the density ratio ρ_1/ρ_2 .

697 Figure 13 shows the mixing layer thickness $h^{*\text{exp}}$ (figure 13a), growth rate $\dot{h}^{*\text{exp}}/h^{*\text{exp}}$
 698 (figure 13b) and the estimated entrainment term u'^* (figure 13c) as a function of time, for all
 699 experiments, grouped by density ratio. As suggested by figure 6, the mixing layer evolution,
 700 due to the RT instability, is mainly dictated by the density ratio between the drop and its
 701 surroundings. In the entrainment stage (typically $t^* > 15$), a larger initial density difference
 702 promotes entrainment by the RT instability through an increased entrainment term u'^* (figure
 703 13c), leading to an increased mixing layer growth rate (figure 13a,b). The local Reynolds
 704 number in the mixing layer depends on the density ratio and decreases with time, but is
 705 typically in the range 1 – 200.

706 5.3. Thinning stage, entrainment stage and transition time scale

707 We now focus on the numerical solution of the coupled ordinary differential equations 4.8,
 708 5.3 and 5.6. Figure 14 shows the geometrical thinning term and the entrainment term in
 709 equation 5.3, as well as the buoyancy term and the drag term in equation 5.6, as a function
 710 of time. Based on this figure, we identify two stages in the evolution of the mixing layer.

711 The first stage, referred to as the thinning stage, is defined by a negative growth rate of the
 712 mixing layer, *i.e.* $\dot{h}^* < 0$ in figure 14a. Its dynamics is controlled by the geometrical evolution
 713 of the crater, with the geometrical thinning term much larger than the RT entrainment term
 714 in equation 5.3 (figure 14a). Since the crater deceleration is large at early times, the buoyancy
 715 term prevails over the drag term in the buoyancy–drag equation 5.6 (figure 14b). However,
 716 this does not influence the evolution of the mixing layer since the geometrical thinning term
 717 dominates.

718 Neglecting the entrainment term u'^* and assuming $h^*(1) = h_0^* = 0.62$ as initial condition
 719 for the thickness of the mixing layer, the solution of the differential equation 5.3 is

$$720 \quad h^* = \left(R^{*3} - 1 + (h_0^* + 1)^3 \right)^{1/3} - R^*. \quad (5.7)$$

721 Using the power-law solution of R^* (equation 4.10), we obtain an analytical solution for h^*
 722 in the geometrical phase (figure 12, dashed lines). Assuming $h^* \ll R^*$, which is reasonable
 723 after a few time units (*e.g.* figure 4), equation 5.3 simplifies as

$$724 \quad \dot{h}^* = -2\dot{R}^* \frac{h^*}{R^*}. \quad (5.8)$$

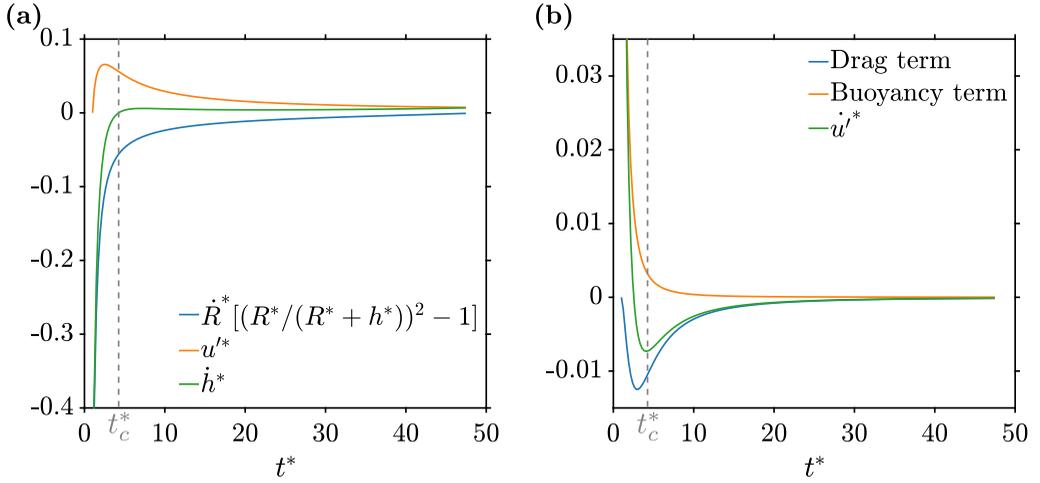


Figure 14: (a) Geometrical thinning term (first term on the right-hand side), entrainment term (second term on the right-hand side) and \dot{h}^* (left-hand side) in equation 5.3, as a function of time. (b) Buoyancy term (first term on the right-hand side), drag term (second term on the right-hand side) and \dot{u}'^* (left-hand side) in equation 5.6, as a function of time. In this numerical solution, $\rho_1/\rho_2 = 1.8$, $Fr = 572$, $\phi = 0.39$, $\xi = 0.47$, $\beta = 0.33$ and $C = 1.34$.

725 The solution then takes the form

$$726 \quad h^* = \frac{h_0^*}{R^{*2}}. \quad (5.9)$$

727 These solutions correspond to the conservation of the initial volume of the impactor, *i.e.* a
 728 sphere of unit dimensionless radius. Using the power-law solution of equation 4.10, equation
 729 5.9 also gives a power-law solution for the mixing layer thickness, velocity and acceleration

$$730 \quad \begin{cases} h^* = h_0^* [Q(t^* - 1) + 1]^{-4/5} \\ \dot{h}^* = -\frac{4}{5} h_0^* Q [Q(t^* - 1) + 1]^{-9/5} \\ \ddot{h}^* = \frac{36}{25} h_0^* Q^2 [Q(t^* - 1) + 1]^{-14/5} \end{cases}, \quad (5.10)$$

731 where $Q = \left(\frac{25}{6} \frac{\phi}{\xi} \frac{\rho_1}{\rho_2}\right)^{1/2}$. Consequently, a power-law solution for the mixing layer growth rate
 732 is

$$733 \quad \frac{\dot{h}^*}{h^*} = -\frac{4}{5} Q [Q(t^* - 1) + 1]^{-1}. \quad (5.11)$$

734 These solutions (figure 12, dotted lines) depend on the density ratio ρ_1/ρ_2 and on the
 735 correction parameters ϕ and ξ , through Q . They explain the early-time evolution of the
 736 mixing layer thickness.

737 The second stage, referred to as the entrainment stage, is defined by a positive growth rate of
 738 the mixing layer, *i.e.* $\dot{h}^* > 0$ in figure 14a. In this stage, spreading and entrainment are similar
 739 in magnitude. The dynamics is then governed by a balance between residual geometrical
 740 effects and entrainment produced by the RT instability. In this stage, the crater deceleration
 741 slows down and hence the buoyancy term quickly decreases. The rate of entrainment is
 742 therefore limited by the drag term (figure 14b). At late times, when the crater size is close
 743 to reaching its maximum, the geometrical thinning term vanishes and the dynamics is only
 744 controlled by the entrainment term (figure 14a).

745 In this stage, we use the approximation 5.8 for the geometrical term to simplify equation

746 **5.3.** We also neglect the buoyancy in equation 5.6 assuming $\frac{2}{3} \frac{\Delta\rho_0}{\rho_2} / (R^{*2} h^*) \ll 1$. These
 747 assumptions, respectively, correspond to the vanishing crater deceleration and $\Delta\rho \ll \rho_2$
 748 during the entrainment stage. With these assumptions, equations 5.3 and 5.6 become

$$749 \quad \begin{cases} \dot{h}^* = -2 \frac{\dot{R}^*}{R^*} h^* + u'^* \\ \dot{u}'^* = -C \frac{u'^{*2}}{h^*} \end{cases} . \quad (5.12)$$

750 Assuming a 2/5 power-law solution for R^* (equation 4.10) and using $h^*(1) = h_0^*$ and
 751 $\dot{h}^*(1) = \dot{h}_0^*$ as initial conditions, the solution to equation 5.12 is

$$752 \quad h^* = h_0^* [1 + Q(t^* - 1)]^{-4/5} \left\{ 1 + K \left[[1 + Q(t^* - 1)]^{9/5} - 1 \right] \right\}^{\frac{1}{1+C}}, \quad (5.13)$$

753 where $K = \frac{1}{9}(C + 1)(4 + \frac{5}{Q} \frac{\dot{h}_0^*}{h_0^*})$ (figure 12, dash-dotted lines). The value $C = 0.71$ required
 754 to fit experimental data is smaller than the value obtained by fitting experimental data with
 755 the full numerical model (equations 4.10, 5.3 and 5.6). We explain this difference with
 756 the assumptions made to obtain the analytical solution, *i.e.* approximated geometrical term,
 757 neglected buoyancy term, $\Delta\rho \ll \rho_2$ and 2/5 power-law for R^* . Using a larger value of C in
 758 the analytical solution results in an underestimate of the layer thickness in the entrainment
 759 stage.

760 The transition time t_c between the thinning stage and the entrainment stage is the time at
 761 which the growth rate changes sign and the mixing layer thickness reaches a local minimum
 762 (figure 15a). This corresponds to the time at which geometrical effects exactly balance the
 763 entrainment produced by the RT instability.

764 We measure the transition time in experiments and in figure 15b we compare the
 765 experimental values with the transition time obtained from the numerical model. Although
 766 uncertainties on the transition time are significant due to the extrinsic variability of
 767 experiments, the numerical model is rather consistent with experimental data.

