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1 **Numerical investigation of spreading time in droplet impact with**
2 **spherical surfaces: from physical analysis to data-driven**
3 **prediction model**

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5
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ABSTRACT

Spreading time, the time that an impacting droplet attains the maximum wetting area on a solid surface, plays a critical role in many engineering applications particularly where heat transfer or chemical reactions are involved. Although the impact dynamics of a droplet significantly differ across the different spreading regimes depending on various collision parameters, it still remains unclear (i) how the spreading time changes for each spreading regime, and (ii) how the target curvature can affect the spreading time. In the present study, the spreading time during droplet impact on a spherical target is systematically studied at the three different spreading regimes for a wide range of impact parameters (Weber number, equilibrium contact angle, Ohnesorge number and droplet-to-target size ratio). The changes of spreading time depending on the impact parameters and underlying physical mechanisms are analyzed in detail at the level of three different spreading regimes. Our results show that the spreading time, proper time scales, dominant impact parameters and associated physical behaviors all significantly and non-linearly change across the three spreading regimes. An improved prediction model for the spreading time is also proposed for each regime, which is now based on only the controllable variables and has an explicit form. The effect of target curvature on the spreading time is further investigated, and finally, a data-driven prediction model is proposed to represent the complicated and non-linear nature of the spreading time broadly across the three spreading regimes.

Keywords: droplet impact; spreading time; multiphase flow simulation; prediction model; data-driven analysis

47 I. INTRODUCTION

48

49 Collisions of droplets with **solid surfaces** can be found not only in a broad range of
50 modern technologies but also in everyday life [1,2]. Since droplet collision phenomena can
51 play crucial roles in the performance of diverse engineering applications such as spray
52 cooling and chilling [3], spray drying [4], wet-scrubbers [5], fluid catalytic cracking (FCC)
53 [6], trickle-bed reactors [7], microfluidic-based encapsulation of biomaterials [8] and tablet
54 coating [9], to name a few examples, their collision processes have naturally been of great
55 interest across various fields of chemical, mechanical, pharmaceutical, food and bio
56 industries [1,9-11]. In general, the collision behavior is dominated by a complicated
57 competition and interaction among inertial, viscous, and capillary forces (sometimes
58 including the gravitational force as well) at very small length and time scales [12,13].

59 For droplet impact problems with a solid surface much effort has focused on
60 investigating the maximum spreading extent (normally characterized by the maximum
61 spreading diameter) [14-17] and the spreading time [18-21]. The former measures “how
62 much a droplet can spread over a target surface”, while the latter quantifies “how long it takes
63 to reach the maximum spreading state”. Both quantities usually play critical roles in
64 determining process performance and efficiency particularly in the context of controlling the
65 droplet deposition process [1,12,19]. For example, the maximum spreading extent of a
66 catalyst particle significantly affects the catalyst efficiency as well as the system performance
67 of trickle-bed reactors [7] and the spreading time also plays an important role in collision
68 processes where heat transfer (e.g., droplet solidification) [14] and/or chemical reactions
69 (e.g., cracking reaction in the FCC system) [22] are involved.

70 After Chandra and Avedisian’s work [23] which has initiated a series of energy-balance-

71 based theoretical approaches for estimating the maximum spreading diameter of droplets
72 numerous studies have been devoted to the spreading dynamics of droplets on a solid surface.
73 Based on the concept of the energy-balance between the collision instant and the maximum
74 spreading state Pasandideh-Fard *et al.* [24] proposed more refined length and time scales to
75 evaluate the energy loss in the boundary layer where the viscous dissipation occurs and their
76 approach has indeed provided a cornerstone for many later theoretical energy-balance-based
77 approaches [25,26]. This type of analysis was also widely utilized for estimating the rebound
78 criteria of bouncing droplets on a non-wettable surface [25] as well as the maximum
79 spreading extent [23-26]. Afterward, Roisman [27] considered the full Navier-Stokes
80 equations to estimate the flow field inside the droplet and proposed a semi-empirical model
81 using the experimentally fitted data. Later, his model was experimentally demonstrated to be
82 one of the most accurate models for the maximum spreading extent for a wide range of
83 Weber number and surface wettability [14]. Wilderman *et al.* [17] presented an interesting
84 analysis for high-speed drop impact problems. They showed that roughly one-half of the
85 initial (impact) kinetic energy can be converted into surface energy owing to a universal head
86 loss regardless of detailed impact conditions and energy dissipation mechanisms. More
87 recently some other sources causing extra energy loss during the spreading stage of droplets
88 have been found to affect the maximal spreading extent, e.g., residual kinetic energy in the
89 form of vortical flow inside a droplet [16] or energy loss due to ‘interface relaxation’ near a
90 solid surface [18]. Although there is still a debate on the maximum spreading extent [28], in
91 general, scaling laws of $We^{1/4}$ derived from momentum conservation [16] and $We^{1/2}$ derived
92 from energy conservation [29] have been experimentally shown to hold for the inviscid
93 regime (for less viscous droplets) whereas a $Re^{1/5}$ rule [30] has worked well for a viscous
94 regime (for highly viscous droplets).

95 There have also been several attempts to examine spreading dynamics for droplet

96 collision on curved surfaces such as a spherical target. Bakshi *et al.* [31] experimentally
97 showed that the film thickness at the collision center undergoes three different phases: (i)
98 initial drop-deformation, (ii) inertia-dominated and (iii) viscosity-dominated. The
99 dimensionless film thickness profiles were collapsed on a single curve for the first and second
100 phases. Malgarinos *et al.* [32] used volume-of-fluid (VOF) simulation to study the problem of
101 drop collision on spherical targets. They investigated the boundary separating the coating
102 from the rebound regime, the film thickness, and the wetted area. Afterward, they further
103 performed extensive VOF simulations to study drop impact on a spherical particle including
104 phase change and reaction phenomena in a high temperature environment ($T \geq 800$ K) [33].
105 Zhang *et al.* presented two-dimensional [34] and three-dimensional [35] Lattice-Boltzmann
106 simulations. They examined the effects of Weber and Reynolds numbers on the film thickness
107 at the collision center, and their simulation results reproduced Bakshi *et al.*'s [31]
108 experimental observations well. Liang *et al.* [36] experimentally studied drops impacting on
109 wetted spherical targets, and they showed the drop-to-target curvature ratio played a critical
110 role in not only the spreading characteristics but also determining the splashing threshold.
111 Mitra *et al.* [37], Liu *et al.*, [38], and Khurana *et al.* [39] proposed energy-balance-based
112 prediction models for evaluating the spreading extent on a spherical surface and their models
113 reproduced their experimental observations well. Yoon and Shin [40] also proposed scaling
114 laws and an empirical correlation for the maximum spreading diameter of a droplet colliding
115 with a spherical surface for a wide range of liquid viscosity. More recently, Yoon *et al.* [41]
116 proposed a data-driven prediction model for the maximum spreading that can be applied to
117 both a flat surface and a spherical surface. Although these studies reviewed above have
118 provided useful relations to estimate the maximum spreading extent as well as have notably
119 shed light on our understanding of the physical mechanisms of spreading droplets, they have
120 mainly focused on the maximum spreading extent itself rather than on the spreading time.

121 We now review studies in the context of the spreading time which has been relatively
 122 less well-explored compared to the maximum spreading extent. Although it has long been
 123 understood that the contact time of a droplet is limited by the Rayleigh oscillation time [42]
 124 which is a function of the drop's density, diameter and surface tension coefficient, the energy-
 125 balance-based approaches, at the very early stage, simply approximated the spreading time τ
 126 ignoring the effect of surface tension. Chandra and Avedisian [23] considered τ to be $\tau = D_d /$
 127 V_{ini} , where D_d and V_{ini} are the initial droplet diameter and the impact velocity, respectively.
 128 This is indeed the advective characteristic time scale which has normally been used to
 129 nondimensionalize the physical time scale [31], and also can be interpreted as the required
 130 time for a droplet to reach zero height from its initial diameter, D_d , at a constant impact
 131 velocity V_{ini} . Pasandideh-Fard *et al.* [24] modified the spreading time to be $\tau = 8D_d / 3V_{ini}$
 132 based on a simple geometric assumption and mass conservation and their model has been
 133 widely used in many other theoretical studies [26,43]. Recently, Antonini and Amirfazli [14]
 134 experimentally showed that the spreading time τ is considerably affected by surface
 135 wettability. Lee *et al.* [44] also presented significant effects of liquid properties on τ . They
 136 replaced the initial drop diameter (D_d) in Pasandideh-Fard *et al.*'s model [24] with D_m (where
 137 D_m is the maximum spreading diameter), suggesting an empirical model, $\tau = b(D_m/V_{ini})$,
 138 where b is the ratio of the surface tension coefficient between a working droplet and a water
 139 droplet. Huang and Chen [18] and Wang *et al.* [45] also adopted a similar spreading time as
 140 the model of Lee *et al.* [44]. Lin *et al.* [19] presented a remarkable experimental study for the
 141 spreading time and systematically analyzed the effects of various collision parameters (e.g.,
 142 impact velocity, liquid properties and surface wettability) on τ for a very wide range of
 143 collision conditions. They also proposed a new empirical model for the spreading time as a
 144 function of Weber number, i.e., $\tau^* = 0.92We^{-0.43}$. Here, τ^* is the rescaled spreading time using
 145 a modified capillary time scale $(\rho D_m^3/8\sigma)^{0.5}$ based on the maximal spreading diameter, D_m ,

146 where ρ and σ are the density of a droplet and the surface tension coefficient, respectively. Du
147 *et al.* [21] also proposed a similar empirical model ($\tau^* = 1.47We^{-0.44}$) as Lin *et al.*'s model
148 [19]. More recently, Aksoy *et al.* [20] showed the effect of liquid viscosity on the spreading
149 time and proposed an empirical model as a function of both Weber and Reynolds number, i.e.,
150 $\tau^* = 2We^{-0.45}Re^{-0.09}$.

151 Even considering the numerous studies heretofore undertaken, it still remains difficult to
152 comprehensively understand how the spreading time of a droplet can be affected by various
153 collision parameters for a wide range of impact parameters. In particular, one can find
154 significant knowledge gaps as follows:

155

156 (i) It is well-known that drop spreading dynamics vary significantly across the
157 different spreading regimes. For example, a drop's physical behavior and its
158 maximum spreading extent change considerably across inviscid and viscous
159 regimes [16,29,30]. However, it is still unclear how the spreading time (τ)
160 changes for different spreading regimes, because the existing studies have usually
161 studied the global behavior of spreading time at a general level, not focused on
162 detailed changes and associated physical mechanisms at the level of each
163 different spreading regime.

164

165 (ii) Although the significant effects from impact parameters (e.g., We, Re, surface
166 wettability, etc.) on the spreading time have been evidently demonstrated by
167 many experiments [14,18-20,44], it is not easy yet to model the spreading time
168 using those impact parameters. All recent scaling laws (or empirical models)
169 reviewed above [18-21,44,45] have still been a function of the maximum
170 spreading diameter (D_m) which is usually unknown. Since the effects of impact

171 parameters are already reflected in D_m [14-17,25,27], D_m cannot be considered as
172 a controllable independent variable to model τ resulting in the limitation of direct
173 “physical interpretation” of relations between the impact parameters and the
174 spreading time. In addition, the existing models [18-21,44,45] cannot be
175 explicitly solved due to the multiple unknown variables (i.e., the maximum
176 spreading diameter and the spreading time) leading to inconvenience for users.

177
178 (iii) The effect of surface curvature on the spreading time has not been systematically
179 investigated so far. Although a few recent numerical studies [46,47] have
180 demonstrated that the surface curvature can affect the spreading time of the
181 droplet collision system, it is still unclear that how the spreading time changes
182 depending on the surface curvature for different spreading regimes. In addition,
183 to the best of our knowledge, there has not yet been a prediction model for the
184 spreading time which can consider the effect of surface curvature.

185
186 In this study, we expand our previous work [41] which examined the maximum
187 spreading extent, now focusing on the systematic investigation of the spreading time which is
188 another important design parameter in the droplet collision system. Based on our verified
189 computational framework, effects of important impact parameters (i.e., collision velocity,
190 surface wettability, liquid viscosity, surface tension and surface curvature) on the spreading
191 time are quantified for a wide range of impact parameters. In particular, detailed changes in
192 the spreading time and associated underlying mechanisms are analyzed for three different
193 spreading regimes to better understand the complicated physical behaviors. A new prediction
194 model for the spreading time using only controllable independent variables is presented for
195 each regime, and a data-driven prediction model for the entire spreading regime is also

196 proposed broadly covering all ranges of impact parameters considered in the present study.

197 The remainder of the current paper is organized as follows: Section II provides a brief
198 introduction of the simulation methods used herein. Section III quantitatively and
199 systematically examines the spreading time of a droplet collision system. Physical analysis
200 and prediction models are also presented in this section. The major findings are summarized
201 in Section IV.