768 We now derive an approximate power-law for the transition time. We assume that $h^* \ll R^*$
 769 and hence we use the approximation 5.8 for the geometrical term to solve equation 5.3. We
 770 also simplify equation 5.6 assuming $\frac{2}{3} \frac{\Delta\rho_0}{\rho_2} / (R^{*2} h^*) \ll 1$. These assumptions correspond to
 771 $\Delta\rho \ll \rho_2$, which is a reasonable assumption at the transition time. With these assumptions,
 772 equations 5.3 and 5.6 can be combined to give

$$773 \quad \ddot{h}^* + 2(2C + 1) \frac{\dot{R}^*}{R^*} \dot{h}^* + 2(2C - 1) \frac{\dot{R}^{*2}}{R^{*2}} h^* + C \frac{\dot{h}^{*2}}{h^*} + 2h^* \frac{\ddot{R}^*}{R^*} + \beta \frac{2}{3} \frac{\Delta\rho_0}{\rho_2} \frac{\ddot{R}^*}{R^{*2} h^*} = 0. \quad (5.14)$$

774 At the critical transition time t_c , $\dot{h}^*(t = t_c) = 0$. Using the power-law solutions for R^*
 775 (equation 4.10) and h^* (equation 5.10) in the thinning stage, we obtain from equation 5.14 a
 776 scaling for the dimensionless transition time t_c^* ,

$$777 \quad t_c^{*\text{lsq}} = c_6 \left\{ 1 + \frac{1}{Q} \left[\left(4h_0^{*2} \frac{C+1}{\beta \frac{\Delta\rho_0}{\rho_2}} \right)^{5/6} - 1 \right] \right\}, \quad (5.15)$$

778 where $c_6 = 1.36 \pm 0.07$ is a least squares best-fit prefactor obtained from experimental data
 779 (figure 15c). As observed experimentally in figure 13, the predicted transition time (equation
 780 5.15) decreases with the density contrast between the fluid of the impacting drop and the
 781 pool. The larger the density contrast, the quicker entrainment effects become comparable to
 782 geometrical effects.

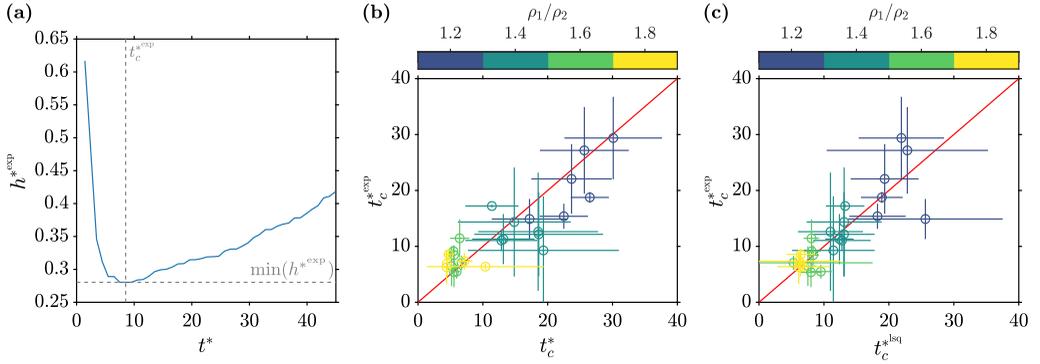


Figure 15: (a) Experimental mixing layer thickness $h^{*\text{exp}}$ as a function of time, for a single experiment. The experimental transition time $t_c^{*\text{exp}}$ between the thinning stage and the entrainment stage is defined as the time at which the mixing layer thickness reaches a local minimum. (b) Experimental transition time scale $t_c^{*\text{exp}}$ as a function of the transition time t_c^* predicted by the buoyancy–drag model (equations 4.8, 5.3 and 5.6). (c) Experimental transition time scale $t_c^{*\text{exp}}$ as a function of the transition time scaling $t_c^{*\text{lsq}}$ (equation 5.15). Colours scale as the density ratio ρ_1/ρ_2 .

783

5.4. Maximum mixing layer thickness

784 The maximum mixing layer thickness h_{max}^* first depends on the growth rate of the RT
 785 instability in the entrainment stage, an increased initial density difference leading to an
 786 increased mixing layer growth rate (figure 13). The maximum thickness also depends on
 787 the time window available for the mixing layer to actually grow without being affected by
 788 geometrical effects. The transition time t_c^* and the maximum opening time t_{max}^* correspond,
 789 respectively, to the lower and upper limits of the available time window. Since t_c^* and t_{max}^* are,
 790 respectively, a decreasing function (equation 5.15) and an increasing function (equation 4.16)
 791 of the density ratio, we expect an increased density ratio to expand the time window available
 792 for entrainment, leading to an increased maximum mixing layer thickness, consistent with
 793 figure 13a. Since t_{max}^* also increases with the Froude number, we also expect the available
 794 time window and the maximum layer thickness to increase with the Froude number.

795 We first compare the experimental maximum mixing layer thickness $h_{max}^{*\text{exp}}$ with the
 796 maximum thickness h_{max}^* obtained from the model (figure 16a). We obtain a good agreement.
 797 We then fit experimental data with the power-law scaling

$$798 \quad h_{max}^{*\text{lsq}} = c_7 \left(\frac{\rho_1}{\rho_2} \right)^{c_8} Fr^{c_9}, \quad (5.16)$$

799 where $c_7 = 0.04 \pm 0.02$, $c_8 = 2.3 \pm 0.2$ and $c_9 = 0.21 \pm 0.06$ (figure 16b). We search a scaling
 800 law as a function of ρ_1/ρ_2 because t_{max}^* and t_c^* are strong functions of ρ_1/ρ_2 (equations
 801 4.16 and 5.15). Scaling 5.16 is consistent with the predicted influence of the density ratio
 802 and the Froude number on the mixing layer growth rate and the time window available for
 803 entrainment. We indeed obtain a maximum mixing layer thickness increasing with both the
 804 initial density ratio ($c_8 > 0$) and the Froude number ($c_9 > 0$). The fact that c_8 is significantly
 805 larger than c_9 is also consistent with the qualitative observations of figure 6, *i.e.* a maximum
 806 mixing layer thickness increasing mainly with the initial density ratio.

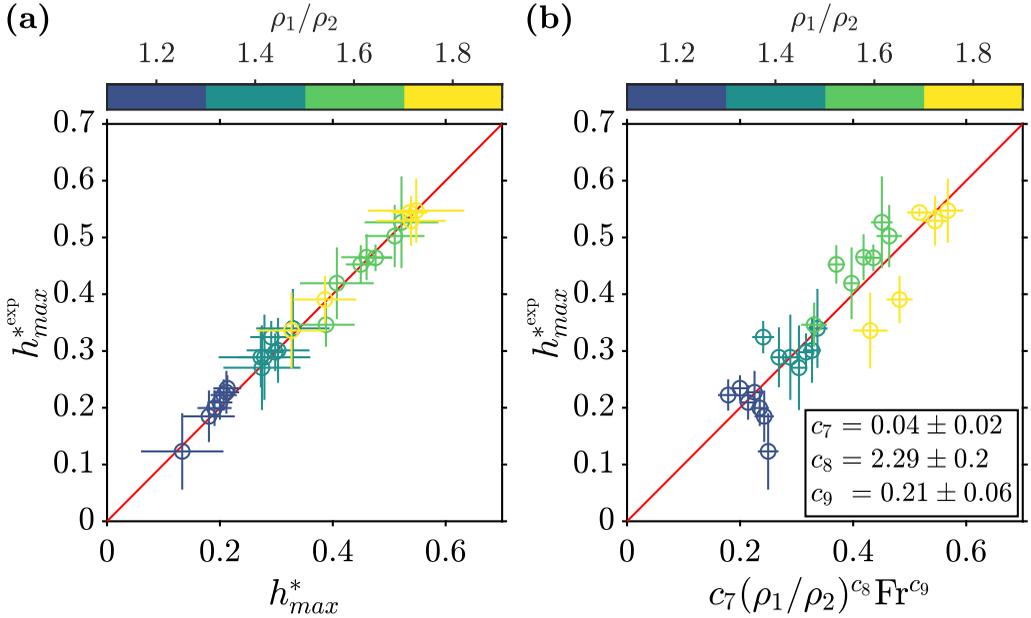


Figure 16: (a) Experimental normalized maximum mixing layer thickness h_{max}^{*exp} as a function of the maximum thickness h_{max}^* predicted by the buoyancy–drag model (equations 4.8, 5.3 and 5.6). (b) Experimental normalized maximum mixing layer thickness h_{max}^{*exp} as a function of the least square best-fit power-law scaling h_{max}^{*lsq} (equation 5.16). Colours scale as the density ratio ρ_1/ρ_2 .

807 6. Geophysical implications

808 We now apply our results to obtain a prediction of the amount of metal–silicate equilibration
 809 following an impact on a magma ocean. After the impact, the metal core of the colliding body
 810 migrates toward the planetary core due to the density contrast with the surrounding silicates
 811 (Rubie *et al.* 2015). Part of the migration occurs in a fully molten magma ocean where the
 812 metal is expected to descend as a turbulent thermal and equilibrate with silicates (Deguen
 813 *et al.* 2011, 2014). The metal phase then undergoes a vigorous stirring (Lherm & Deguen
 814 2018), leading to its fragmentation (Landeau *et al.* 2014; Wacheul *et al.* 2014; Wacheul &
 815 Le Bars 2018) into centimetric drops (Stevenson 1990; Karato & Murthy 1997; Rubie *et al.*
 816 2003; Ichikawa *et al.* 2010). However, these turbulent thermal models assume that the metal
 817 cores are released as a compact volume in the magma ocean. In contrast, recent investigations
 818 show that the impactor core equilibrates with silicates during the impact stage, prior to the
 819 fall in the magma ocean (Kendall & Melosh 2016; Landeau *et al.* 2021). Our experiments
 820 confirm this result and show that an RT instability develops during the opening of the crater,
 821 possibly equilibrating metal and silicates.

822 In order to estimate the equilibration produced by the RT instability, we estimate the mass of
 823 ambient silicates entrained by the RT instability that is likely to equilibrate with the impacting
 824 metal core. For the sake of simplicity, we assume that the mass of entrained silicates fully
 825 equilibrates, *i.e.* mixes, with the metal of the impactor. When an impactor with a radius R_i , a
 826 volume fraction of metal f_m , a metal core density ρ_m and a silicate mantle density ρ_s , impacts
 827 a planetary target, the dimensionless mass of equilibrated silicates is $\Delta = M_s/M_m$, where
 828 $M_m = f_m \rho_m (4/3) \pi R_i^3$ is the mass of the metal core and $M_s = \rho_s [2\pi R_{max}^2 h_{max} - (4/3) \pi R_i^3]$
 829 is the mass of entrained silicates (Deguen *et al.* 2014). After non-dimensionalization, the

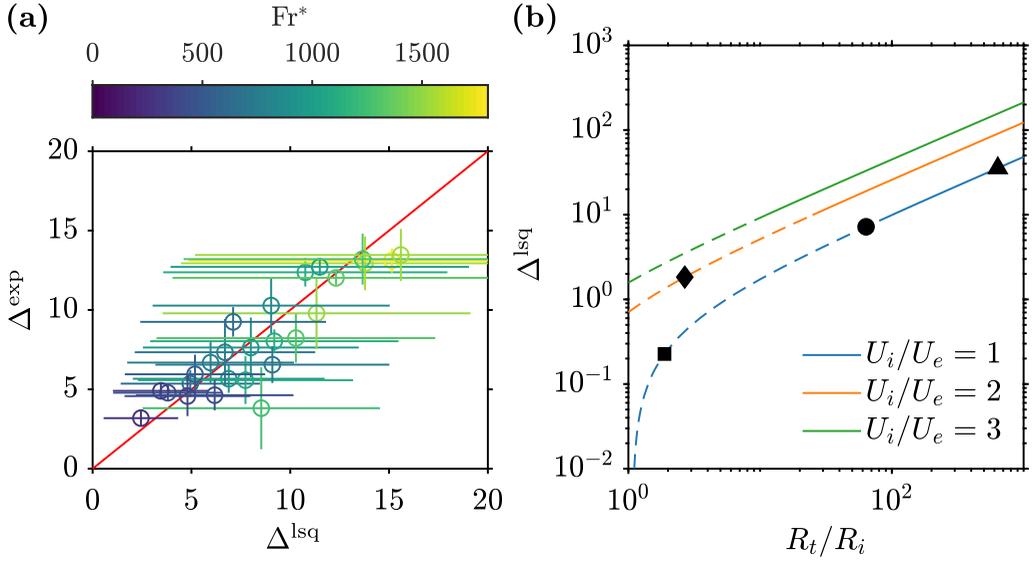


Figure 17: (a) Experimental mass of equilibrated silicates Δ^{exp} , as a function of the mass of equilibrated silicates scaling Δ^{lsq} (equation 6.2), using $f_m = 1$. Colours scale as the modified Froude number Fr^* . (b) Mass of equilibrated silicates scaling Δ^{lsq} as a function of the target-to-impactor radius R_t/R_i (equation 6.3), for several impact velocities U_i and using $f_m = 0.16$ and $\rho_m/\rho_s = 2$. Symbols and lines: triangle, impactor of 10 km in radius onto a Earth-sized target; circle, impactor of 100 km in radius onto a Earth-sized target; square, canonical Moon-forming impact with a Mars-sized impactor (Canup 2004); diamond, fast-spinning Earth Moon-forming impact with a fast ($U_i = 2U_e$) and small ($R_i/R_t = 0.3$) impactor (Cuk & Stewart 2012); dashed lines correspond to an extrapolated range of Froude number, *i.e.* $Fr < 200$, which is outside of the experimental Froude number range.

830 mass of silicates equilibrated with metal during crater opening is

$$831 \quad \Delta = \frac{\rho_s}{\rho_m} \left(\frac{3}{2} \frac{1}{f_m} R_{max}^* h_{max}^* - 1 \right). \quad (6.1)$$

832 Using scaling laws for R_{max}^* (equation 4.14) and h_{max}^* (equation 5.16), we obtain the
833 following scaling law for the mass of equilibrated silicates

$$834 \quad \Delta^{\text{lsq}} = \frac{\rho_s}{\rho_m} \left(\frac{1}{f_m} c_{10} \left(\frac{\bar{\rho}}{\rho_s} \right)^{c_{11}} Fr^{c_{12}} - 1 \right) \quad (6.2)$$

835 where $c_{10} = 0.07 \pm 0.03$, $c_{11} = 2.8 \pm 0.2$ and $c_{12} = 0.67 \pm 0.06$. In this scaling, the density
836 ratio is defined with $\bar{\rho}/\rho_s$, where $\bar{\rho} = \rho_m f_m + \rho_s (1 - f_m)$ is the mean density ratio of the
837 impactor, because this scaling derives from the crater size and the maximum mixing layer
838 thickness scaling laws (equations 4.14 and 5.16, respectively), which indeed use the mean
839 density of the impactor. We validate this scaling law against experimental data in figure 17a,
840 using $f_m = 1$ since the drop is a one-phase fluid. The predicted values of the dimensionless
841 mass of equilibrated silicates Δ^{lsq} are indeed close to the experimental values Δ^{exp} . Relatively
842 large error bars of Δ^{lsq} mainly result from the uncertainty on the prefactor coefficient of h_{max}^{lsq}
843 (equation 5.16). This means that the following geophysical applications have to be considered
844 carefully given the uncertainties on the scaling law coefficients.

845 In the context of planetary impacts, the Froude number is given by

$$846 \quad Fr = 2 \frac{R_t U_i^2}{R_i U_e^2} \quad (6.3)$$

847 where $U_e = \sqrt{2gR_t}$ is the escape velocity and R_t is the radius of the target planet. The
 848 impact velocity of colliding bodies during accretion is typically one to three times the escape
 849 velocity (Agnor *et al.* 1999; Agnor & Asphaug 2004), which means that the Froude number
 850 depends mainly on the target-to-impactor radius.

851 Using equations 6.2 and 6.3, we calculate the estimated mass of silicates equilibrated with
 852 the impacting core during crater opening Δ^{lsq} as a function of the target-to-impactor radius
 853 R_t/R_i (figure 17b). We use $f_m = 0.16$ and $\rho_m/\rho_s = 2$ to match the internal structure of a
 854 differentiated impactor (Canup 2004). Since the Froude number increases with the target-
 855 to-impactor radius, it means that smaller colliding bodies will produce more equilibration,
 856 relative to their size, than giant impactors. For example, impactors with a 10 km and 100 km
 857 radius (figure 17b, triangle and circle, respectively) will then equilibrate with 35.5 and 7.2
 858 times its own mass, respectively.

859 Several giant impact scenarios have been proposed to explain the formation of the Moon
 860 (*e.g.* Canup 2004; Cuk & Stewart 2012). We expect the canonical impact scenario with a
 861 Mars-sized impactor (Canup 2004) to equilibrate with 0.2 times its own mass during this
 862 crater opening stage (figure 17b, square). In contrast, we predict that the fast-spinning and
 863 smaller impactor proposed by Cuk & Stewart (2012) equilibrates with 1.8 times its own
 864 mass (figure 17b, diamond). These giant impacts scenarios involve small target-to-impactor
 865 radius, corresponding to an extrapolated range of Froude number, *i.e.* $Fr < 200$ (figure 17b,
 866 dashed lines), which is outside of the experimental Froude number range used to constrain the
 867 scaling. In addition, this small target-to-impactor radius is very sensitive to the uncertainty on
 868 the scaling law coefficients. The mass of equilibrated silicates extrapolated for large impactor
 869 thus has to be considered carefully.

870 Recent experiments estimate the mass of equilibrated silicates during the impact, con-
 871 sidering both the crater formation, its collapse into an upward jet and the collapse of the
 872 jet (Landeau *et al.* 2021). The 10 km impactor, the 100 km radius impactor, the canonical
 873 Moon-forming impactor and the fast-spinning Earth impactor, respectively, equilibrate with
 874 1155, 74, 1.5 and 12 times the impactor mass. This means that the fraction of silicates
 875 equilibrated during the opening stage of the crater, in comparison with the whole cratering
 876 process including the jet formation and collapse, are 3%, 10%, 13% and 15%, respectively.
 877 This is in agreement with an impact-induced equilibration mostly dominated by the collapse
 878 of the jet (Landeau *et al.* 2021).

879 Experiments including the formation and collapse of the jet (Landeau *et al.* 2021) have
 880 been done with relatively small density contrasts ($\rho_1/\rho_2 < 1.1$) and the effect of the RT
 881 instability was not as important as in the present experiments. At a given Froude number,
 882 *e.g.* $Fr = 300$, the fluid of the impactor equilibrates with 4.8 and 13.3 times its own mass
 883 of ambient fluid, if $\rho_1/\rho_2 = 1.1$ and $\rho_1/\rho_2 = 1.8$, respectively (equation 6.2). This means
 884 that the mass of ambient fluid likely to equilibrate with the impactor is 2.8 times larger in
 885 our experiments close to the metal–silicate density ratio than in the experiments of Landeau
 886 *et al.* (2021). We thus implicitly assume that equilibration produced by the RT instability and
 887 the jet are independent, whereas the equilibration produced by the jet is actually promoted by
 888 the RT instability and the dispersion of the impactor during crater opening. In other words,
 889 we assume that equilibration produced by the RT instability and the jet is combined in an
 890 additive way, whereas it is likely to be a multiplicative process. This means that we probably
 891 underestimate the influence of the jet on the overall equilibration for the large density ratios.