202

203 II. COMPUTATIONAL FORMULATION

204

205 Since we primarily focus on understanding the physical characteristics of the spreading
206 time, and have utilized the same computational framework as in our previous studies
207 [40,41,48,49], we provide here a brief introduction of our numerical methods used in this
208 study rather than a fully detailed explanation. For more algorithmic details and relevant
209 techniques on our simulation methods, and for our various benchmark tests including grid
210 convergence characteristics, readers can refer to our previous work [40,41,48,49]. Note also
211 that our computational framework has been extensively applied to diverse droplet collision
212 cases for various types of solid surface including flat substrates [49], spherical objects
213 [40,41,48], cylindrical targets [50] and has also been streamlined (using an adaptive-mesh-
214 refinement approach [51]) and parallelized [52].

215 For incompressible flows the governing equations can be applied to all three phases
216 (solid, liquid and gas) and are solved on a fixed Cartesian (Eulerian) grid as a single-field
217 formulation:

218

$$219 \quad \nabla \cdot \mathbf{u} = 0 \quad (1)$$

220

$$\rho \left[\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} \right] = -\nabla P + \rho \mathbf{g} + \nabla \cdot \mu (\nabla \mathbf{u} + \nabla \mathbf{u}^T) + \mathbf{F} \quad (2)$$

222

223 where \mathbf{u} , P and \mathbf{g} are the velocity vector, the pressure, and the gravitational acceleration,
224 respectively. ρ is the density and μ is the viscosity. The surface tension force \mathbf{F} is considered
225 only at the phase (gas-liquid) interface and can be modeled by the following hybrid
226 representation based on the continuum-surface-force (CSF) formulation [53]:

227

$$\mathbf{F} = \sigma \kappa_H \nabla I \quad (3)$$

229

230 where κ_H is the curvature field. Note that I is the indicator function which has the
231 characteristics of the Heaviside function, varying from 0 (zero) in one phase (liquid droplet)
232 to 1 (one) in the other phase (ambient air).

233 The physical properties of both phases (i.e., the ambient air and the liquid droplet) can
234 also be assigned by using the indicator function I as follows:

235

$$\rho = \rho_d + (\rho_a - \rho_d) I \quad (4)$$

237

$$\mu = \mu_d + (\mu_a - \mu_d) I \quad (5)$$

239

240 Here, the subscript “d” denotes droplet, whereas “a” denotes air. Note that this common
241 approach has been broadly applied to various multiphase-flow simulations and more detailed
242 techniques dealing with the surface tension force and the physical properties near the gas-
243 liquid phase interface can be found in Unverdi and Tryggvason [54] and Brackbill *et al.* [53].

244 Our dynamic contact angle modeling technique is given in Appendix A, and further details of
245 the discretization of the governing equations, time-stepping, and other solution techniques
246 used in the present study can also be found in Yoon *et al.* [41] and Yoon and Shin [48].

247 One of the most essential algorithms for this type of multiphase-flow simulation is an
248 interface tracking method. Here we use the level contour reconstruction method (LCRM) [55-
249 57] to track the gas-liquid phase interface. The LCRM is a hybrid method benefitting from
250 two of the most well-established and popular interface tracking methods, i.e., level set [58]
251 and front tracking [54]. In the LCRM the phase interface is basically represented by using
252 (Lagrangian) moving marker elements (lines in 2D simulations and triangles in 3D
253 simulations, as in the original front tracking method) but the moving elements are re-meshed
254 using the distance function which is the key feature of the level set method. Such a
255 reconstruction procedure can allow us to avoid the algorithmic complexity of dealing with the
256 connectivity among the moving interface elements which is the most well-known
257 shortcoming of the original front tracking method. Note that each moving element can always
258 be implicitly (naturally) interconnected without any further artificial connection technique
259 because the reconstruction procedure is performed at each boundary between cells (i.e., at
260 cell faces) on the fixed Cartesian grid. Note also that the LCRM retains the tracking
261 capability of the original front tracking method (i.e., accurate representation of the phase
262 interface using moving markers) while avoiding the complex algorithmic difficulties in
263 handling the topology changes (e.g., deformation, pinch-off or coalescence of the phase
264 interface). For more details on the LCRM readers can refer to Shin and Juric [55,56].

265 Figure 1 illustrates the simulation geometry and boundary conditions used in this study.
266 All simulations are performed in a (two-dimensional) axi-symmetric domain. The axi-
267 symmetric boundary condition is set for the left boundary whereas open boundary conditions
268 are set for the right, upper and lower boundaries. The lengths of the physical domain are set

269 as $R_L = 7.5$ (in the radial r direction) and $Z_L = 20$ (in the axial z direction) times the droplet
 270 radius, respectively, which are sufficiently large compared to the size of the droplet. **Note that**
 271 **near the splashing threshold [59,60] the droplet rim may exhibit considerable asymmetric**
 272 **behavior along the circumferential direction (e.g., fingering or cusp of lamella [12,13]).**
 273 **Therefore, we set the simulation conditions sufficiently below the splashing threshold**
 274 **(specific conditions will be described later) to ensure that our axi-symmetric formulation can**
 275 **be a reasonable approach to reducing the necessary computational resources [41,51]. A**
 276 droplet impacts onto a stationary spherical **target** at an initial impact velocity, (V_{ini}) and D_s is
 277 the diameter of the **spherical target** where the subscript “s” denotes the **target surface**.
 278

279 **III. RESULTS AND DISCUSSION**

280
 281 The spreading dynamics of a droplet on a solid surface is primarily dominated by a
 282 complicated time-and-space dependent interplay among inertial, viscous, and capillary effects
 283 [12,13] and to characterize their relative importance, three non-dimensional variables, i.e., the
 284 Weber number ($We = \rho_d V_{ini}^2 D_d / \sigma$), the Reynolds number ($Re = \rho_d V_{ini} D_d / \mu_d$) and the
 285 Ohnesorge number [$Oh = \mu_d / (\rho_d \sigma D_d)^{0.5}$] have usually been used. Because Oh can also be
 286 represented by using We and Re (i.e., $Oh = We^{0.5} / Re$) the spreading dynamics can basically
 287 be characterized by two of these numbers. Here we use We and Oh as impact parameters with
 288 two other variables, i.e., the equilibrium contact angle (θ_{eqi}) and the droplet-to-**target** size
 289 ratio (Ω) which characterize surface wettability and the curvature of the **target surface**,
 290 respectively.

291 Three characteristic time scales are considered, i.e., the advective time, $t_a = D_d / V_{ini}$
 292 [30,61], the capillary time, $t_c = (\rho_d D_d^3 / 8\sigma)^{0.5}$ [16,62], and the viscous time. $t_v = \mu_d D_d / \sigma$ [16,62].
 293 The physical time scale, t , can be non-dimensionalized by the three characteristic time scales

294 above and the non-dimensional advective, capillary, and viscous time scales are denoted by
295 t_a^* ($t_a^* = t / t_a$), t_c^* ($t_c^* = t / t_c$), t_v^* ($t_v^* = t / t_v$), respectively.

296 To sufficiently cover the broad range of impact parameters, and to systematically
297 examine the effects of impact parameters on the spreading time, we consider a total of 2400
298 collision cases. The droplet diameter and the density are fixed at $D_d = 2$ mm and $\rho_d = 998.2$
299 kg/m³, respectively. The physical properties of the gas phase (ambient air) and the surface
300 tension coefficient are also fixed ($\rho_a = 1.2$ kg/m³, $\mu_a = 0.000018$ N s/m², and $\sigma = 0.0728$ N/m).

301 Four collision parameters, i.e., We , Oh , θ_{eqi} and Ω are varied by controlling the impact
302 velocity (V_{ini}), droplet viscosity (μ_d), equilibrium contact angle (θ_{eqi}) and the diameter of the
303 target (D_s), respectively. The splashing threshold is usually described by We and Oh [59,63],
304 and splashing can be observed more easily on a small spherical target compared to that on a
305 flat substrate [63]. Since splashing can occur near $We \sim 150 - 160$ if $\Omega \sim 1/4 - 1/2$ (for
306 millimetric water drops, $Oh \sim 0.0026$) [48,63], we consider a range of Weber number below
307 110 ($1 \leq We \leq 110$) to ensure that the collision outcomes would not fall into the splashing
308 region. Therefore, a total of 12 cases of We (i.e., $We = 1, 2, 4, 6, 8, 10, 20, 30, 50, 70, 90,$ and
309 110) are selected (the initial impact velocity V_{ini} is varied from 0.190 to 2.003 m/s).

310 Also, we considered a broad range of Oh number ($0.0013 \leq Oh \leq 0.7869$), a total of 10
311 cases of Oh (i.e., $Oh = 0.0013, 0.0026, 0.0052, 0.0104, 0.0262, 0.0525, 0.1049, 0.2620,$
312 0.5246 and 0.7869) are selected (the droplet viscosity μ_d is varied from 0.0005 to 0.3 N s/m²).
313 Note that such a range of droplet viscosity considered herein corresponds to a range of 0.5 –
314 300 times the viscosity of a water droplet which covers a range of Oh of most practical
315 interest among diverse engineering fields, e.g., gasoline ($Oh \sim 0.0018$), water ($Oh \sim 0.0026$),
316 ethyl alcohol ($Oh \sim 0.0063$), squalene ($Oh \sim 0.05$), and printable ink ($Oh \sim 0.1-1.0$) with the
317 given droplet diameter (2 mm) [64,65]. 5 cases of θ_{eqi} (i.e., $\theta_{eqi} = 20^\circ, 55^\circ, 90^\circ, 125^\circ$ and
318 160°) are also selected to cover a wide range of surface wettability from hydrophilic to

319 (super) hydrophobic surfaces.

320 The impact physics on a spherical **target** is practically identical to that on a flat substrate
321 if a droplet collides onto a sufficiently large **spherical target** which has a diameter 10 times
322 larger than the droplet (i.e., $\Omega \leq 1/10$) [38,47]. Note that this assumption has been
323 experimentally [38] and numerically [40,41] confirmed. Conversely, collision dynamics can
324 differ greatly on a sufficiently small **spherical target** because a **target** can be fully coated by
325 the droplet. In this case the typical maximum spreading state cannot be observed and full
326 coating can usually be observed where $\Omega \geq 1/1.5$ [39,48]. Therefore, we here examine the
327 effect of **target** size between those two thresholds ($1/10 \leq \Omega \leq 1/1.5$) and 4 cases of the
328 droplet-to-**target** size ratio (i.e., $\Omega = 1/10, 1/4, 1/3, 1/2$) are selected.

329 As a result, a total of 2400 collision cases (12 Weber numbers, 10 Ohnesorge numbers, 5
330 surface wettabilities and 4 droplet-to-**target** size ratios) are considered in the current study.
331 This large dataset will also be used to train a deep neural network to formulate a data-driven
332 prediction model as a non-linear regressor as well as to systematically examine the effects of
333 4 collision parameters (We , Oh , θ_{eq} and Ω) on the physical behavior of the spreading time.
334 For more details of the simulation settings readers can also refer to our previous study
335 [40,41].

336

337 **A. Review of typical spreading dynamics**

338 The schematic diagram in Figure 2 summarizes the different spreading regimes for the
339 droplet-solid collision system depending on its physical behavior. If a droplet collides onto a
340 solid surface with sufficiently high inertia, splashing can occur where tiny re-atomized
341 droplets are generated [12,13]. It has been found that higher We and lower Oh normally
342 promote splashing [59] but the detailed threshold for the onset of splashing can also be
343 affected by many other collision environments, e.g., surface roughness [63], target curvature

344 [63] and ambient gas pressure [60], etc. Conversely, if a droplet impacts onto a solid surface
345 with very low inertia compared to surface tension, i.e., at very low Weber number ($We \leq 1$), a
346 droplet can behave like an elastic ball, sometimes bouncing off a solid surface numerous
347 times [66]. This region is the so-called elastic regime.

348 Between those two thresholds (i.e., above the elastic threshold and under the splashing
349 threshold), in general, a droplet spreads over a solid surface and then reaches a maximum
350 spreading state. Afterward, the droplet can retract and is eventually deposited on a surface (or
351 can also bounce off a surface depending on surface hydrophobicity) [12,13]. Such a
352 spreading regime can typically be categorized into the following three sub-regimes depending
353 on the underlying physical mechanisms: (i) capillary-driven spreading at relatively low
354 Weber number [41,61], (ii) inertia-driven but capillary-limited (inviscid) at high We and low
355 Oh [16,40,62] and (iii) inertia-driven but viscous-limited at high We and high Oh [16,40,62].

356 If We is sufficiently small (i.e., We on the order of 10^0), spreading dynamics is primarily
357 driven by capillary effects acting on the three-phase triple contact line [61,67,68]. At such a
358 capillary-driven spreading regime (i.e., regime 1 in Fig.2, also denoted as “CD regime”
359 hereafter), surface wettability and the associated wetting nature can play more significant
360 roles in the spreading process rather than inertial effects due to the low Weber number
361 [18,41]. At a very early stage of contact the apparent contact angle (θ_{ap}) is nearly 180° which
362 is far from its equilibrium contact angle (θ_{eqi}). Therefore, the capillary force which can be
363 scaled as $F_c \sim \sigma R_w (\theta_{ap} - \theta_{eqi})$ (where R_w is the contact radius) starts to act near the contact line
364 and drive the spreading process. Note that this capillary effect can make a droplet spread over
365 a solid surface even if the initial collision velocity is zero and it has also been experimentally
366 [15] and numerically [41] confirmed that the maximum spreading extent of millimetric
367 droplets can be larger than the initial droplet diameter in cases of zero impact velocity. In
368 Fig.2 the CD regime is indicated by the blue region.