892 7. Conclusion

893 In this paper, we use a backlighting method and LIF to visualise the crater boundary and the
894 mixing layer produced around the cavity after the impact of a drop on a deep liquid pool.
895 We show that crater deceleration after impact is responsible for a density-driven perturbation
896 at the drop–pool interface. We interpret these perturbations as a spherical RT instability.
897 We derive an energy conservation model for the crater radius evolution (equation 4.8) and
898 compare it with backlight experiments. In particular, we obtain scaling laws for the maximum
899 crater radius (equations 4.13 and 4.18) and the crater opening time (equations 4.16 and 4.21).
900 We also derive a mixing layer evolution model (equations 5.3 and 5.6) involving two stages.
901 The mixing layer dynamics is first controlled by the geometrical evolution of the crater,
902 then by the balance between residual geometrical effects and entrainment produced by the
903 RT instability. We obtain scaling laws for the transition time scale between stages (equation
904 5.15) and the maximum mixing layer thickness (equation 5.16). From our results, we derive
905 scaling laws for equilibration between metal and silicates during a planetary impact onto a
906 magma ocean.

907 In order to validate the extrapolation of our experimental scaling laws to giant impacts,
908 experiments at lower Froude number involving large volume impactors (*e.g.* Landeau
909 *et al.* 2021) are required. This would allow us to investigate the possible effect of the
910 Reynolds number on the mixing layer. Furthermore, several physical aspects neglected in
911 our experiments have to be investigated experimentally or numerically, in order to examine
912 their effect on the cratering and equilibration dynamics. Immiscibility effects may change the
913 equilibration dynamics, in particular with the fragmentation of the metal phase. Furthermore,
914 the viscosity and diffusivity contrasts may influence thermal and chemical transfers between
915 phases, as well as the RT instability wavelength. Compressibility effects are significant
916 during the opening stage of the crater, with a Mach number larger than unity (Kendall
917 & Melosh 2016). They include the propagation of an impact shock wave and melting
918 processes. Experiments neglect these effects with a Mach number much smaller than unity.
919 Compressibility may influence the crater evolution and equilibration dynamics following the
920 impact, in particular the evolution of the mixing layer during crater opening. Nonetheless,
921 the flow velocity becomes very rapidly subsonic because the kinetic energy of the impactor
922 is quickly distributed over an increasingly large volume and because this kinetic energy is
923 converted into gravitational potential energy. Oblique impact effects are also expected since
924 the probability of a vertical impact is less than the likelihood of an oblique impact, with
925 a maximum probability for a 45° angle (Shoemaker 1961). The asymmetry caused by an
926 oblique impact modifies the dynamics of the mixing layer with the growth of a shear instability
927 around the cavity. The crater then collapses, producing an oblique jet with an angle similar to
928 the impact angle, but in the opposite direction. These features modifying the dynamics of the
929 mixing layer during crater opening and the properties of the jet following the collapse of the
930 crater crucially affect the overall equilibration of the metal phase. Differentiated impactors
931 effects are also expected since most of the Earth mass was accreted from differentiated bodies
932 (Kleine *et al.* 2002; Scherstén *et al.* 2006). Experiments currently fail to reproduce an impactor
933 with a differentiated fluid core. The existence of a mantle around the impactor’s core may
934 influence the mixing layer dynamics by changing the distribution of the metal phase around
935 the crater and the dynamics of the RT instability during crater opening. Depending on whether
936 the impactor core directly descends in the magma ocean or is entrained into the upward jet,
937 the equilibration of the metal phase may be modified. The role of impactor differentiation
938 may be investigated numerically, or experimentally using compound drops (*e.g.* Blanken
939 *et al.* 2021). However, since the compound drop experiments involve immiscible water–oil
940 drops, it will be challenging to study the mixing properties following the impact. Finally,

941 pressure and temperature effects are expected during the impact, in particular regarding
 942 large energetic impacts such as the Moon-forming giant impact. These effects neglected in
 943 experiments may result in an increased miscibility of metal and silicate phases (Morard
 944 & Katsura 2010), with potential consequences on the fragmentation of the metal phase.
 945 This may affect chemical equilibration and increase the iron content in the magma ocean.
 946 Temperature contrasts may also slightly influence the buoyancy of the metal phase, which
 947 may affect the impact dynamics in a minor way.

948 **Supplementary data.** Supplementary movies are available at
 949 <https://doi.org/10.1017/jfm.2019...>

950 **Acknowledgements.** We thank M. Moulin for the help with the design and construction of the experimental
 951 apparatus and J. Vatteville for the help with the imaging equipment. We thank three anonymous reviewers
 952 for their valuable comments that improved significantly the manuscript.

953 **Funding.** This work was supported by the European Research Council (ERC) under the European Unions
 954 Horizon 2020 research and innovation programme (grant number 716429). ISTerre is part of Labex
 955 OSUG@2020 (ANR10 LABX56). Partial funding for this research was provided by the Center for Matter
 956 at Atomic Pressure (CMAP), a National Science Foundation (NSF) Physics Frontier Center, under award
 957 PHY-2020249. Any opinions, findings, conclusions or recommendations expressed in this material are those
 958 of the author(s) and do not necessarily reflect those of the National Science Foundation.

959 **Declaration of interests.** The authors report no conflict of interest.

960 **Author ORCID.** V. Lherm, <https://orcid.org/0000-0001-5814-0637>; R. Deguen, <https://orcid.org/0000-0002-6883-0655>; T. Alboussière, <https://orcid.org/0000-0002-3692-899X>; M. Landeau, <https://orcid.org/0000-0003-4163-1311>

963 **Author contributions.** V.L., R.D. and T.A. derived the model and V.L. conducted the experiments. All
 964 authors contributed to analysing data and reaching conclusions. V.L. wrote the paper, with contributions
 965 from all co-authors.

966 **Appendix A. Early-time wavelength**

967 We convert the experimental instability wavelength into an equivalent spherical harmonic
 968 degree and we compare with an experimental scaling law. From the number of plumes n ,
 969 counted on the hemispherical section of the density interface at $t^* = 10$, we estimate the
 970 instability wavelength as $\lambda = \pi R/n$. We then obtain the corresponding spherical harmonics
 971 degree of maximum instability l_{max} using the Jeans relation (Jeans 1923)

$$972 \quad \sqrt{l_{max}(l_{max} + 1)} = \frac{2\pi R}{\lambda}. \quad (\text{A } 1)$$

973 In a thin layer configuration with no surface tension, the preferred instability wavelength
 974 is the minimum of the layer thickness and of a length scale obtained by balancing buoyancy,
 975 inertia and viscous forces. In the following, we assume that the instability wavelength depends
 976 on the buoyancy–inertia–viscosity balance. On the one hand, the balance between buoyancy
 977 and inertia is valid at small harmonic degrees l , with an instability growth rate scaling
 978 as $l^{1/2}$ (Rayleigh 1899; Taylor 1950). On the other hand, the balance between buoyancy
 979 and viscosity is valid at large harmonic degrees l , with an instability growth rate scaling
 980 as $1/l$ (Chandrasekhar 1961). The instability growth rate thus reaches a maximum when
 981 inertia, buoyancy and viscous forces are of the same order of magnitude. The corresponding
 982 instability wavelength is $\lambda \sim \{v_2^2/[(\Delta\rho_0/\rho_2)\dot{R}^*]\}^{1/3}$, which after non-dimensionalization
 983 gives

$$984 \quad \lambda^* \sim \left(\frac{\Delta\rho_0}{\rho_2}\dot{R}^*\right)^{-1/3} Re^{-2/3}. \quad (\text{A } 2)$$

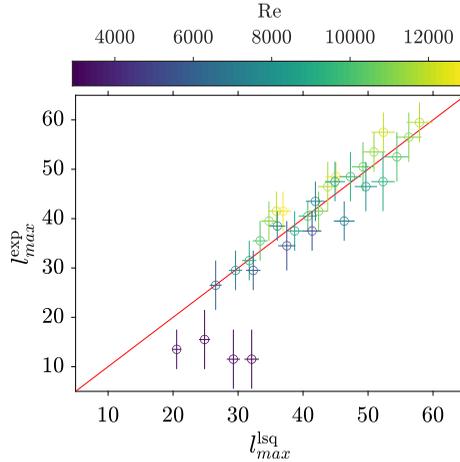


Figure 18: Experimental harmonic degree of maximum instability l_{max}^{exp} , as a function of the harmonic degree of maximum instability scaling l_{max}^{lsq} (equation A 3). Colours scale as the Reynolds number Re .

985 Assuming that \ddot{R}^* and R^* scale, respectively, as $R_{max|Bo \rightarrow \infty}^* / t_{max|Bo \rightarrow \infty}^{*2}$ and $R_{max|Bo \rightarrow \infty}^*$
 986 (see equations 4.13 and 4.16) and using the Jeans relation (equation A 1), equation A 2 gives
 987 a scaling for the degree of maximum instability

$$988 \quad l_{max}^{lsq} = c_{13} \phi^{1/4} \xi^{-1/3} Fr^{-1/12} \left(\frac{\rho_1}{\rho_2} \right)^{1/4} \left(\frac{\Delta \rho_0}{\rho_2} \right)^{1/3} Re^{2/3}, \quad (\text{A } 3)$$

989 where $c_{13} = 0.177 \pm 0.005$ is a least squares best-fit prefactor obtained from experimental
 990 data (figure 18).

991 Figure 18 shows the harmonic degrees of maximum instability measured in all experiments,
 992 as a function of scaling A 3 for l_{max} . Experimental data agree with this scaling, except for
 993 Froude numbers smaller than 100. This corresponds to Reynolds numbers smaller than
 994 4000. In this case, the crater differs from the hemispherical shape and from the purely radial
 995 acceleration assumed in the scaling. Furthermore, the fact that experimental data scale with
 996 the buoyancy–inertia–viscosity scaling (equation A 3) indicates that the preferred wavelength
 997 is indeed set by the buoyancy–inertia–viscosity balance rather than by the layer thickness.

998 Appendix B. Energy partitioning and kinetic energy correction

999 Figure 19a shows the correction parameter ϕ as a function of the Froude number. Since
 1000 several energy sinks such as crown energy and viscous dissipation are neglected in the
 1001 model, we expect ϕ to be smaller than unity. In our experiments, we find $\phi = 0.38 \pm 0.04$,
 1002 in agreement with previous works where experimental data are fitted using a partitioning
 1003 coefficient in the range 0.2 – 0.6, depending on the Froude number (Engel 1966; Olevson
 1004 1969; Leng 2001).

1005 In our experiments, Fr and We are highly correlated since Bo does not vary by a large
 1006 amount. The coefficient ϕ is found to be a decreasing function of both Fr and We , which
 1007 scales in particular as

$$1008 \quad \phi = Fr^{-0.156 \pm 0.001}, \quad (\text{B } 1)$$

1009 and is relatively independent of the density ratio and the drop size. This implies that as
 1010 the impactor inertia increases, the relative importance of the neglected energy sink terms

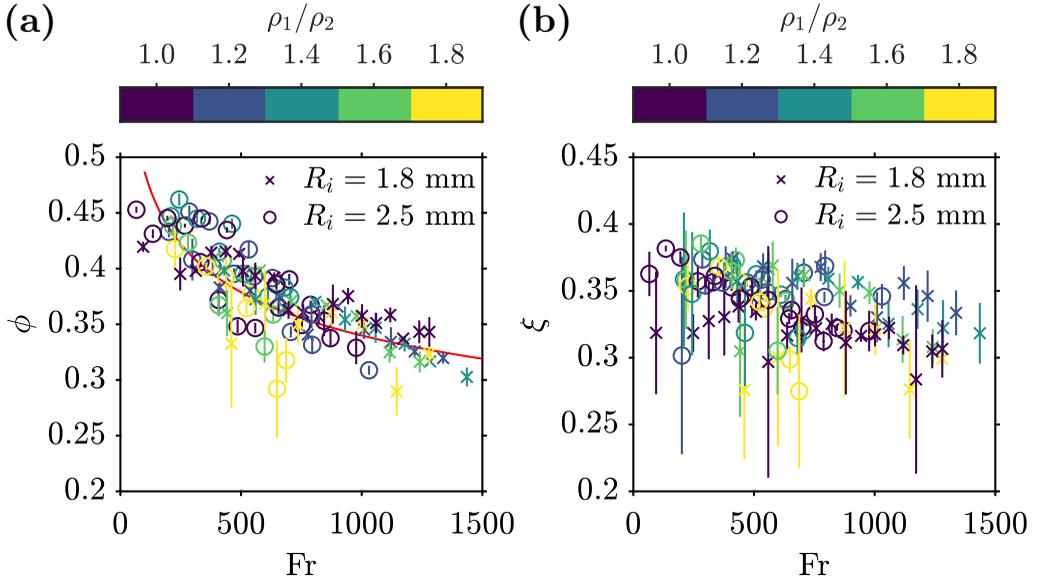


Figure 19: Energy partitioning parameter ϕ (a) and crater kinetic energy correction parameter ξ (b), as a function of the Froude number Fr . The solid line gives the best-fit power-law scaling (equation B 1). Colours scale as the density ratio ρ_1/ρ_2 . Circles and crosses correspond, respectively, to large and small drop size series.

1011 increases. This may be explained by a change in the energy balance between the crater
 1012 energy and the crown energy (Olevson 1969). As the impactor inertia increases, the relative
 1013 importance of the surface energy of the crater and the crown decreases, while the potential
 1014 energy of the crater and the kinetic energy of the crown increases, resulting in a global
 1015 increase of the crown energy to the expense of the crater. According to Olevson (1969), the
 1016 energy within the crown increases with Fr faster than the energy within the crater, which
 1017 would imply that ϕ is a decreasing function of Fr . The drop deformation upon impact may
 1018 also increase with impactor inertia and with it the energy required for this deformation,
 1019 decreasing to this extent the energy delivered to the pool.

1020 Figure 19b shows the kinetic energy correction parameter ξ , as a function of the Froude
 1021 number. It accounts for the difference between the deliberately simplified velocity potential
 1022 used in the model 4.3 and the true flow. Since the crater boundary is not hemispherical and
 1023 the crown is necessarily generated by a tangential velocity field, the true velocity potential
 1024 cannot be purely radial (Engel 1967; Bisighini *et al.* 2010). Parameter ξ is very likely a
 1025 function of time, but we assume it to be constant. In our experiments, the kinetic energy
 1026 correction parameter is smaller than unity with $\xi = 0.34 \pm 0.03$. This means that the velocity
 1027 model overestimates the crater kinetic energy in the energy balance. We do not observe any
 1028 resolvable trend between ξ , Fr and ρ_1/ρ_2 .

1029 Appendix C. Buoyancy prefactor and drag coefficient

1030 Figure 20 shows the fitted buoyancy prefactor β and drag coefficient C for each experiment.
 1031 The mean value of the buoyancy prefactor is $\beta = 0.32 \pm 0.17$. This is smaller than the
 1032 value $\beta = 1$ found by Dimonte (2000) for plane layer experiments. This difference may come
 1033 from the hemispherical shape or the finite thickness of the dense layer in our experiments.
 1034 Given error bars, we find that β shows no resolvable trend with the density ratio or the Froude
 1035 number. However, variations may exist and require further investigations.

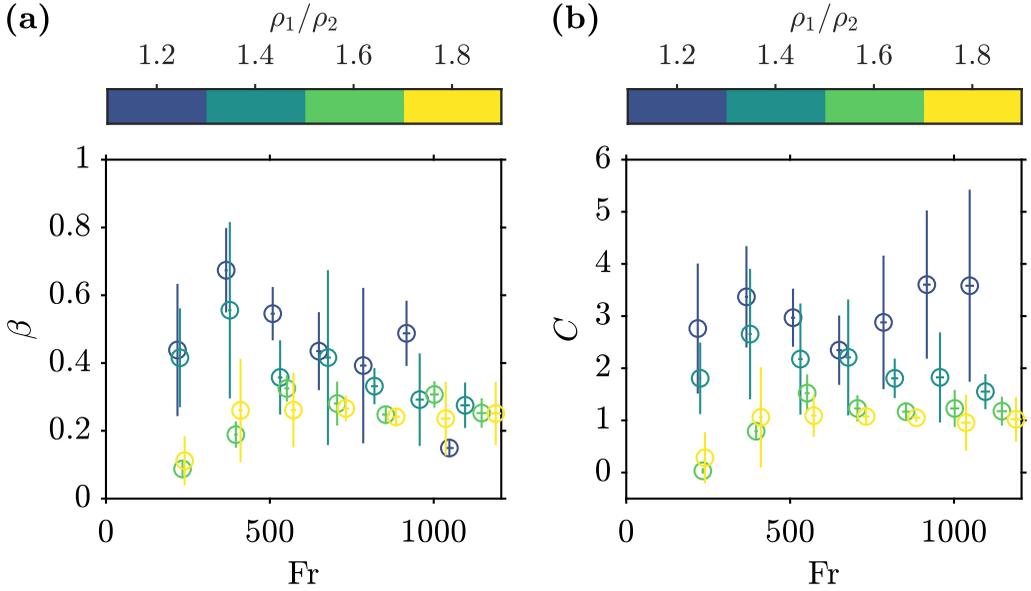


Figure 20: Buoyancy prefactor β (a) and drag coefficient C (b), as a function of the Froude number Fr . Colours scale as the density ratio ρ_1/ρ_2 .

1036 The mean value of the drag coefficient is $C = 1.9 \pm 1.1$. This value agrees with the value
 1037 $C = 2.5 \pm 0.6$ obtained for constant, variable and impulsive accelerations of a plane mixing
 1038 layer (Dimonte 2000). We find that C decreases when the density ratio increases (figure
 1039 20b). The mean value of the drag coefficient at $\rho_1/\rho_2 = 1.2$ and $\rho_1/\rho_2 = 1.8$ is, respectively,
 1040 $C = 3.1 \pm 1.1$ and $C = 1.1 \pm 0.4$. As for the buoyancy prefactor, C shows no resolvable trend
 1041 with the Froude number.

1042 In order to further compare the values of β and C with the results of (Dimonte 2000), we
 1043 now consider two simplified end-member acceleration histories: a constant acceleration and
 1044 an impulsive acceleration.

1045 In the case of a plane layer under constant acceleration \ddot{R} , the solution to the buoyancy–
 1046 drag equation 5.4 is $h = \alpha(\Delta\rho/\bar{\rho})\ddot{R}t^2$ (e.g. Dimonte 2000), where α is an empirical prefactor
 1047 and $\bar{\rho} = (\rho_1 + \rho_2)/2$. The parameter α can be expressed as a function of the parameters β
 1048 and C by looking for a solution of equation 5.4. Taking $u' = \dot{h}$, as appropriate for a plane
 1049 layer, solving gives

$$1050 \quad \alpha = \frac{\beta}{2 + 8C \frac{\rho_2}{\rho_1 + \rho_2}}. \quad (\text{C } 1)$$

1051 Figure 21a shows α , calculated for each experiment and compares the results to the
 1052 homogeneous buoyancy–drag model of Dimonte (2000) for $C = 1$, $C = 2$ and $C = 3$. The
 1053 mean value $\alpha = 0.04 \pm 0.01$ in our experiments is smaller than the values $\alpha = 0.05$ – 0.07 that
 1054 are obtained at the same density ratio with immiscible fluids and under constant acceleration
 1055 (Dimonte & Schneider 2000). The observed values of α are also smaller than the predictions
 1056 from the homogeneous model of Dimonte (2000), in particular at large density ratios. For
 1057 example, at $\rho_1/\rho_2 = 1.8$ the drag coefficient found in our experiments is approximately
 1058 $C = 1$ (figure 20b), which leads to an overestimate of α by a factor 2 in figure 21a. This may
 1059 be a consequence of the variable acceleration, but also of the spherical interface, miscibility
 1060 and the finite thickness of the dense layer.