369 If We is sufficiently large [i.e., $We \geq O(10^1)$] but Oh is small [i.e., $Oh \leq O(10^{-2})$], the
370 spreading process is primarily governed by inertial effects but is terminated by the capillary
371 limit [16,40,62]. At this inertia-driven but capillary-limited (inviscid) regime (i.e., regime 2 in
372 Fig.2, also denoted as “IC regime” hereafter), the effects of surface wettability and viscous
373 dissipation play minor roles in the spreading process, due to the high We and low Oh
374 numbers. In other words, a droplet can spread over a surface until the interfacial deformation
375 is no longer available. The spreading dynamics in the IC regime can also be characterized by
376 the ejection of a thin liquid sheet (i.e., jetting lamella) at the initial stage [69], and
377 considerable deformation of the leading edge outer rim can be expected during the
378 subsequent spreading stage [27,40,70]. The maximum spreading extent can be scaled by the
379 Weber number, and $We^{1/4}$ [16,40,62] and $We^{1/2}$ [29] rules have been shown to work well. In
380 Fig.2 the IC regime is indicated by the red region.

381 On the other hand, if both We and Oh are sufficiently large [i.e., $We \geq O(10^1)$ and $Oh \geq$
382 $O(10^{-2})$], the spreading process is primarily governed by inertial effects but is terminated by
383 the viscous limit [16,40,62]. At this inertia-driven but viscous-limited spreading regime (i.e.,
384 regime 3 in Fig.2, also denoted as “IV regime” hereafter), the effect of viscous dissipation
385 plays a significant role whereas the effect of surface wettability is still negligible due to the
386 high We and Oh numbers. In other words, a droplet can spread over a surface until all the
387 initial impact (kinetic) energy is dissipated by viscous damping. No lamella ejection is
388 expected in this viscous regime [19] owing to the strong viscous resistance which hinders the
389 interfacial deformation. The maximum spreading extent can be scaled by the Reynolds
390 number and a $Re^{1/5}$ rule has been found to hold well [16,29,40,62]. In Fig.2 the IV regime is
391 indicated by the green region.

392 Note that, in the present study, we focus on these three spreading regimes (i.e., CD, IC,
393 and IV regimes). The other collision regimes (elastic regime and splashing regime) are

394 beyond the current scope due to their different collision phenomena and different area of
395 interest among the scientific communities.

396 Figure 3 depicts the typical spreading processes of droplets for the three different
397 spreading regimes summarized in Fig.2 above. The morphological evolutions of spreading
398 droplets for the CD, IC, and IV regimes are presented in Fig.3(a), (b), and (c), respectively.
399 Note that the (non-dimensional) advective time scale (t_a^*) is used for all cases in Fig.3, since
400 here we simply focus on describing the typical spreading processes rather than addressing
401 their detailed physical interpretations. Note also that the collision cases where $\Omega = 1/10$
402 which can reasonably be regarded as a flat substrate [38,47] are considered in Fig.3 due to its
403 generality. As seen in Fig.3(a), the capillary effect acting near the contact line drives the
404 spreading process for the CD regime. The upper-part of the droplet can nearly maintain its
405 spherical shape owing to the low Weber number (low impact velocity), whereas a capillary
406 wave is formed above the contact line due to the local interfacial curvature [see $t_a^* = 0.34$ in
407 Fig.3(a)] [18,19,68]. As the capillary waves propagate to the central part of the droplet, the
408 droplet shows a staircase-like pyramidal structure [18,19,68] without significant interface
409 deformation ($0.34 \leq t_a^* \leq 0.50$). After the maximum spreading state, the oscillatory motion of
410 the central part is observed in the vertical direction and then the droplet reaches an
411 equilibrium state ($0.50 \leq t_a^* \leq 5.00$).

412 In the IC regime, as seen in Fig.3(b), a thin liquid sheet (lamella) is ejected above the
413 solid surface due to the strong inertial effect overwhelming the surface tension without
414 significant viscous resistance [see $t_a^* = 0.25$ in Fig.3(b)]. The droplet rapidly spreads over the
415 surface and the thin lamella continues consuming liquid mass, forming a thick rim at the
416 leading edge of the lamella ($0.80 \leq t_a^* \leq 1.74$) [27,40,61,70]. **As the lamella spreads on a
417 surface, the leading edge of the lamella is detached from the surface due to its strong radial
418 motion at $t_a^* = 0.80$.** After the maximum spreading state, the droplet starts its retraction

419 motion, driven by the capillary effect, to reduce the increased surface energy during the
420 spreading stage ($1.74 \leq t_a^* \leq 8.00$).

421 In the IV regime, as explained above, no notable lamella ejection is observed [19]. The
422 interfacial deformation is significantly hindered by the strong viscous damping. After the
423 maximum spreading state ($t_a^* = 0.98$), the central part of the droplet slightly and slowly rises
424 again, reaching the equilibrium state, but no meaningful recoiling motion is seen in terms of
425 its wetted area. Note that all these results showing the evolutions of the droplet morphologies
426 for the three spreading regimes depicted in Fig.3 are consistent with the experimental
427 observations of Lin *et al.* [19].

428 Figure 4 quantifies the temporal evolution of the spreading extent characterized by the
429 non-dimensional spreading diameter β^* ($\beta^* = \beta/D_d$ where β denotes the arc length of the
430 droplet in contact with a solid **target**, see the inset in Fig.4) of the droplets for the three
431 collision cases shown in Fig.3. Note that, for the CD and the IC regimes, detecting the
432 spreading time with the maximum value of β^* can be straightforward because the droplet
433 starts its recoiling motion soon after the maximum spreading state (see black solid line and
434 blue dashed line). The corresponding time for the maximum value of β^* is visually evident.
435 Conversely, as shown in Fig.3(c), usually no obvious recoiling stage occurs for the IV regime
436 due to the high viscous resistance (see red dotted line in Fig.4). In this case, detecting the
437 spreading time can be quite confusing and even a very small disturbance in β^* can lead to a
438 critical bias for measuring the spreading time. Therefore, to avoid such a complexity, the time
439 when a droplet attains 95 % of the maximum value of β^* is defined as the spreading time as
440 marked by three triangles in Fig.4 (see also green arrow). Note that, hereafter, the non-
441 dimensional forms of the spreading time are denoted by τ_a (if the advective time scale is
442 used), τ_c (if the capillary time scale is used) and τ_v (if the viscous time scale is used).

443

444 **B. Validations**

445 In the above subsection, we addressed the typical spreading dynamics for the three
446 different spreading regimes. Before we start analyzing the spreading time for each regime in
447 more detail in this subsection, we validate our simulation method by comparisons with
448 existing experimental results to check its capability for simulating the droplet collision
449 phenomena in terms of the spreading time.

450 Figure 5 compares our simulation results with the experimental data of Mitra *et al.* [2]
451 and Banitabaei and Amirfazli [71] focusing on the collision behavior on spherical targets for
452 a broad range of We (We = 0.9 – 155.5). First, Fig.5(a) depicts the temporal variation of β^*
453 and the morphological evolution of the droplet for the low Weber number case (We = 0.9).
454 The insets in Fig.5(a) shows the interfacial shapes captured by the experiment (on the right
455 side of each panel) [2] and by the current simulation (on the left side of each panel) at three
456 different time instants ($t_a^* = 0.24, 0.47, \text{ and } 0.71$). As seen, our simulation result shows a
457 good agreement with the experimental observation. Fig.5(b) depicts the temporal variation of
458 the non-dimensional film thickness, h^* [$h^* = h/D_d$, where h is the liquid film thickness
459 measured on the impact center, see the inset on the lower-left corner in Fig.5(b)] and the
460 morphological evolutions of the droplet for the high Weber number case (We = 155.5). The
461 other insets in Fig.5(b) show the interfacial shapes captured by the experiment [71] and by
462 the current simulation at four different time instants ($t_a^* = 0.30, 0.60, 1.05 \text{ and } 1.95$). Again,
463 our simulation result shows a good agreement with the experimental observation. Since the
464 liquid film thickness is one of the most sensitive parameters to the mesh size in this type of
465 numerical simulations [72], we here present our grid-convergence test result as well in
466 Fig.5(b). As seen in the dashed gray box, our simulated h^* almost converges at 32 CPR (cells
467 per radius) as in our previous study [48,51]. Overall, our simulation result shows an excellent
468 agreement with the experimental result [71] in terms of both qualitative and quantitative

469 comparisons. Note also that all simulation cases presented in the current study have been
470 performed with 64 CPR to be on the safe side.

471 In figure 6 we further validate our simulation results focusing on the quantification of the
472 spreading time for a broad range of We and Oh. Fig.6(a) depicts the non-dimensional
473 spreading time, τ_a , for the low Oh case (Oh = 0.0026 for a water drop). As seen, our
474 simulation result is consistent with the experimentally measured data of Huang and Chen
475 [18]. A small deviation can be attributed to the difference of the definition of spreading time
476 (i.e., we defined the spreading time as the time that a droplet attains 95% of the maximum
477 spreading extent. See Fig.4 and its explanation). In Fig.6(b), we compare the simulated τ_a
478 with the existing numerical result [21] and empirical correlation [19] for the high Oh case
479 (Oh = 0.62 in the case where the droplet is 240 times the viscosity of a water drop). Our
480 result again shows an excellent agreement with the existing results [19,21].

481 Further validation test results for drop impact problems are also extensively provided in
482 Appendix B, which were presented in our previous study. From Figs.5 – 6 and Appendix B,
483 we can conclude that our simulation methods used in this study can reasonably reproduce the
484 experimental data for a wide range of collision parameters in the context of both qualitative
485 and quantitative comparisons. Note that the current numerical framework has been widely
486 applied to various droplet-solid collision problems and has been thoroughly validated against
487 many existing experimental data [40,41,48,49,51]. Note also that the impact parameters
488 considered in our current and previous validation tests reasonably cover the practical
489 conditions used in the current study ($0 \leq We \leq 110$, $20^\circ \leq \theta_{eqi} \leq 160^\circ$, $0.0013 \leq Oh \leq 0.7869$
490 and $1/10 \leq \Omega \leq 1/2$).

491

492 **C. Effects of collision parameters on the spreading time and prediction**
493 **models**

494 We now analyze the spreading time for the three different spreading regimes addressed
495 above, i.e., the capillary-driven (CD), the inertia-driven but capillary-limited (IC) and the
496 inertia-driven but viscous-limited (IV) regimes in more detail and present a prediction model
497 for each regime. Note that the collision cases for a flat surface ($\Omega = 1/10$) are chosen first for
498 demonstrating our results owing to its generality. Then we further extend our analysis
499 considering the effect of the droplet-to-target size ratio (Ω).

500 We first examine the spreading time for the CD regime ($We \leq 10^1$). At this low We
501 region the spreading dynamics is primarily governed by its wetting nature and the associated
502 capillary effect where surface wettability plays a dominant role. Therefore, the capillary time
503 scale, t_c^* , can be an appropriate time scale for analyzing physical characteristics of the
504 spreading time in this regime [16,19].

505 Figure 7(a), (b), and (c) depict the effects of We , θ_{eqi} and Oh on the spreading profiles,
506 respectively. As seen in Fig.7(a), the spreading time (τ_c) basically decreases as We increases
507 whereas the maximum value of β^* (denoted as β_{max}^* hereafter) increases with We . This result
508 is natural because higher spreading velocity can be attained at the high We cases [19], thus a
509 shorter spreading time is needed. Note that this is also consistent with the experimental
510 observation of Lin *et al.* [19]. For the cases shown in Fig.7(a), τ_c , where $We = 10$ ($\tau_c = 0.69$),
511 is about 21 % shorter compared to that for the case where $We = 2$ ($\tau_c = 0.87$). In Fig.7(b), as
512 expected, the dominant effect of the surface wettability is seen. The droplet can spread for a
513 longer time over the surface as θ_{eqi} decreases (i.e., as the surface becomes hydrophilic) and
514 can attain a larger spreading extent as well owing to the strongly wettable nature of the
515 hydrophilic surface. This result is also consistent with the existing experimental data [19,43].
516 For the cases shown in Fig.7(b), τ_c , where $\theta_{eqi} = 20^\circ$ ($\tau_c = 1.38$), is almost double the case
517 where $\theta_{eqi} = 160^\circ$ ($\tau_c = 0.70$).