1061 In the case of an impulsive acceleration, the buoyancy term in equation 5.4 is negligible

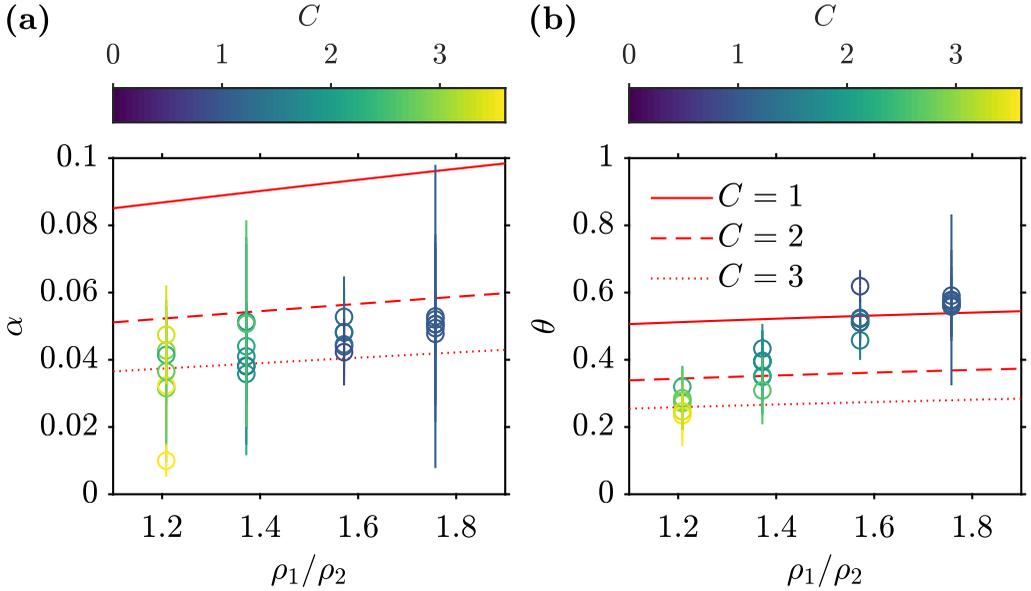


Figure 21: Constant acceleration prefactor α (a) and impulsive acceleration exponent θ (b) as a function of the density ratio ρ_1/ρ_2 . Colours scale as the experimental drag coefficient C . Solid lines, dashed lines and dotted lines correspond to the homogeneous buoyancy–drag model of Dimonte (2000) for $C = 1$, $C = 2$ and $C = 3$, respectively.

1062 since $\ddot{R} = 0$. Neglecting geometrical effects, *i.e.* $u' = \dot{h}$ and assuming that $\bar{\rho} = (\rho_1 + \rho_2)/2$,
 1063 the solution is then given by $h = h_0\tau^\theta$, where $\tau = u'_0 t / \theta h_0 + 1$ and h_0 and u'_0 are initial values
 1064 (*e.g.* Dimonte 2000). The exponent is then

$$1065 \quad \theta = \frac{1}{1 + 2C \frac{\rho_2}{\rho_1 + \rho_2}}. \quad (\text{C } 2)$$

1066 Figure 21b shows θ , calculated for each experiments and compares the results to the
 1067 buoyancy–drag model of Dimonte (2000). The mean value $\theta = 0.4 \pm 0.1$ is close to the
 1068 values $\theta = 0.2 - 0.3$ that are obtained at the same density ratio between immiscible fluids and
 1069 under an impulsive acceleration (Dimonte & Schneider 2000). The exponent θ increases with
 1070 the density ratio, consistently with the homogeneous model of Dimonte (2000) estimated at
 1071 consistent values of the drag coefficient C (figure 20b). Since the acceleration of the crater
 1072 evolves as $t^{-8/5}$, the acceleration is relatively close to be impulsive, explaining the good
 1073 agreement between our experiments and the impulsive acceleration model.

REFERENCES

- 1074 AGBAGLAH, GILOU, JOSSEAND, CHRISTOPHE & ZALESKI, STÉPHANE 2013 Longitudinal instability of a liquid
 1075 rim. *Physics of Fluids* **25** (2), 022103.
- 1076 AGNOR, CRAIG & ASPHAUG, ERIK 2004 Accretion Efficiency during Planetary Collisions. *The Astrophysical*
 1077 *Journal* **613** (2), L157–L160.
- 1078 AGNOR, CRAIG B., CANUP, ROBIN M. & LEVISON, HAROLD F. 1999 On the Character and Consequences of
 1079 Large Impacts in the Late Stage of Terrestrial Planet Formation. *Icarus* **142** (1), 219–237.
- 1080 ALLEN, ROBERT FRANCIS 1975 The role of surface tension in splashing. *Journal of Colloid and Interface*
 1081 *Science* **51** (2), 350–351.
- 1082 ARNETT, W. DAVID, BAHCALL, JOHN N., KIRSHNER, ROBERT P. & WOOSLEY, STANFORD E. 1989 Supernova
 1083 1987A. *Annual Review of Astronomy and Astrophysics* **27** (1), 629–700.
- 1084 BADRO, JAMES, AUBERT, JULIEN, HIROSE, KEI, NOMURA, RYUICHI, BLANCHARD, INGRID, BORENSZTAJN,

- 1085 STEPHAN & SIEBERT, JULIEN 2018 Magnesium Partitioning Between Earth's Mantle and Core and its
1086 Potential to Drive an Early Exsolution Geodynamo. *Geophysical Research Letters* **45** (24), 13,240–
1087 13,248.
- 1088 BADRO, JAMES, BRODHOLT, JOHN P., PIET, HÉLÈNE, SIEBERT, JULIEN & RYERSON, FREDERICK J. 2015 Core
1089 formation and core composition from coupled geochemical and geophysical constraints. *Proceedings*
1090 *of the National Academy of Sciences* **112** (40), 12310–12314.
- 1091 BALAKRISHNAN, KAUSHIK & MENON, SURESH 2011 Characterization of the Mixing Layer Resulting from
1092 the Detonation of Heterogeneous Explosive Charges. *Flow, Turbulence and Combustion* **87** (4),
1093 639–671.
- 1094 BELL, GEORGE I 1951 Taylor instability on cylinders and spheres in the small amplitude approximation. *Los*
1095 *Alamos Scientific Laboratory, Los Alamos, NM, Technical Report No. LA-1321* .
- 1096 BERBEROVIĆ, EDIN, VAN HINSBERG, NILS P., JAKIRLIĆ, SUAD, ROISMAN, ILIA V. & TROPEA, CAMERON 2009
1097 Drop impact onto a liquid layer of finite thickness: Dynamics of the cavity evolution. *Physical Review*
1098 *E* **79** (3), 036306.
- 1099 BISIGHINI, ALFIO, COSSALI, GIANPIETRO E., TROPEA, CAMERON & ROISMAN, ILIA V. 2010 Crater evolution
1100 after the impact of a drop onto a semi-infinite liquid target. *Physical Review E* **82** (3), 036319.
- 1101 BLANKEN, NATHAN, SALEEM, MUHAMMAD SAEED, THORAVAL, MARIE-JEAN & ANTONINI, CARLO 2021
1102 Impact of compound drops: A perspective. *Current Opinion in Colloid & Interface Science* **51**,
1103 101389.
- 1104 CANUP, ROBIN M. 2004 Simulations of a late lunar-forming impact. *Icarus* **168** (2), 433–456.
- 1105 CANUP, ROBIN M. 2012 Forming a Moon with an Earth-like Composition via a Giant Impact. *Science*
1106 **338** (6110), 1052–1055.
- 1107 CHAMBERS, J. 2010 Terrestrial Planet Formation. In *Exoplanets* (ed. S. Seager), pp. 297–317. Tucson,
1108 Arizona, USA: University of Arizona Press.
- 1109 CHANDRASEKHAR, S. 1955 The character of the equilibrium of an incompressible fluid sphere of variable
1110 density and viscosity subject to radial acceleration. *The Quarterly Journal of Mechanics and Applied*
1111 *Mathematics* **8** (1), 1–21.
- 1112 CHANDRASEKHAR, S. 1961 *Hydrodynamic and Hydromagnetic Stability*. Oxford: Clarendon.
- 1113 CORGNE, ALEXANDRE, KESHAV, SHANTANU, FEI, YINGWEI & McDONOUGH, WILLIAM F. 2007 How much
1114 potassium is in the Earth's core? New insights from partitioning experiments. *Earth and Planetary*
1115 *Science Letters* **256** (3), 567–576.
- 1116 CUK, MATIJA & STEWART, SARAH T. 2012 Making the Moon from a Fast-Spinning Earth: A Giant Impact
1117 Followed by Resonant Despinning. *Science* **338** (6110), 1047–1052.
- 1118 DALZIEL, S. B., LINDEN, P. F. & YOUNGS, D. L. 1999 Self-similarity and internal structure of turbulence
1119 induced by Rayleigh–Taylor instability. *Journal of Fluid Mechanics* **399**, 1–48.
- 1120 DEGUEN, RENAUD, LANDEAU, MAYLIS & OLSON, PETER 2014 Turbulent metal–silicate mixing, fragmentation,
1121 and equilibration in magma oceans. *Earth and Planetary Science Letters* **391**, 274–287.
- 1122 DEGUEN, RENAUD, OLSON, PETER & CARDIN, PHILIPPE 2011 Experiments on turbulent metal–silicate mixing
1123 in a magma ocean. *Earth and Planetary Science Letters* **310** (3), 303–313.
- 1124 DIMONTE, GUY 1999 Nonlinear evolution of the Rayleigh–Taylor and Richtmyer–Meshkov instabilities.
1125 *Physics of Plasmas* **6** (5), 2009–2015.
- 1126 DIMONTE, GUY 2000 Spanwise homogeneous buoyancy-drag model for Rayleigh–Taylor mixing and
1127 experimental evaluation. *Physics of Plasmas* **7** (6), 2255–2269.
- 1128 DIMONTE, GUY & SCHNEIDER, MARILYN 2000 Density ratio dependence of Rayleigh–Taylor mixing for
1129 sustained and impulsive acceleration histories. *Physics of Fluids* **12** (2), 304–321.
- 1130 EMMONS, H. W., CHANG, C. T. & WATSON, B. C. 1960 Taylor instability of finite surface waves. *Journal of*
1131 *Fluid Mechanics* **7** (2), 177–193.
- 1132 ENGEL, OLIVE G. 1961 Collisions of Liquid Drops With Liquids. *Tech. Rep.*. National Bureau of Standards,
1133 Washington, D. C.
- 1134 ENGEL, OLIVE G. 1962 Collisions of Liquid Drops With Liquids. Part 2 - Crater Depth in Fluid Impact.
1135 *Tech. Rep.*. National Bureau of Standards, Gaithersburg, MD.
- 1136 ENGEL, OLIVE G. 1966 Crater Depth in Fluid Impacts. *Journal of Applied Physics* **37** (4), 1798–1808.
- 1137 ENGEL, OLIVE G. 1967 Initial Pressure, Initial Flow Velocity, and the Time Dependence of Crater Depth in
1138 Fluid Impacts. *Journal of Applied Physics* **38** (10), 3935–3940.
- 1139 FEDORCHENKO, ALEXANDER I. & WANG, AN-BANG 2004 On some common features of drop impact on liquid
1140 surfaces. *Physics of Fluids* **16** (5), 1349–1365.
- 1141 FISCHER, REBECCA A., NAKAJIMA, YOICHI, CAMPBELL, ANDREW J., FROST, DANIEL J., HARRIES, DENNIS,

- 1142 LANGENHORST, FALKO, MIYAJIMA, NOBUYOSHI, POLLOK, KILIAN & RUBIE, DAVID C. 2015 High
1143 pressure metal–silicate partitioning of Ni, Co, V, Cr, Si, and O. *Geochimica et Cosmochimica Acta*
1144 **167**, 177–194.
- 1145 GIELEN, MARISE V., SLEUTEL, PASCAL, BENSCHOP, JOS, RIEPEN, MICHEL, VORONINA, VICTORIA, VISSER,
1146 CLAAS WILLEM, LOHSE, DETLEF, SNOEIJER, JACCO H., VERSLUIS, MICHEL & GELDERBLOM, HANNEKE
1147 2017 Oblique drop impact onto a deep liquid pool. *Physical Review Fluids* **2** (8), 083602.
- 1148 HAYNES, WILLIAM M. 2016 *CRC Handbook of Chemistry and Physics*. CRC Press.
- 1149 HOLSAPPLE, K A 1993 The Scaling of Impact Processes in Planetary Sciences. *Annual Review of Earth and*
1150 *Planetary Sciences* **21** (1), 333–373.
- 1151 ICHIKAWA, H., LABROSSE, S. & KURITA, K. 2010 Direct numerical simulation of an iron rain in the magma
1152 ocean. *Journal of Geophysical Research* **115** (B01404), 1–12.
- 1153 JACOBS, J. W. & SHEELEY, J. M. 1996 Experimental study of incompressible Richtmyer–Meshkov instability.
1154 *Physics of Fluids* **8** (2), 405–415.
- 1155 JAIN, UTKARSH, JALAAL, MAZIYAR, LOHSE, DETLEF & VAN DER MEER, DEVARAJ 2019 Deep pool water-
1156 impacts of viscous oil droplets. *Soft Matter* **15** (23), 4629–4638.
- 1157 JALAAL, MAZIYAR, KEMPER, DAVE & LOHSE, DETLEF 2019 Viscoplastic water entry. *Journal of Fluid*
1158 *Mechanics* **864**, 596–613.
- 1159 JEANS, JAMES HOPWOOD 1923 The propagation of earthquake waves. *Proceedings of the Royal Society of*
1160 *London. Series A, Containing Papers of a Mathematical and Physical Character* **102** (718), 554–574.
- 1161 KARATO, SHUN-ICHIRO & MURTHY, VR. 1997 Core formation and chemical equilibrium in the Earth—I.
1162 Physical considerations. *Physics of the Earth and Planetary Interiors* **100** (1), 61–79.
- 1163 KELLER, JOSEPH B. & KOLODNER, IGNACE 1954 Instability of Liquid Surfaces and the Formation of Drops.
1164 *Journal of Applied Physics* **25** (7), 918–921.
- 1165 KENDALL, JORDAN D. & MELOSH, H. J. 2016 Differentiated planetesimal impacts into a terrestrial magma
1166 ocean: Fate of the iron core. *Earth and Planetary Science Letters* **448**, 24–33.
- 1167 KLEINE, T., MÜNKER, C., MEZGER, K. & PALME, H. 2002 Rapid accretion and early core formation on
1168 asteroids and the terrestrial planets from Hf–W chronometry. *Nature* **418** (6901), 952.
- 1169 KRECHETNIKOV, ROUSLAN & HOMS, GEORGE M. 2009 Crown-forming instability phenomena in the drop
1170 splash problem. *Journal of Colloid and Interface Science* **331** (2), 555–559.
- 1171 LABROSSE, S., HERNLUND, J. W. & COLTICE, N. 2007 A crystallizing dense magma ocean at the base of the
1172 Earth's mantle. *Nature* **450** (7171), 866.
- 1173 LANDEAU, M., DEGUEN, R. & OLSON, P. 2014 Experiments on the fragmentation of a buoyant liquid volume
1174 in another liquid. *Journal of Fluid Mechanics* **749**, 478–518.
- 1175 LANDEAU, MAYLIS, DEGUEN, RENAUD, PHILLIPS, DOMINIC, NEUFELD, JEROME A., LHERM, VICTOR &
1176 DALZIEL, STUART B. 2021 Metal–silicate mixing by large Earth-forming impacts. *Earth and Planetary*
1177 *Science Letters* **564**, 116888.
- 1178 LENG, LIOW JONG 2001 Splash formation by spherical drops. *Journal of Fluid Mechanics* **427**, 73–105.
- 1179 LHERM, V. & DEGUEN, R. 2018 Small-Scale Metal/Silicate Equilibration During Core Formation: The
1180 Influence of Stretching Enhanced Diffusion on Mixing. *Journal of Geophysical Research: Solid*
1181 *Earth* **123** (12), 10,496–10,516.
- 1182 LHUISSIER, H., SUN, C., PROSPERETTI, A. & LOHSE, D. 2013 Drop Fragmentation at Impact onto a Bath of
1183 an Immiscible Liquid. *Physical Review Letters* **110** (26), 264503.
- 1184 LI, JIE & AGEE, CARL B. 1996 Geochemistry of mantle–core differentiation at high pressure. *Nature*
1185 **381** (6584), 686.
- 1186 LIN, HAO, STOREY, BRIAN D. & SZERI, ANDREW J. 2002 Rayleigh–Taylor instability of violently collapsing
1187 bubbles. *Physics of Fluids* **14** (8), 2925–2928.
- 1188 LINDEN, P. F., REDONDO, J. M. & YOUNGS, D. L. 1994 Molecular mixing in Rayleigh–Taylor instability.
1189 *Journal of Fluid Mechanics* **265**, 97–124.
- 1190 LINDL, JOHN 1998 *Inertial Confinement Fusion: The Quest for Ignition and Energy Gain Using Indirect*
1191 *Drive*. American Inst. of Physics.
- 1192 LOMBARDINI, M., PULLIN, D. I. & MEIRON, D. I. 2014 Turbulent mixing driven by spherical implosions.
1193 Part 1. Flow description and mixing-layer growth. *Journal of Fluid Mechanics* **748**, 85–112.
- 1194 MACKLIN, W. C. & METAXAS, G. J. 1976 Splashing of drops on liquid layers. *Journal of Applied Physics*
1195 **47** (9), 3963–3970.
- 1196 MELOSH, H. J. 1989 *Impact Cratering: A Geologic Process*. New York: Oxford University Press.
- 1197 MESHKOV, E. E. 1969 Instability of the interface of two gases accelerated by a shock wave. *Fluid Dynamics*
1198 **4** (5), 101–104.