518 It is interesting to observe the effect of Oh on τ_c . As depicted in Fig.7(c), for relatively

519 low Oh cases ($Oh \leq 0.0104$, see 4 dashed lines), no meaningful effect of Oh is seen because
 520 the spreading process is primarily driven by the capillary effect and is also terminated by the
 521 capillary limit due to the negligible role of the viscous resistance at such low Oh numbers
 522 [16,40]. Note also that this trend is consistent with the existing experimental observation of
 523 Lee *et al.* [15] who showed that there was no notable difference of the spreading time
 524 between a water droplet ($Oh \sim 0.0026$) and a glycerol mixture ($Oh \sim 0.0263$, about 10 times
 525 the viscosity of a water droplet) under the given surface conditions ($\theta_{eqi} = 52^\circ - 61^\circ$) in the
 526 low We region ($We < 2.5$). Conversely, the spreading profiles start to be affected by Oh as Oh
 527 further increases above $Oh \geq 0.0262$ [see solid lines in Fig.7(c)] and the recoiling stages
 528 disappear for much higher Oh cases ($Oh \geq 0.1049$) owing to the significant viscous damping.
 529 The interesting point is that no notable difference of τ_c can be found even for higher Oh cases
 530 (τ_c is almost constant for the entire range of Oh). We found that this is because Oh primarily
 531 affects the spreading velocity, not the spreading time, for the higher Oh region. To check the
 532 effect of Oh on the spreading velocity in more detail, in Fig.7(d) and (e), we plot the non-
 533 dimensional spreading velocity V_s^* (defined as $V_s^* = \beta_{max}^* / \tau_c$) and the non-dimensional
 534 maximum spreading extent β_{max}^* for the cases shown in Fig.7(c), respectively. As seen, V_s^* is
 535 nearly constant for the low Oh region (roughly where $Oh \leq 0.01$), whereas it significantly
 536 decreases as Oh increases in the high Oh region ($Oh \geq 0.05$). Since the effect of Oh on β_{max}^*
 537 (this result in Fig.7(e) is also consistent with the existing experimental result [16]) evidently
 538 shows exactly the same trend as the effect of Oh on V_s^* , τ_c can be considered almost
 539 independent of Oh for the entire range of Oh.

540 It is also interesting to compare the spreading mechanism for this CD regime with the
 541 retraction mechanism for high-speed impact cases. In fact, both processes (i.e., spreading of
 542 low-speed impact and retraction of high-speed impact) are presumably based on the same
 543 physical mechanism, i.e., capillary-driven motion under viscous resistance. In the high-speed

544 impact case, at the beginning of the retraction stage (i.e., right after the maximum spreading
545 state), the droplet shows a severely deformed interfacial shape and its motion is almost
546 stopped. Then we begin to see the capillary-driven recoiling motion under viscous resistance,
547 to reduce its increased surface energy caused by surface deformation during the former
548 spreading stage [62]. Since the spreading stage in the CD regime is also driven by the
549 capillary effect under viscous resistance, one can expect that there should essentially and
550 naturally be consistent physical characteristics between the spreading behavior of the low-
551 speed impact in the CD regime and the retraction behavior of the high-speed impact cases.
552 We found that the spreading velocity in the CD regime analyzed in Fig.7(d) demonstrates the
553 same trend as the retraction characteristics which were experimentally observed by Bartolo *et*
554 *al.* [62]. They showed that the retraction rate of a droplet is not affected by Oh where $Oh <$
555 0.05 , whereas it is strongly reduced as Oh increases where $Oh > 0.05$. Their result shows an
556 excellent agreement with our findings described in Figs.7(c) and (d).

557 In Fig.7 above, we observed the obvious effects of We and θ_{eqi} on τ_c , whereas the effect
558 of Oh can be negligible. Therefore, τ_c in the CD regime can be properly represented by We
559 and θ_{eqi} and **we propose a prediction model:**

560

$$561 \quad \tau_c = (1.66 - 0.31\theta_{eqi}) We^{-0.24} \quad (6)$$

562

563 In figure 8, a total of 600 τ_c data are plotted with the **prediction model** in Eq.(6). As seen,
564 all τ_c data for a range of $We \leq 10$ can nearly be collapsed on the single curve (see black solid
565 line), regardless of Oh. Above the threshold of $We \sim 10$ (i.e., beyond the CD regime) τ_c
566 cannot be properly represented by Eq.(6).

567 We now examine the spreading time for the IC regime [$We \geq O(10^1)$ and $Oh \leq O(10^{-2})$].
568 At this high We but low Oh region, the spreading dynamics is primarily dominated by the

569 strong inertial effect but terminated by the capillary limit [16,40,62]. Therefore, the capillary
570 time scale, t_c^* , can still be applicable for this regime [16,19].

571 Figure 9(a) depicts the effect of We on the spreading profiles. τ_c first decreases as We
572 increases for the CD regime (see 5 dashed lines) as shown in Fig.7(a) above, but it increases
573 again if We further increases ($We \geq 30$) as the collision dynamics sufficiently enter the IC
574 regime (see 6 solid lines). This somewhat interesting result can be attributed to the severe
575 deformation of the leading edge of the lamella and its relaxation phenomena [40]. As the
576 droplet spreads over the surface the leading edge of the lamella continues deforming. The tip
577 of the lamella is detached from the surface, and it also continues growing as a rim by
578 consuming the liquid mass from the central region of the droplet [27,70]. Note that such a
579 deformed rim plays an important role in determining the maximum wetting area [70] since
580 the volume of the rim is not small [27] and the maximum wetting area is attained during the
581 relaxation period of the rim [40]. Therefore, a more deformed shape at higher We leads to the
582 longer relaxation time of the rim and eventually the longer spreading time. Fan *et al.* [73] also
583 experimentally showed that the spreading time of a water droplet on a spherical target first
584 decreases as We increases, then it again increases with We if $We \geq 36.5$, showing an
585 agreement with our results in Fig.9(a). However, the authors think it would be ideal if our
586 physical interpretation of the underlying mechanism shown in Fig.(b) could also be
587 confirmed by additional experiments.

588 Fig.9(b) illustrates such a deformation and relaxation process of the rim in more detail.
589 For the relatively low We case ($We = 20$), as depicted on the left side of Fig.9(b), the droplet
590 continues its spreading process until it reaches the maximum spreading state and no
591 significant deformation of the lamella edge is observed because the spreading dynamics does
592 not fully enter the IC regime yet. Conversely, for the high We case where $We = 110$ [see the
593 right side of Fig.9(b)], a severe deformation of the rim detached from the surface is clearly

594 captured at $t_c^* = 0.60\tau_c$ (see black line). Then the rim begins its relaxation to reduce the
 595 surface energy and eventually reaches the maximum spreading state (see pink line). Note that
 596 the relaxation period takes quite a long time (about 40 % of τ_c in this case) eventually leading
 597 to the delayed τ_c .

598 Fig.9(c) and (d) show the effects of θ_{eqi} and Oh on the spreading profiles, respectively.
 599 As seen, no meaningful effect of θ_{eqi} on τ_c is observed, because the effect of the surface
 600 wettability is overwhelmed by the strong inertial effect during most of the spreading stage
 601 [27,41]. Although the surface wettability plays a non-negligible role in the spreading profiles
 602 near the maximum spreading state and after τ_c (i.e., during the retraction stage) where the
 603 spreading velocity and the associated inertial effect are at a minimal level, the effect of θ_{eqi} on
 604 τ_c is almost negligible until the droplet reaches nearly the maximum spreading state. For this
 605 case τ_c can be considered to be a constant value regardless of the surface wettability. The
 606 effect of Oh on τ_c is also not meaningful if Oh is sufficiently small. As seen in Fig.9(d), τ_c is
 607 not affected by Oh if $Oh < 10^{-2}$ (see 3 solid lines, i.e., at IC regime), whereas τ_c decreases as
 608 Oh increases if $Oh > 10^{-2}$ (see 3 dashed lines, i.e., beyond the IC regime). This result is also
 609 consistent with the existing result of Zhu *et al.* [46] who demonstrated that the spreading
 610 dynamics is independent of Oh for low Oh, whereas it is strongly affected by Oh for high Oh
 611 under similar We conditions as the present study ($We \geq 25$). Their proposed threshold ($Oh \sim$
 612 0.008) is also consistent with the current observation ($Oh \sim 0.01$).

613 In Fig.9 above, we observed the obvious effect of We on τ_c , whereas the effects of θ_{eqi}
 614 and Oh can be negligible. Therefore, τ_c in the IC regime can be properly represented by We
 615 and **we propose a prediction model:**

616

$$617 \quad \tau_c = 0.31We^{0.18} \quad (7)$$

618

619 In figure 10, a total of 150 τ_c data are plotted with the **prediction model** in Eq.(7). As
620 seen, all τ_c data for a range of $We > 30$ can nearly be collapsed on the single curve (see black
621 solid line) regardless of θ_{eqi} and Oh.

622 We now examine the spreading time for the IV regime [$We \geq O(10^1)$ and $Oh \geq O(10^{-2})$].
623 At this high We and Oh region the spreading dynamics is primarily dominated by the strong
624 inertial effect but terminated by the viscous limit [16,40,62]. Therefore, the viscous time
625 scale, t_v^* , would be the proper scale in this regime.

626 Figure 11(a), (b), and (c) depict the effects of We, θ_{eqi} , and Oh on the spreading profiles,
627 respectively. As seen in Fig.11(a), the spreading time, τ_v , naturally decreases as We increases
628 due to the higher spreading velocity [19]. The delayed spreading time shown in the IC regime
629 (caused by the relaxation period of the deformed lamella rim) is not observed in this IV
630 regime because the interfacial deformation is significantly restricted by the strong viscous
631 damping. No meaningful effect of θ_{eqi} on τ_v is seen [see Fig.11(b)] due to its minor role
632 compared to the strong inertial and viscous effects. Conversely, as expected, the dominant
633 effect of Oh on τ_v is obviously detected in Fig.11(c). Since the spreading stage is terminated
634 much earlier for higher Oh cases under the given We due to increased viscous resistance (i.e.,
635 the initial kinetic energy is more quickly dissipated [44]), τ_v is significantly reduced for
636 higher Oh conditions. Note that the effect of Oh on τ_v observed in Fig.11(c) is also consistent
637 with the existing experimental observation [19,44].

638 In Fig.11 above, we observed the obvious effects of We and Oh on τ_v , whereas the effect
639 of θ_{eqi} can still be negligible. Therefore, τ_v in the IV regime can be properly represented by
640 We and Oh, and **we propose a prediction model**:

641

$$642 \tau_v = 0.3Oh^{-1.13}We^{-0.22} \quad (8)$$

643

644 In figure 12, a total of 250 τ_v data are plotted with the **prediction model** in Eq.(8). As
645 seen, all τ_v data for a range of $We > 30$ can nearly be collapsed on the single curve (see black
646 solid line), regardless of θ_{eqi} .

647

648 **D. Effect of **target** curvature and data-driven prediction model**

649 In the above subsection, we addressed the effects of the three collision parameters (We ,
650 θ_{eqi} , and Oh) on the spreading time for the three different spreading regimes (CD, IC, and IV
651 regimes) and a **prediction model** was presented for each regime for the cases where $\Omega = 1/10$
652 (for a flat surface). We now expand our analysis considering the surface curvature which has
653 been relatively less explored. The effect of the **target** curvature characterized by the droplet-
654 to-**target** size ratio (Ω) on the spreading time and a new prediction model for the spreading
655 time which can be applied to the entire range of the 4 collision parameters will be presented
656 in this subsection.

657 Figure 13(a), (b), and (c) depict the effects of Ω on the spreading profiles for the CD, the
658 IC, and the IV regime, respectively. In the CD regime, as seen in Fig.13(a), τ_c increases with
659 Ω (on the smaller **target**). This behavior is attributed to the reduced capillary force acting near
660 the contact line for the smaller **targets**. As described above, at this low We level, the dominant
661 force driving the spreading process is the capillary force F_c [which can be scaled as $F_c \sim$
662 $\sigma R_w(\theta_{ap} - \theta_{eqi})$] formed at the initial stage of contact. At the instant that a droplet contacts a
663 solid surface, its apparent contact angle (θ_{ap}) is almost close to 180° and the difference
664 between the apparent contact angle and the equilibrium contact angle results in the capillary
665 force F_c on the contact line. Since such an apparent contact angle at the initial contact stage is
666 formed at a lower level on the smaller **target** [see θ_{ap2} and yellow dashed line in the zoom-in
667 of the dashed black box, in Fig.13(d)] compared to the bigger **target** (see θ_{ap1}), $\theta_{ap} - \theta_{eqi}$ and
668 F_c are reduced on the smaller **targets** eventually leading to lower spreading velocity and

669 longer spreading time. For the cases shown in Fig.13(a), τ_c where $\Omega = 1/2$ ($\tau_c = 0.99$) is about
670 21 % longer compared to that for the flat surface ($\Omega = 1/10$, $\tau_c = 0.82$).

671 As seen in Fig.13(b), the effect of Ω on τ_c is apparently observed also in the IC regime,
672 i.e., τ_c increases on the smaller **targets**. Note that the trend of the effect of Ω on τ_c in the IC
673 regime is also consistent with many existing studies [40,46,47] and this behavior is attributed
674 to the reduced lamella deformation on the smaller **target** [40]. As already shown in Fig.3(b)
675 above, a thin lamella sheet is ejected at the early stage of spreading. Although the motion of
676 bulk liquid is primarily dominated by inviscid inertial effects [16,40,62], the dynamics in the
677 vicinity of the lamella jet is locally but considerably affected by viscous effects [20,69]. Since
678 the smaller **targets** provide more space in the vertical direction as in an expanding channel
679 flow for the lamella sheet being ejected [i.e., higher ζ is provided on smaller **targets** as
680 illustrated in the zoom-in of the dashed red box in Fig.13(d)] the local velocity of the radially
681 expanding lamella jet is reduced on the smaller **targets**. Therefore, the lamella jet develops
682 less sharply, having a more rounded and bluff shape [40,69]. This allows the droplet to
683 expend its kinetic energy more efficiently for the spreading itself rather than the rim
684 deformation detached from the surface eventually leading to the longer spreading time as
685 well as the longer spreading extent compared to those on the larger **targets** (under the same
686 kinetic energy level). A more detailed physical mechanism for the reduced lamella
687 deformation on smaller **targets** can be found in the existing numerical study [40]. For the
688 cases shown in Fig.13(b), τ_c , where $\Omega = 1/2$ ($\tau_c = 0.80$) is about 45 % longer compared to that
689 for the flat surface ($\Omega = 1/10$, $\tau_c = 0.55$).