- 1199 MIKAEILIAN, KARNIG O. 1990 Rayleigh-Taylor and Richtmyer-Meshkov instabilities and mixing in stratified
1200 spherical shells. *Physical Review A* **42** (6), 3400–3420.
- 1201 MIKAEILIAN, KARNIG O. 2016 Viscous Rayleigh-Taylor instability in spherical geometry. *Physical Review E*
1202 **93** (2), 023104.
- 1203 MONTEUX, J., JELLINEK, A. M. & JOHNSON, C. L. 2011 Why might planets and moons have early dynamos?
1204 *Earth and Planetary Science Letters* **310** (3), 349–359.
- 1205 MORARD, G. & KATSURA, T. 2010 Pressure–temperature cartography of Fe–S–Si immiscible system.
1206 *Geochimica et Cosmochimica Acta* **74** (12), 3659–3667.
- 1207 MORTON, DAVID, RUDMAN, MURRAY & JONG-LENG, LIOW 2000 An investigation of the flow regimes resulting
1208 from splashing drops. *Physics of Fluids* **12** (4), 747–763.
- 1209 NAKAJIMA, MIKI, GOLABEK, GREGOR J., WÜNNEMANN, KAI, RUBIE, DAVID C., BURGER, CHRISTOPH,
1210 MELOSH, HENRY J., JACOBSON, SETH A., MANSKE, LUKAS & HULL, SCOTT D. 2021 Scaling laws
1211 for the geometry of an impact-induced magma ocean. *Earth and Planetary Science Letters* **568**,
1212 116983.
- 1213 OKAWA, TOMIO, SHIRAISHI, TAKUYA & MORI, TOSHIAKI 2006 Production of secondary drops during the
1214 single water drop impact onto a plane water surface. *Experiments in Fluids* **41** (6), 965.
- 1215 OLEVSON, K. L. R. 1969 Energy balances for transient water craters. *US Geol. Surv. Prof. Pap D* (650),
1216 189–194.
- 1217 ORON, D., ARAZI, L., KARTOON, D., RIKANATI, A., ALON, U. & SHVARTS, D. 2001 Dimensionality dependence
1218 of the Rayleigh–Taylor and Richtmyer–Meshkov instability late-time scaling laws. *Physics of Plasmas*
1219 **8** (6), 2883–2889.
- 1220 PLESSET, M. S. 1954 On the Stability of Fluid Flows with Spherical Symmetry. *Journal of Applied Physics*
1221 **25** (1), 96–98.
- 1222 PROSPERETTI, ANDREA 1977 Viscous effects on perturbed spherical flows. *Quarterly of Applied Mathematics*
1223 **34** (4), 339–352.
- 1224 PROSPERETTI, A & OGUZ, H N 1993 The Impact of Drops on Liquid Surfaces and the Underwater Noise of
1225 Rain. *Annual Review of Fluid Mechanics* **25** (1), 577–602.
- 1226 RAY, BAHNI, BISWAS, GAUTAM & SHARMA, ASHUTOSH 2015 Regimes during liquid drop impact on a liquid
1227 pool. *Journal of Fluid Mechanics* **768**, 492–523.
- 1228 RAYLEIGH, JOHN WILLIAM STRUTT 1899 *Scientific Papers (Vol. 2: 1881-1887)*. Cambridge University Press.
- 1229 REIN, MARTIN 1993 Phenomena of liquid drop impact on solid and liquid surfaces. *Fluid Dynamics Research*
1230 **12** (2), 61–93.
- 1231 RICHTMYER, ROBERT D. 1960 Taylor instability in shock acceleration of compressible fluids. *Communications*
1232 *on Pure and Applied Mathematics* **13** (2), 297–319.
- 1233 RIGHTER, K. 2011 Prediction of metal–silicate partition coefficients for siderophile elements: An update and
1234 assessment of PT conditions for metal–silicate equilibrium during accretion of the Earth. *Earth and*
1235 *Planetary Science Letters* **304** (1), 158–167.
- 1236 RUBIE, DAVID C., FROST, DANIEL J., MANN, UTE, ASAHARA, YUKI, NIMMO, FRANCIS, TSUNO, KYUSEI,
1237 KEGLER, PHILIP, HOLZHEID, ASTRID & PALME, HERBERT 2011 Heterogeneous accretion, composition
1238 and core–mantle differentiation of the Earth. *Earth and Planetary Science Letters* **301** (1), 31–42.
- 1239 RUBIE, D. C., MELOSH, H. J., REID, J. E., LIEBSKE, C. & RIGHTER, K. 2003 Mechanisms of metal–silicate
1240 equilibration in the terrestrial magma ocean. *Earth and Planetary Science Letters* **205** (3), 239–255.
- 1241 RUBIE, D. C., NIMMO, F. & MELOSH, H. J. 2015 9.03 - Formation of the Earth's Core. In *Treatise on*
1242 *Geophysics (Second Edition)* (ed. Gerald Schubert), pp. 43–79. Oxford: Elsevier.
- 1243 RUDGE, JOHN F., KLEINE, THORSTEN & BOURDON, BERNARD 2010 Broad bounds on Earth's accretion and
1244 core formation constrained by geochemical models. *Nature Geoscience* **3** (6), 439.
- 1245 SCHERSTÉN, ANDERS, ELLIOTT, TIM, HAWKESWORTH, CHRIS, RUSSELL, SARA & MASARIK, JOZEF 2006 Hf–W
1246 evidence for rapid differentiation of iron meteorite parent bodies. *Earth and Planetary Science Letters*
1247 **241** (3), 530–542.
- 1248 SCHMIDT, WOLFRAM 2006 From tea kettles to exploding stars. *Nature Physics* **2** (8), 505–506.
- 1249 SHOEMAKER, EUGENE M. 1961 Interpretation of Lunar Craters. In *Physics and Astronomy of the Moon* (ed.
1250 Zdenek Kopal), pp. 283–359. Academic Press.
- 1251 SIEBERT, JULIEN, BADRO, JAMES, ANTONANGELI, DANIELE & RYERSON, FREDERICK J. 2012 Metal–silicate
1252 partitioning of Ni and Co in a deep magma ocean. *Earth and Planetary Science Letters* **321–322**,
1253 189–197.
- 1254 SIEBERT, JULIEN, CORGNE, ALEXANDRE & RYERSON, FREDERICK J. 2011 Systematics of metal–silicate

- 1255 partitioning for many siderophile elements applied to Earth's core formation. *Geochimica et*
1256 *Cosmochimica Acta* **75** (6), 1451–1489.
- 1257 SOLOMATOV, V. 2015 9.04 - Magma Oceans and Primordial Mantle Differentiation. In *Treatise on Geophysics*
1258 (*Second Edition*) (ed. Gerald Schubert), pp. 81–104. Oxford: Elsevier.
- 1259 STEVENSON, D. J. 1990 Fluid dynamics of core formation. In *Origin of the Earth* (ed. Horton E. Newsom &
1260 John H. Jones), pp. 231–249. New York: Oxford University Press.
- 1261 SUN, YICHENG, ZHOU, HUIQUN, YIN, KUN, ZHAO, MENGQI, XU, SHIJIN & LU, XIANCAI 2018 Transport
1262 Properties of Fe₂SiO₄ Melt at High Pressure From Classical Molecular Dynamics: Implications for
1263 the Lifetime of the Magma Ocean. *Journal of Geophysical Research: Solid Earth* **123** (5), 3667–3679.
- 1264 TAKITA, HARUNA & SUMITA, IKURO 2013 Low-velocity impact cratering experiments in a wet sand target.
1265 *Physical Review E* **88** (2), 022203.
- 1266 TANIGUCHI, H. 1988 Surface tension of melts in the system CaMgSi₂O₆-CaAl₂Si₂O₈ and its structural
1267 significance. *Contributions to Mineralogy and Petrology* **100** (4), 484–489.
- 1268 TAYLOR, GEOFFREY INGRAM 1950 The instability of liquid surfaces when accelerated in a direction
1269 perpendicular to their planes. I. *Proceedings of the Royal Society of London. Series A. Mathematical*
1270 *and Physical Sciences* **201** (1065), 192–196.
- 1271 THOMAS, V. A. & KARES, R. J. 2012 Drive Asymmetry and the Origin of Turbulence in an ICF Implosion.
1272 *Physical Review Letters* **109** (7), 075004.
- 1273 TONKS, W. BRIAN & MELOSH, H. JAY 1993 Magma ocean formation due to giant impacts. *Journal of*
1274 *Geophysical Research: Planets* **98** (E3), 5319–5333.
- 1275 TRYGGVASON, GRÉTAR 1988 Numerical simulations of the Rayleigh-Taylor instability. *Journal of*
1276 *Computational Physics* **75** (2), 253–282.
- 1277 VILLERMAUX, EMMANUEL 2020 Fragmentation versus Cohesion. *Journal of Fluid Mechanics* **898**, P1.
- 1278 WACHEUL, JEAN-BAPTISTE & LE BARS, MICHAEL 2018 Experiments on fragmentation and thermo-chemical
1279 exchanges during planetary core formation. *Physics of the Earth and Planetary Interiors* **276**, 134–
1280 144.
- 1281 WACHEUL, JEAN-BAPTISTE, LE BARS, MICHAEL, MONTEUX, JULIEN & AURNOU, JONATHAN M. 2014
1282 Laboratory experiments on the breakup of liquid metal diapirs. *Earth and Planetary Science Letters*
1283 **403**, 236–245.
- 1284 WALSH, AMANDA M., HOLLOWAY, KRISTI E., HABDAS, PIOTR & DE BRUYN, JOHN R. 2003 Morphology and
1285 Scaling of Impact Craters in Granular Media. *Physical Review Letters* **91** (10), 104301.
- 1286 WILLE, G., MILLOT, F. & RIFFLET, J. C. 2002 Thermophysical Properties of Containerless Liquid Iron up to
1287 2500 K. *International Journal of Thermophysics* **23** (5), 1197–1206.
- 1288 WILLIAMS, JEAN-PIERRE & NIMMO, FRANCIS 2004 Thermal evolution of the Martian core: Implications for
1289 an early dynamo. *Geology* **32** (2), 97–100.
- 1290 WOOD, BERNARD J., WALTER, MICHAEL J. & WADE, JONATHAN 2006 Accretion of the Earth and segregation
1291 of its core. *Nature* **441** (7095), 825.
- 1292 WORTHINGTON, A. M. 1895 *The Splash of a Drop*. London, UK: London, S.P.C.K.
- 1293 YOUNGS, DAVID L. & WILLIAMS, ROBIN J. R. 2008 Turbulent mixing in spherical implosions. *International*
1294 *Journal for Numerical Methods in Fluids* **56** (8), 1597–1603.
- 1295 ZHANG, LI V., BRUNET, PHILIPPE, EGGERS, JENS & DEEGAN, ROBERT D. 2010 Wavelength selection in the
1296 crown splash. *Physics of Fluids* **22** (12), 122105.
- 1297 ZHOU, YE 2017 Rayleigh–Taylor and Richtmyer–Meshkov instability induced flow, turbulence, and mixing.
1298 II. *Physics Reports* **723–725**, 1–160.