690 The effect of Ω on τ_v in the IV regime is depicted in Fig.13(c). As seen, no meaningful
691 effect of Ω on τ_v is observed for this highly viscous regime. Although the maximum
692 spreading extent is slightly increased on the smaller **targets** due to the reduced energy
693 dissipation [40], τ_v is almost constant in terms of its practical measurement.

694 In figure 14, we refine the regime map (shown in Fig.2 above) to summarize the major
695 findings presented so far, including the details of the physical behaviors of the spreading
696 time. The underlying physical mechanism governing the spreading process, the characteristic
697 time scale, the effects of various collision parameters on the spreading time and the
698 **prediction model** (for a flat surface) are also presented for each regime. Note that the effects
699 of 4 collision parameters, i.e., We , θ_{eqi} , Oh , and Ω are indicated by 4 different markers, i.e.,
700 circles, squares, triangles and inverted triangles, respectively. If a certain parameter has a
701 positive (or negative) effect on the spreading time, its effect is indicated by a red (or blue)
702 marker, whereas an empty marker is used if its effect is negligible. For example, if the
703 spreading time increases with We , the effect of We is marked by a filled red circle.
704 Conversely, the effect of We is marked by a filled blue circle if the spreading time decreases
705 as We increases.

706 Fig.14 highlights the complex nature of the spreading time owing to the different
707 physical mechanism for each regime. The effects of the impact parameters (We , θ_{eqi} , Oh and
708 Ω) also play different roles. For instance, the effect of We is positively related to the
709 spreading time for the CD and the IV regime, whereas it is negatively related to the spreading
710 time for the IC regime. In particular, none of the 4 impact parameters shows a consistent
711 effect across the three spreading regimes, which highlights the significant nonlinearity and
712 complexity of the spreading time. The characteristic time scales are also different.

713 Therefore, formulating a single correlation for the spreading time as a broad cross-over
714 across the three different spreading regimes is quite difficult since it should cover all the
715 different physical mechanisms, different characteristic time scales, and different effects of the
716 collision parameters. In addition, if the effect of the **target** curvature (i.e., the effect of Ω) on
717 the spreading time were to be considered together, the problem would be much more
718 complicated. Rather than formulating a single correlation, therefore, we consider here a data-

719 driven approach to model such complex and non-linear behavior of the spreading time across
720 all the spreading regimes.

721 Recently, data-driven approaches have been broadly applied to model complicated fluid
722 phenomena such as turbulent flows [74] or multiphase flow systems [75]. For droplet impact
723 problems as well, the maximum spreading extent [41], the splashing threshold [76], and the
724 lamella area [77] have successfully been modeled using data-driven techniques, based on
725 their powerful ability to correlate non-linear relations among large datasets and reduce the
726 level of complexity.

727 The spreading time can basically be modeled as a function of 4 impact parameters as
728 follows:

729

$$730 \quad \tau_a = f(\text{We}, \text{Oh}, \theta_{\text{eqi}}, \Omega) \quad (9)$$

731

732 Here, the advective time scale is applied for our data-driven prediction model, in order to
733 treat our dataset in a consistent manner. f is the non-linear and complicated (unknown)
734 function between inputs (4 impact parameters) and output (τ_a) data which is now modeled by
735 using an artificial intelligence technique. The multilayer neural network, also called the
736 multilayer perceptron (MLP) [78], is utilized as a nonlinear regressor to model the function f .
737 Figure 15 illustrates the schematic diagram of the typical MLP structure. The MLP consists of
738 the input layer (marked by red nodes), the multiple hidden layers (marked by green nodes)
739 and the output layer (marked by blue nodes). Each layer usually consists of a number of
740 nodes (also called neurons), whereas the last output layer has only one output node. The 4
741 input features (We , Oh , θ_{eqi} and Ω) are provided from the input layer to the first hidden layer
742 and then linear combinations of the input features are calculated on each node. Such linear
743 combinations are nonlinearized by using an activation function and then are delivered to the

744 next hidden layer. Such an operation is called a feedforward procedure and is applied to all
745 the subsequent layers. The last output layer has no nonlinearization operation and provides
746 the final output value.

747 A total of 2400 impact cases (12 Weber numbers, 10 Ohnesorge numbers, 5 surface
748 wettabilities, and 4 droplet-to-**target** size ratios) are utilized to set the dataset. The dataset is
749 randomly split into 3 parts: the training dataset (70 %, 1680 data points), the validation
750 dataset (15 %, 360 data points) and the test dataset (15 %, 360 data points). The training
751 dataset and the test dataset are used to train our MLP model and finally test the accuracy of
752 the trained model, respectively, whereas the validation dataset is utilized to check if
753 overfitting occurs during the training process. More detailed information on our MLP
754 modeling and its training procedure can be found in the **Appendix C**.

755 Figure 16 compares the spreading time predicted by our data-driven model with the true
756 values (measured values from our numerical simulations). 360 data points from the test
757 dataset are plotted for comparison (see black circles). As can be seen, τ_a predicted by our
758 MLP model demonstrates a good agreement with the true dataset, showing that the mean
759 square error (MSE) is on the order of $O(10^{-4})$ ($MSE = 5.23 \times 10^{-4}$). Note that, except for only
760 a few cases, all the tested cases fall inside a deviation range of $\pm 15\%$. Note that all the test
761 data samples shown in Fig.16 have not been used for the training process of our MLP model.
762 To the best of our knowledge, our prediction model is the first data-driven model for the
763 spreading time of the **droplet collision system which can be applied for a curved spherical**
764 **target**. As one of the most universal models, the current model can cover a very broad range
765 of impact parameters across not only the three spreading regimes (i.e., CD, IC, IV regimes
766 and the transition area among them as well), but also a wide range of droplet-to-**target** size
767 ratio (from a flat substrate to a small **spherical target**).

768 **Last but not least, the authors would like to clarify that there is no physics in the current**

769 data-driven prediction model, which can indeed be understood as providing both pros and
770 cons [75]. The fact that such a data-driven approach is not fully based on a physical
771 background might be a reason for which it has been less attractive to some scientific
772 communities compared to conventional approaches supported by strong physical principles
773 [79,80]. However, both approaches share identical goals, i.e., modeling complicated
774 phenomena which usually need to consider many multi-dimensional variables [80,81]. The
775 advantages of data-driven approaches can also be that users would not need to consider
776 physics models alone for those complex system because data-driven models basically don't
777 rely on constraints attributed to physics [75]. This characteristic allows them to work as a
778 useful non-linear regressor. For example, the current data-driven prediction model can predict
779 the spreading time even for transition regimes marked in the gray zones in Fig.14, which is
780 apparently very difficult to predict using Eqs. 6 – 8 above due to their non-linear nature. This
781 clearly shows the advantage of data-driven approaches at least in the context of engineering
782 applications and practical usefulness.

783

784 **IV. CONCLUSION**

785

786 In this study we quantitatively and systematically investigate the spreading time of droplet
787 impact onto a spherical target for a wide range of impact parameters ($0 \leq We \leq 110$, $20^\circ \leq$
788 $\theta_{eqi} \leq 160^\circ$, $0.0013 \leq Oh \leq 0.7869$ and $1/10 \leq \Omega \leq 1/2$). A total of 2400 collision cases are
789 simulated and are used to better understand the physical behavior of the spreading time. This
790 large simulation dataset is also used for training a data-driven prediction model as well as for
791 the derivation of prediction models.

792 For the three different spreading regimes, i.e., the capillary-driven (CD), the inertia-

793 driven but capillary-limited (IC) and the inertia-driven but viscous-limited (IV) regimes, the
794 effects of the impact parameters on the spreading time are analyzed in detail. Not only the
795 spreading time, but also proper time scales, dominant impact parameters and associated
796 physical behaviors all significantly and non-linearly change across the three spreading
797 regimes. The spreading time is primarily dominated by the We and θ_{eqi} in the CD regime
798 whereas only the Weber number plays a meaningful role in the IC regime. In the IV regime it
799 is mainly affected by We and Oh . An improved prediction model is also presented for each
800 regime. To the best of our knowledge, these proposed prediction models (Eqs. 6-8) are the
801 first models for the spreading time using only the controllable independent variables (e.g.,
802 We , Oh and θ_{eqi}) in explicit form which can provide more straightforward solutions compared
803 to the existing models [18-21,44,45].

804 The effect of target curvature on the spreading time is also analyzed in detail for the
805 three different spreading regimes, showing complicated trends due to its non-linear nature.
806 Finally, a data-driven prediction model is proposed to predict the spreading time broadly
807 across the three different spreading regimes. Our proposed data-driven model shows good
808 prediction capability for the full ranges of impact parameters considered in the present study
809 and also shows a good agreement with our simulation dataset.

810 The limitation of the current study is as follows: (i) As in the existing models, prediction
811 models (Eqs. 6-8) are still not derived from physical background. Further efforts are needed
812 to combine “physics-based models” which could be helpful in understanding such
813 complicated behavior within the existing empirical results. (ii) Users need to be careful in
814 using the proposed data-driven model. Although our numerical methods have thoroughly
815 been validated, the number of current data samples (2400 cases) is still insufficient to model
816 the spreading time for the entire spreading regime compared to other common data-driven
817 models trained using “big-data”. Our numerical data used for the training are also confined to

818 the current impact conditions ($0 \leq We \leq 110$, $20^\circ \leq \theta_{\text{eqi}} \leq 160^\circ$, $0.0013 \leq Oh \leq 0.7869$ and
819 $1/10 \leq \Omega \leq 1/2$). Broader ranges of impact parameters, and more precise and abundant data
820 samples can improve the applicability and accuracy of the current model. (iii) While the
821 current axi-symmetric simulation code is sufficiently validated, full three-dimensional
822 simulations are still essential to better understand those complicated physical phenomena. In
823 particular, impact dynamics at high We and low Oh where the liquid sheets and rim can be
824 highly deformed still need to be captured more precisely. We are currently working on these
825 issues.

826

827 **APPENDIX A: Details of the dynamic contact angle**

828 **modeling**

829

830 The dynamic contact angle θ_{dyn} is basically modeled as a function of the equilibrium
831 contact angle and the velocity of the contact line, as in well-known previous numerical
832 studies on drop impact problems [82,83]. In the present study, we use the model of Yokoi *et*
833 *al.* [83]:

834

$$835 \quad \theta_{\text{dyn}}(U_{\text{CL}}) = \begin{cases} \min \left[\theta_{\text{eqi}} + \left(\frac{\text{Ca}}{q_1} \right)^{1/3}, \theta_{\text{mda}} \right], & \text{if } U_{\text{CL}} \geq 0 \text{ (for spreading)} \\ \max \left[\theta_{\text{eqi}} + \left(\frac{\text{Ca}}{q_2} \right)^{1/3}, \theta_{\text{mdr}} \right], & \text{if } U_{\text{CL}} \leq 0 \text{ (for receding)}. \end{cases} \quad (A1)$$

836

837 Ca is the capillary number ($\text{Ca} = \mu U_{\text{CL}} / \sigma$), where U_{CL} is the velocity of the contact line.

838 Contact angle hysteresis is represented by the difference between the allowable minimum

839 (θ_{mdr}) and maximum (θ_{mda}) contact angles, and the dynamic contact angle (θ_{dyn}) can always
840 be maintained between θ_{mdr} and θ_{mda} . q_1 and q_2 are experimentally measured constants, and
841 the same values as in Yokoi *et al.* [83] ($q_1 = 9.0 \times 10^{-9}$ and $q_2 = 9.0 \times 10^{-8}$) are applied to the
842 current study.

843

844 **APPENDIX B: Validation tests in our previous studies**

845

846 We here introduce our validation test results presented in our previous studies. Figure
847 17(a) compares the maximum spreading extent β^*_{max} computed by our simulation with Mao
848 *et al.* [25]’s semi-empirical model which is well-known as one of the most accurate models
849 that can be applied for a wide range of drop viscosity [84]. β^*_{max} shows generally good
850 agreement with the existing model of Mao *et al.* [25] for a wide range of We number ($30 \leq$
851 $We \leq 90$) and Oh number ($0.0013 \leq Oh \leq 0.7869$). Note that a discrepancy shown at the very
852 low Oh region ($Oh \leq 0.0026$) is attributed to the limitation of shape approximation in Mao *et*
853 *al.*’s model [25], which overestimates β^*_{max} if the viscosity is very low and the surface tension
854 force plays an important role [25,84]. Therefore, in figure 17(b), we compare our simulated
855 β^*_{max} again with the well-known viscous-free scaling law proposed by Clanet *et al.* [16]. As
856 can be seen, for cases with very low Oh number ($Oh = 0.0026$), our simulation result shows
857 an excellent agreement with the existing scaling law [16].

858 In figure 17(c), we further compare our simulated time-dependent non-dimensional
859 spreading diameter β^* with three existing experimental results to see if our numerical
860 framework can simulate the impact phenomena well on spherical targets. The solid lines in
861 Fig. 17(c) depict the simulation results, whereas the red squares, blue circles, and black
862 crosses show the experimental results from Mitra *et al.* [2], Liu *et al.* [38], and Khurana *et al.*
863 [39], respectively. Our simulation results also show good agreement with the existing

864 experimental observations in the literature, in the context of dynamic spreading
 865 characteristics. Overall, we concluded that our simulation framework can simulate the drop
 866 impact phenomena on both flat and spherical targets well for a wide range of collision
 867 parameters. Note again that these validation tests shown in figure 17 are not new tests, but
 868 were already presented in our previous study [40].

869

870 **APPENDIX C: Details of MLP training**

871

872 The output value on the n^{th} node in the m^{th} layer is:

873

$$874 \quad a_n^m = g \left(\sum_{k=1}^{N^{m-1}} w_{nl}^m a_l^{m-1} + b_n^m \right) \quad (\text{C1})$$

875

876 Here, N stands for the number of nodes of each layer. The function g denotes the nonlinear
 877 activation function. w_{nl} is the weight between the l^{th} node of the previous layer and the n^{th}
 878 node of the current layer whereas b is the bias.

879 The weights (w) and biases (b) are adjustable coefficients and are first initialized as
 880 random variables and zero, respectively. w and b are continuously updated during the training
 881 process, minimizing the loss function based on the back-propagation algorithm [85]. In the
 882 current study, the mean square error (MSE) is used as a loss function:

883

$$884 \quad \text{MSE} = \frac{1}{M} \sum_{k=1}^M (\tau_a - \tau_{a(p)})^2 \quad (\text{C2})$$

885

886 where M is the number of data points considered and $\tau_{a(p)}$ is the predicted value of the

887 spreading time using the trained MLP model.

888 Three hyperparameters (i.e., the activation function, the number of nodes at each hidden
889 layer and the number of hidden layers) can usually be selected by the user. We chose the
890 common rectified linear unit function (generally called ReLU) as our activation function
891 owing to its well-known capability for deep neural networks [86]. The ReLU can be written
892 as:

$$894 \quad g(x) = \text{ReLU}(x) = \max(x, 0) \quad (\text{C3})$$

895
896 Two other hyperparameters are set by manual search because the current dataset is relatively
897 very simple. Finally, a 4×30 (4 hidden layers and 30 neurons in each hidden layer)
898 architecture is chosen as the MLP network. The ADAM optimization procedure [87,88] is
899 used for the MLP training and the computational implementation for the MLP training is
900 performed using the open-source libraries Keras [89] and TensorFlow [90]. For more details
901 on the deep learning architecture and its training techniques readers can refer to Nielsen [91].

902 The training of our MLP model is carried out until the validation MSE (measured from
903 the validation dataset) reaches a steady state. After 3300 training epochs the validation MSE
904 converges to a steady state of about 4.62×10^{-4} (see Fig.18).

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912 **DECLARATIONS**

913

914 **Ethical Approval**

915 Not applicable.

916

917 **Conflict of interest**

918 The authors declare that they have no conflict of interest.

919

920 **Authors' contributions**

921 Ikroh Yoon led conceptualization, data curation, formal analysis, investigation, validation,
922 visualization, and writing (draft). Seungwon Shin led funding acquisition, methodology,
923 resources, and software, and supported conceptualization, data curation, formal analysis,
924 investigation, validation, visualization, and writing (draft). Damir Juric and Jalel Chergui
925 supported conceptualization, data curation, formal analysis, funding acquisition,
926 investigation, methodology, resources, software, validation, visualization, and writing (draft).
927 All authors equally contributed to reviewing/editing the manuscript.

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935

936 **Data availability**

937 The data that support the findings of this study are available from the corresponding author, I.
938 Yoon, upon reasonable request.

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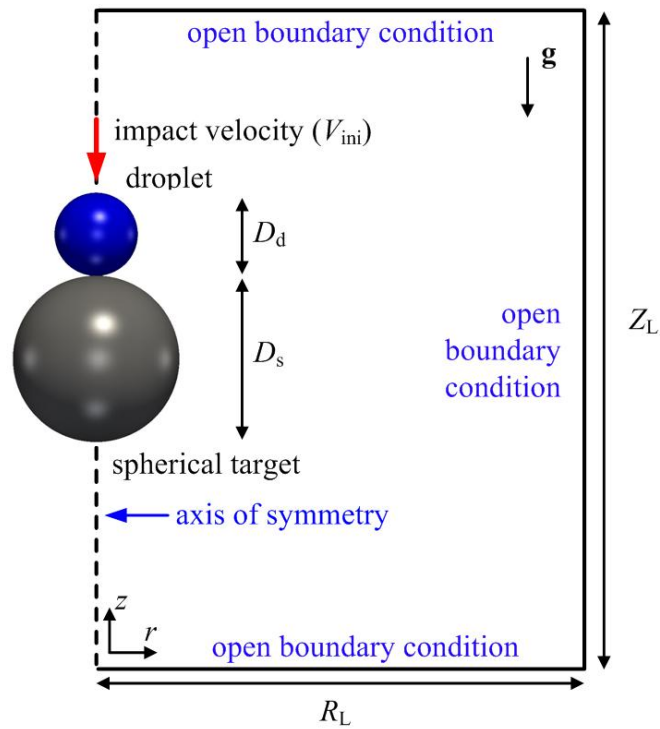
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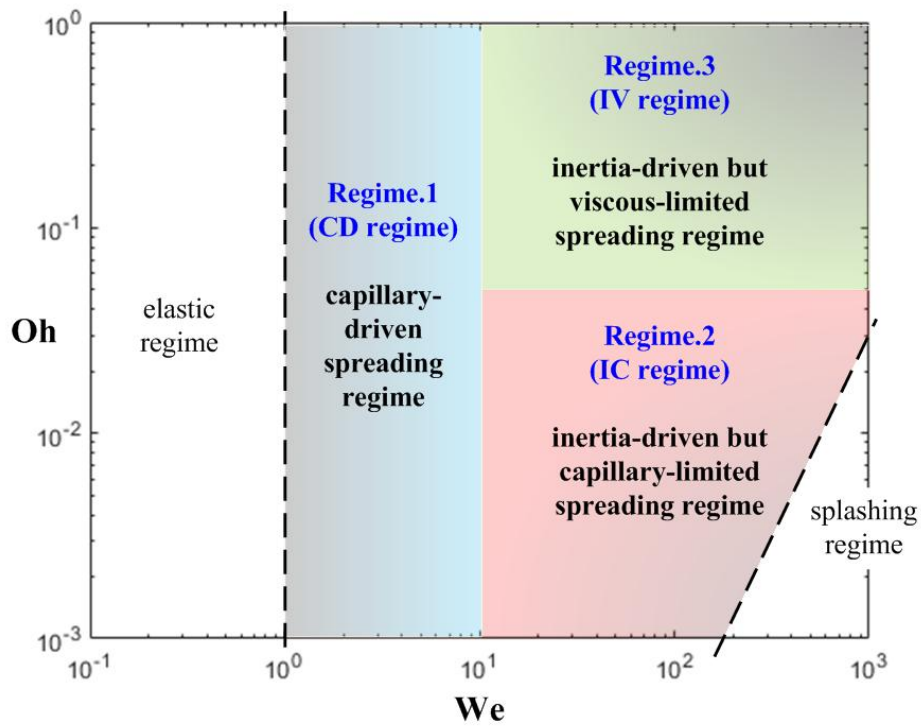
FIGURES



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Fig. 1. Simulation geometry and boundary conditions for the droplet impact system.

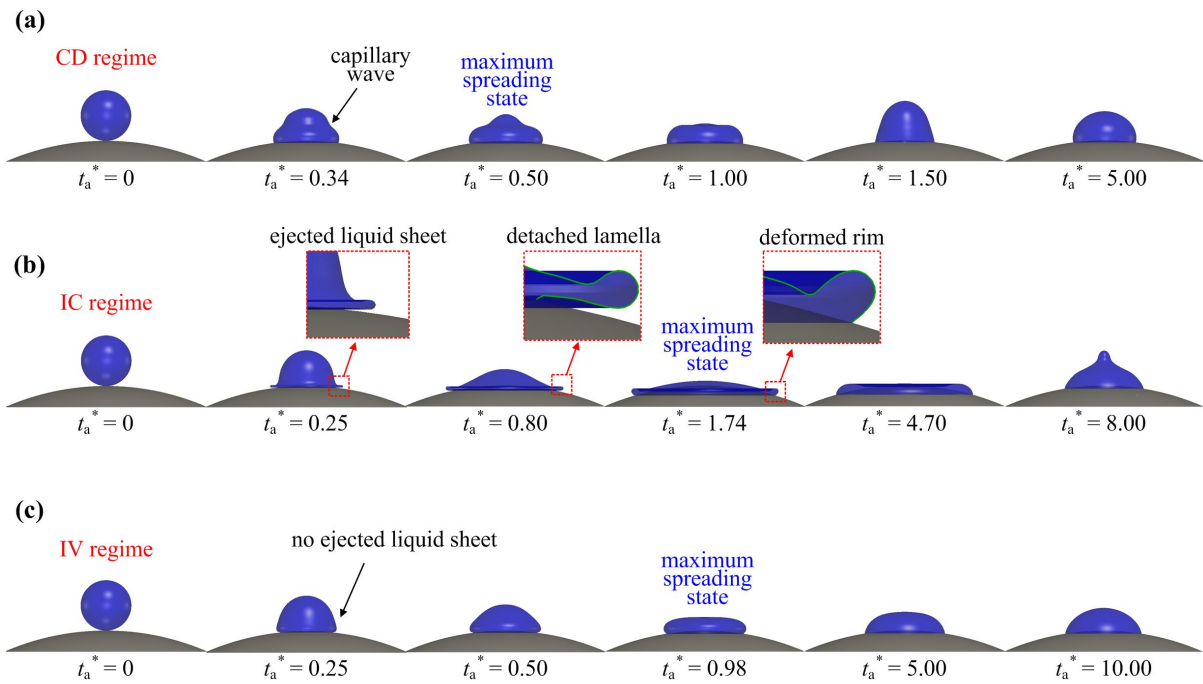
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Fig. 2. Schematic regime-map for droplet spreading dynamics. Three different spreading regimes, i.e., the capillary-driven (CD), the inertia-driven but capillary-limited (IC), and the inertia-driven but viscous-limited (IV) regimes are marked by the blue, red and green areas, respectively.

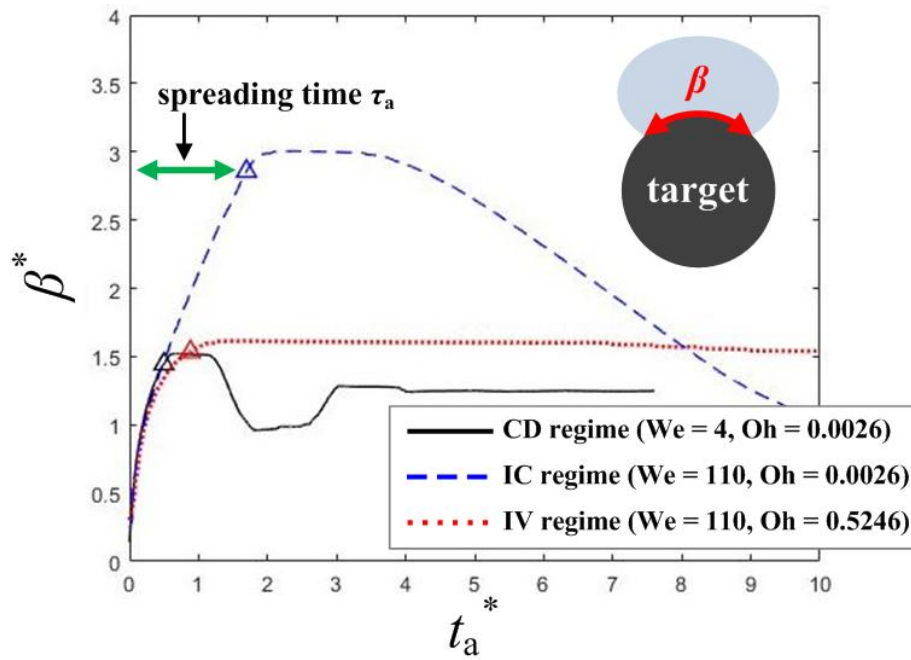
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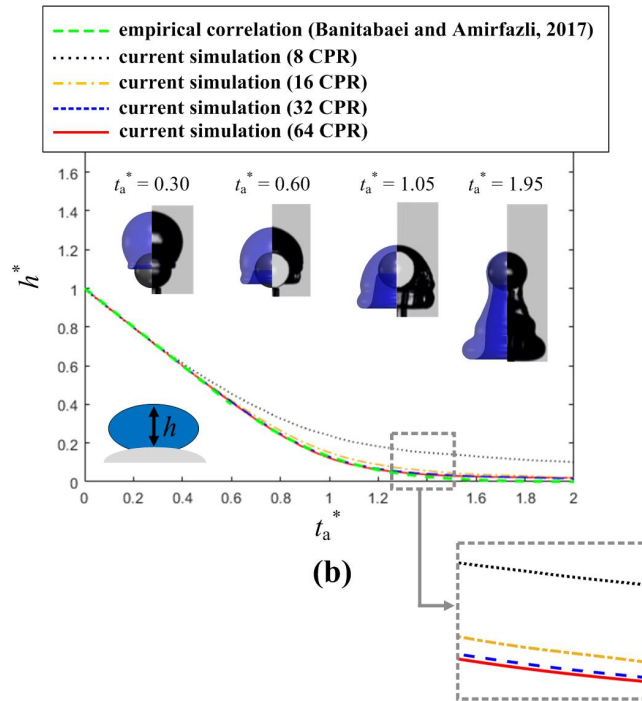
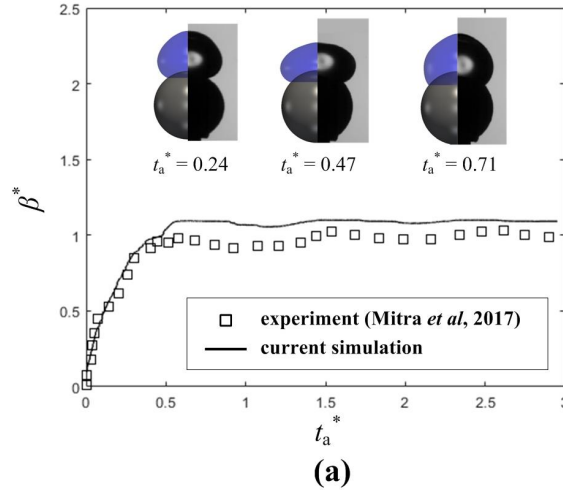
Fig. 3. Typical evolution of droplet interfacial morphology for the three different spreading regimes. (a) CD regime ($We = 4, \theta_{eqi} = 90^\circ$ and $Oh = 0.0026$). (b) IC regime ($We = 110, \theta_{eqi} = 90^\circ$ and $Oh = 0.0026$). (c) IV regime ($We = 110, \theta_{eqi} = 90^\circ$ and $Oh = 0.5246$). In these cases, $\Omega = 1/10$ for generality.

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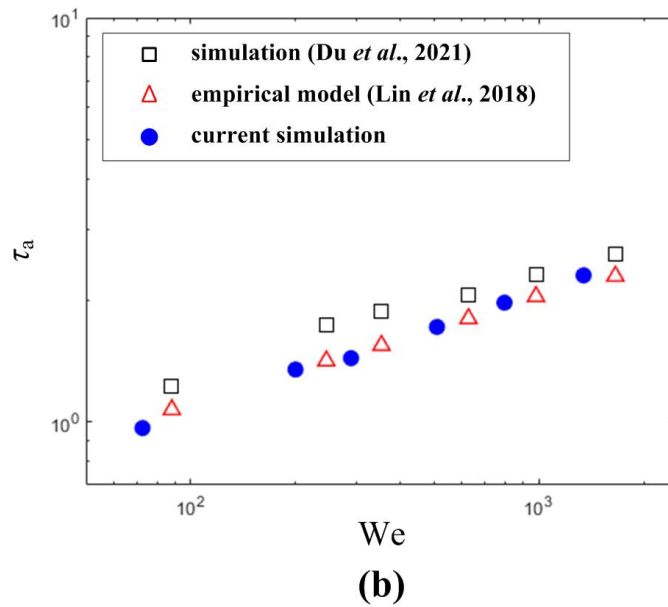
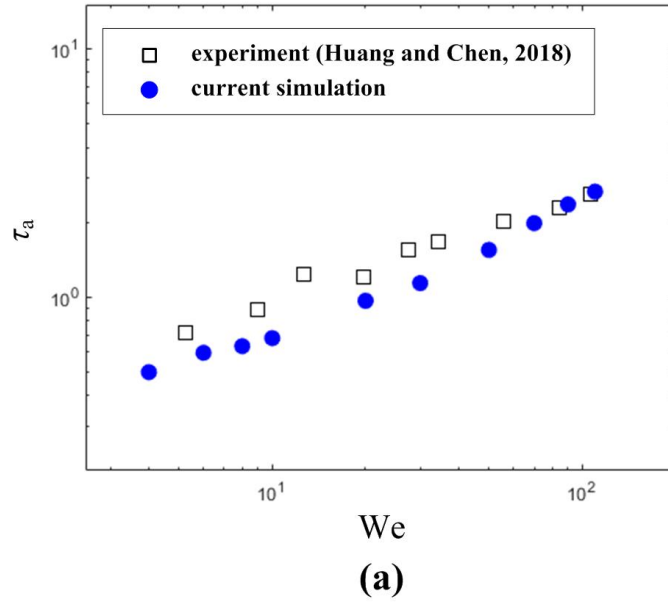
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Fig. 4. Typical evolution of the non-dimensional spreading extent, β^* , for the three different spreading regimes. All collision conditions for each regime are the same as Fig.3. The inset depicts the schematic for measuring β^* . The spreading times (the time that a droplet attains 95 % of the maximum value of β^*) are also marked by the triangles.



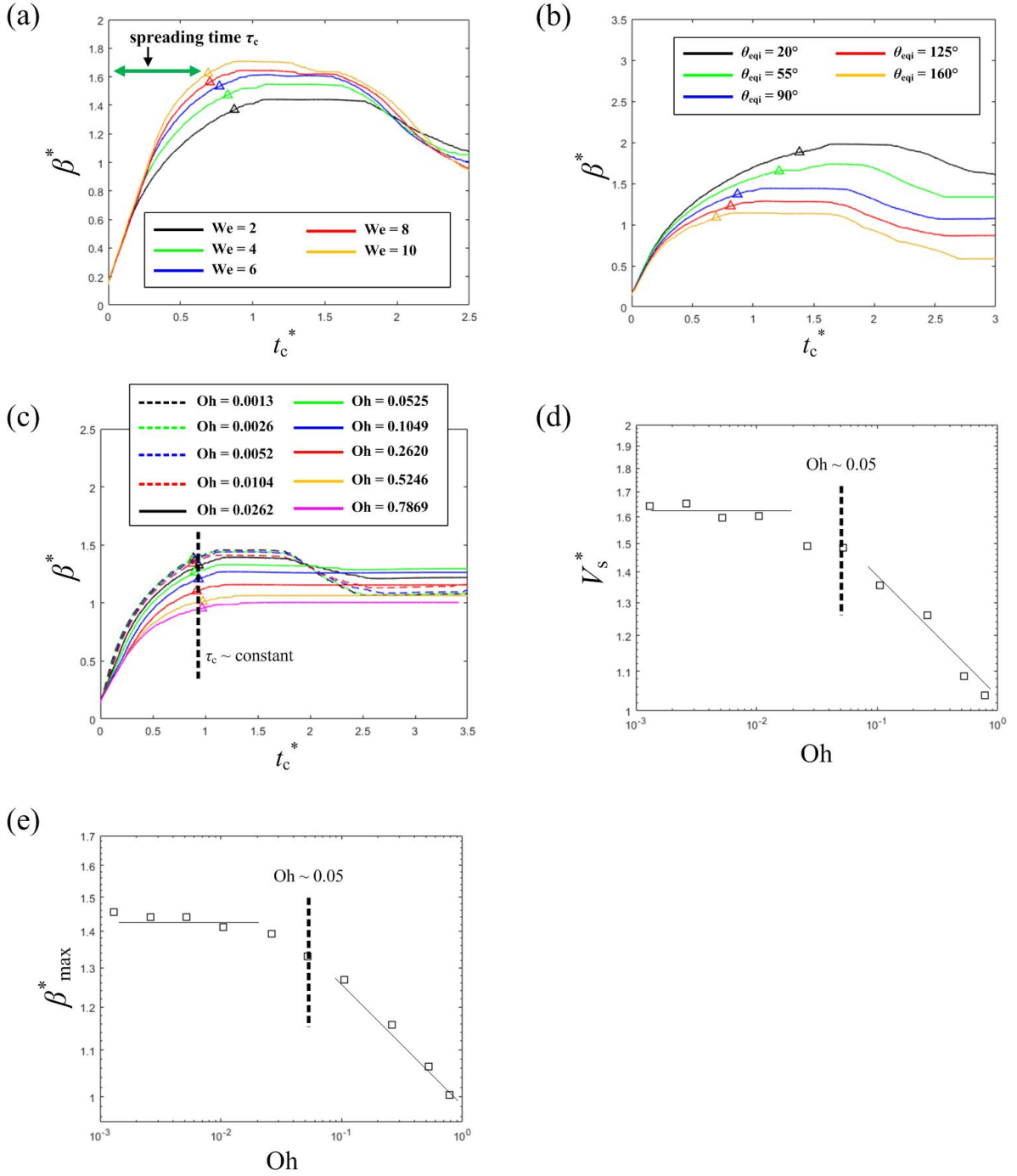
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Fig. 5. Benchmark tests. (a) Evolution of the non-dimensional spreading extent, β^* , for a droplet spreading on a small **spherical target** at low We ($We = 0.9$, $\theta_{eqi} = 85^\circ$, $Oh = 0.0026$ and $\Omega = 0.83$). The insets compare the interfacial shapes between the experiment [2] and the current simulation for 3 different time instants ($t_a^* = 0.24, 0.47, \text{ and } 0.71$). Reproduced with permission from S. Mitra *et al.*, “Interactions in droplet and particle system of near unity size ratio,” *Chem. Eng. Sci.* **170**, 154–175 (2017). Copyright 2017 Elsevier, Ltd. (b) Evolution of the non-dimensional film thickness, h^* , for a droplet spreading on a small **spherical target** at high We ($We = 155.5$, $\theta_{eqi} = 97^\circ$, $Oh = 0.0026$ and $\Omega = 1.7$) simulated by using 4 different grid resolutions (8, 16, 32 and 64 CPR). The insets compare the interfacial shapes between the experiment [71] and the current simulation for 4 different time instants ($t_a^* = 0.30, 0.60, 1.05 \text{ and } 1.95$), whereas the other inset in the lower-left corner depicts the schematic for measuring h . Reproduced with permission from S. A. Banitabaei and A. Amirfazli, “Droplet impact onto a solid sphere: Effect of wettability and impact velocity,” *Phys. Fluids* **29**, 062111 (2017). Copyright 2017 AIP Publishing LLC.



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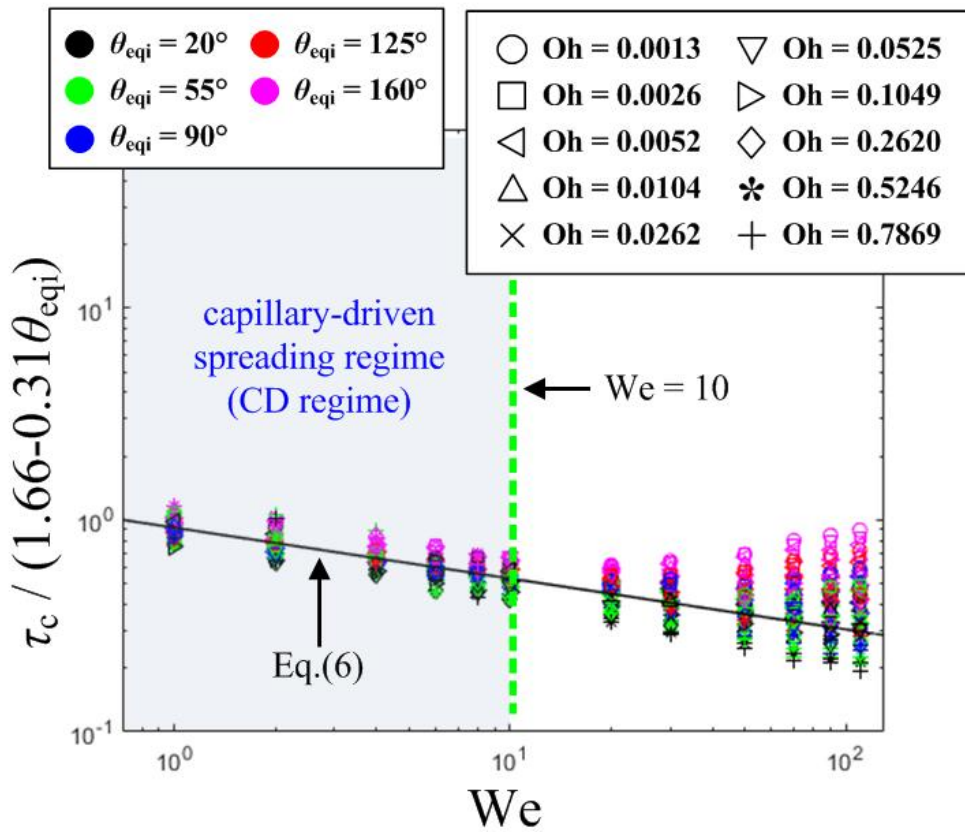
Fig. 6. Benchmark tests. (a) Non-dimensional advective spreading time, τ_a , at low Oh (Oh = 0.0026). The existing experimental result [18] is marked by black squares. (b) τ_a at high Oh (Oh = 0.62). The existing results from the numerical simulation [21] and the empirical correlation [19] are marked by black squares and red triangles, respectively.



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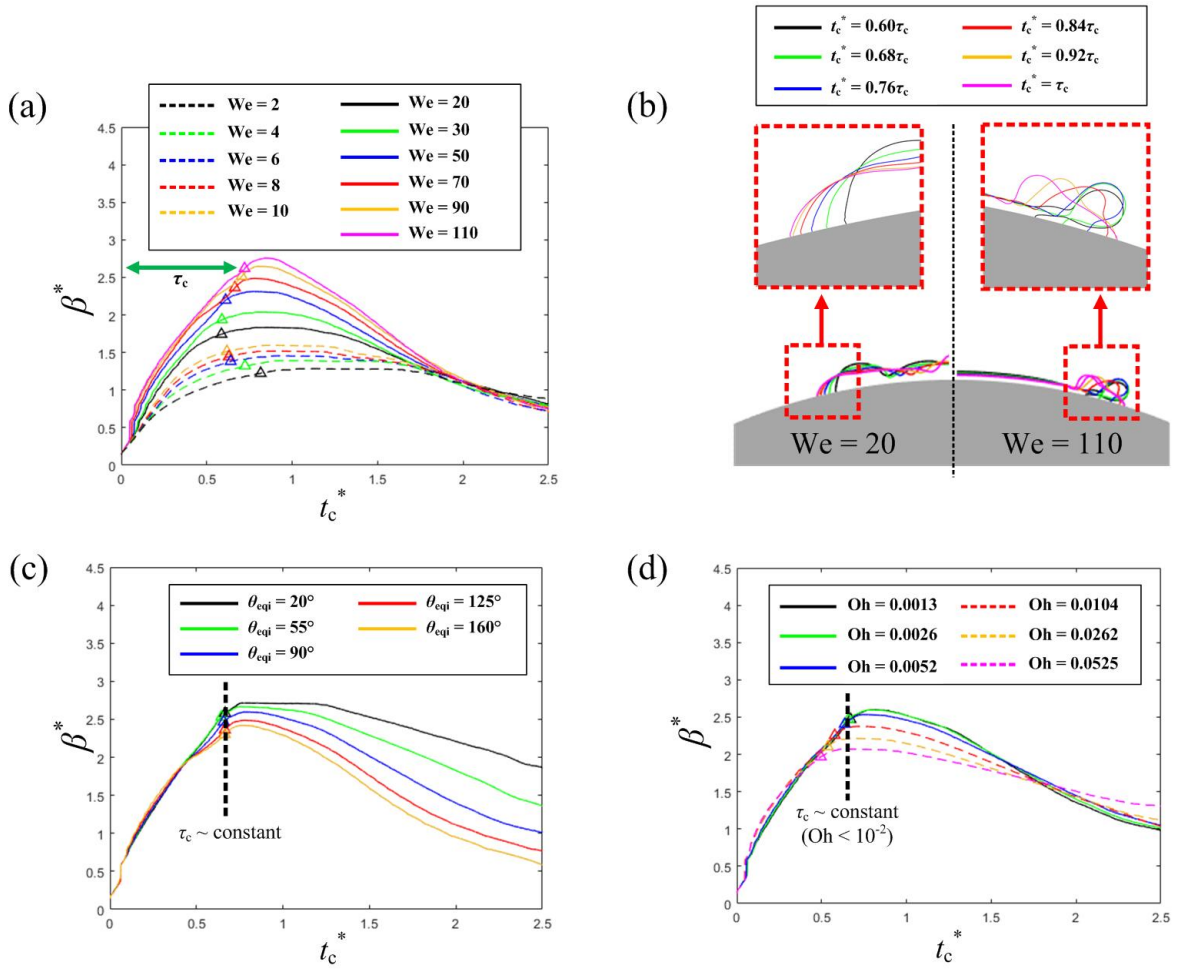
Fig. 7. Effect of the impact parameters on τ_c for the CD regime. (a) Effect of We ($We = 2-10$, $\theta_{eqi} = 90^\circ$ and $Oh = 0.0026$). (b) Effect of θ_{eqi} ($We = 2$, $\theta_{eqi} = 20^\circ-160^\circ$ and $Oh = 0.0026$). (c) Effect of Oh ($We = 2$, $\theta_{eqi} = 90^\circ$ and $Oh = 0.0013-0.7869$). (d) Effect of Oh on the non-dimensional spreading velocity, V_s^* . (e) Effect of Oh on the non-dimensional maximum spreading extent, β_{max}^* . For (d) and (e), all collision conditions are the same as Fig.7(c).

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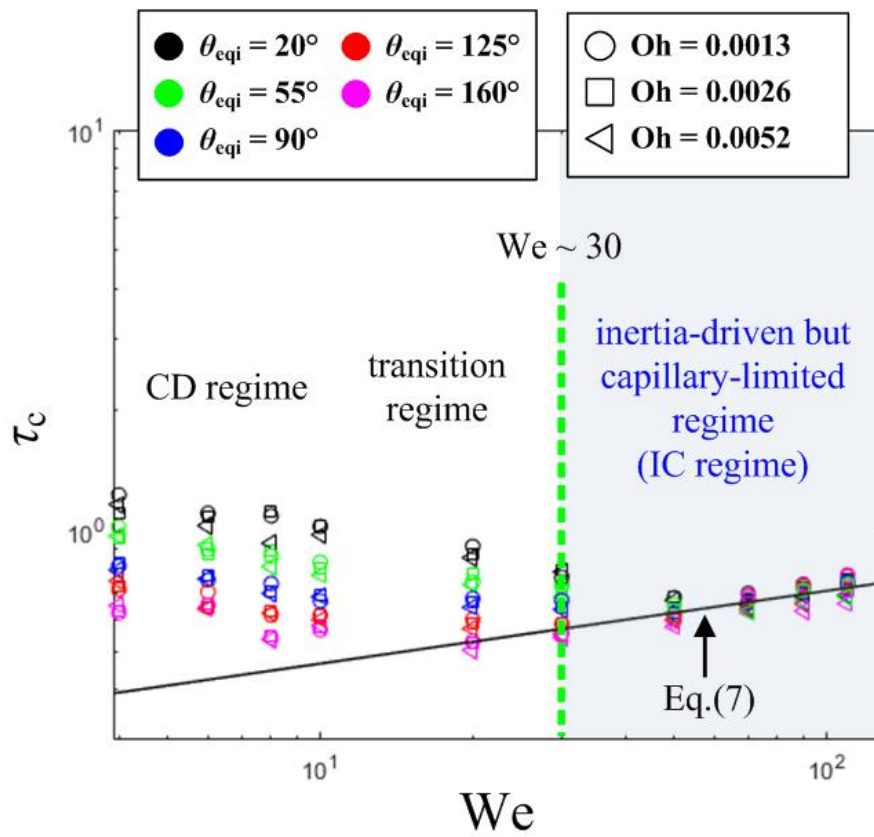
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Fig. 8. A total of 600 τ_c data ($We = 1-110$, $\theta_{eqi} = 20^\circ-160^\circ$ and $Oh = 0.0013-0.7869$). The proposed prediction model, i.e., Eq.(6) for the CD regime is indicated by the black solid line. The green dashed line depicts the threshold for the CD regime.

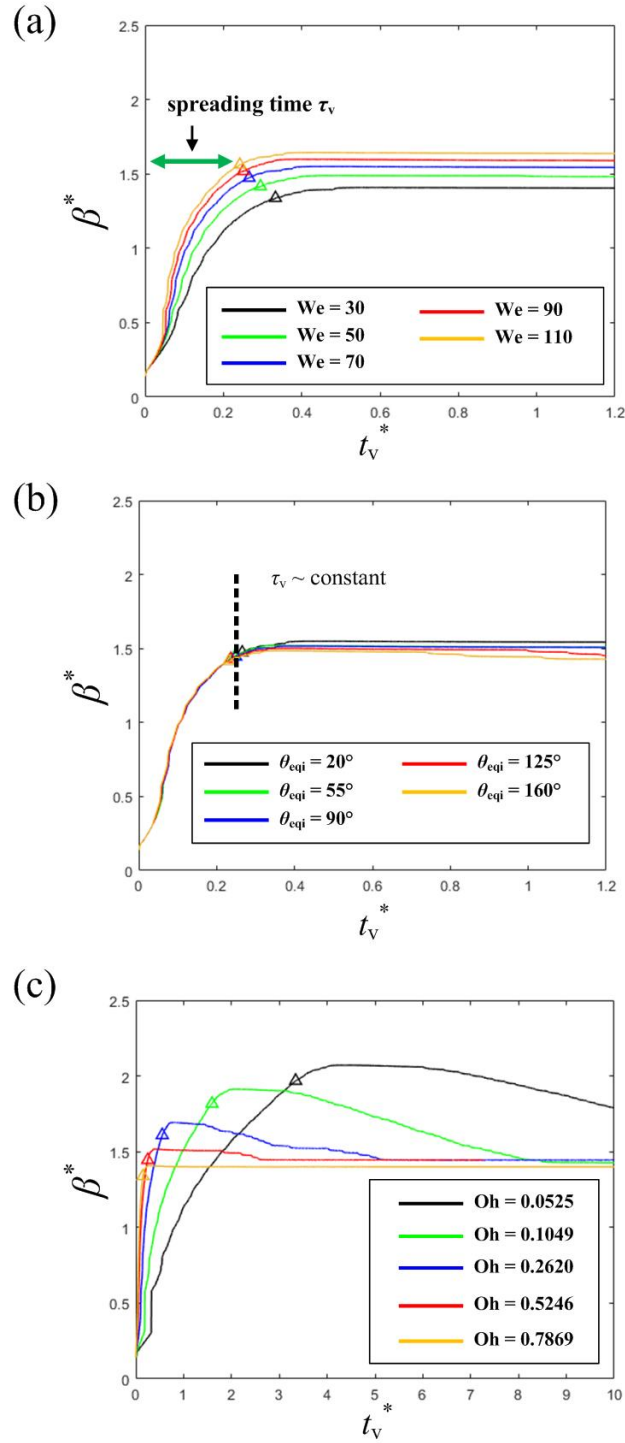


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Fig. 9. Effect of the impact parameters on τ_c for the IC regime. (a) Effect of We (We = 20–110, $\theta_{eqi} = 125^\circ$ and Oh = 0.0026). (b) Effect of We on the interfacial evolution of a droplet and its lamella deformation (We = 20 and 110, $\theta_{eqi} = 20^\circ$ and Oh = 0.0026). (c) Effect of θ_{eqi} (We = 90, $\theta_{eqi} = 20^\circ$ – 160° and Oh = 0.0026). (d) Effect of Oh (We = 70, $\theta_{eqi} = 90^\circ$ and Oh = 0.0013–0.0525).



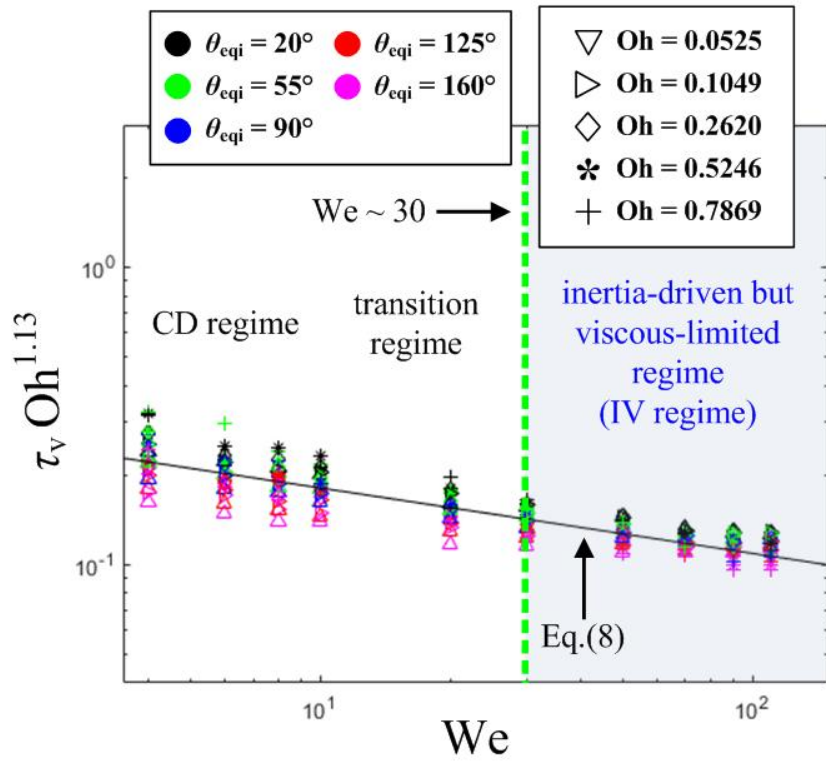
1454 **Fig. 10.** A total of 150 τ_c data ($We = 4\text{--}110$, $\theta_{\text{eqi}} = 20^\circ\text{--}160^\circ$ and $\text{Oh} = 0.0013\text{--}0.0052$). The
 1455 proposed prediction model, i.e., Eq.(7), for the IC regime is indicated by the black solid line.
 1456 The green dashed line depicts the threshold for the IC regime.
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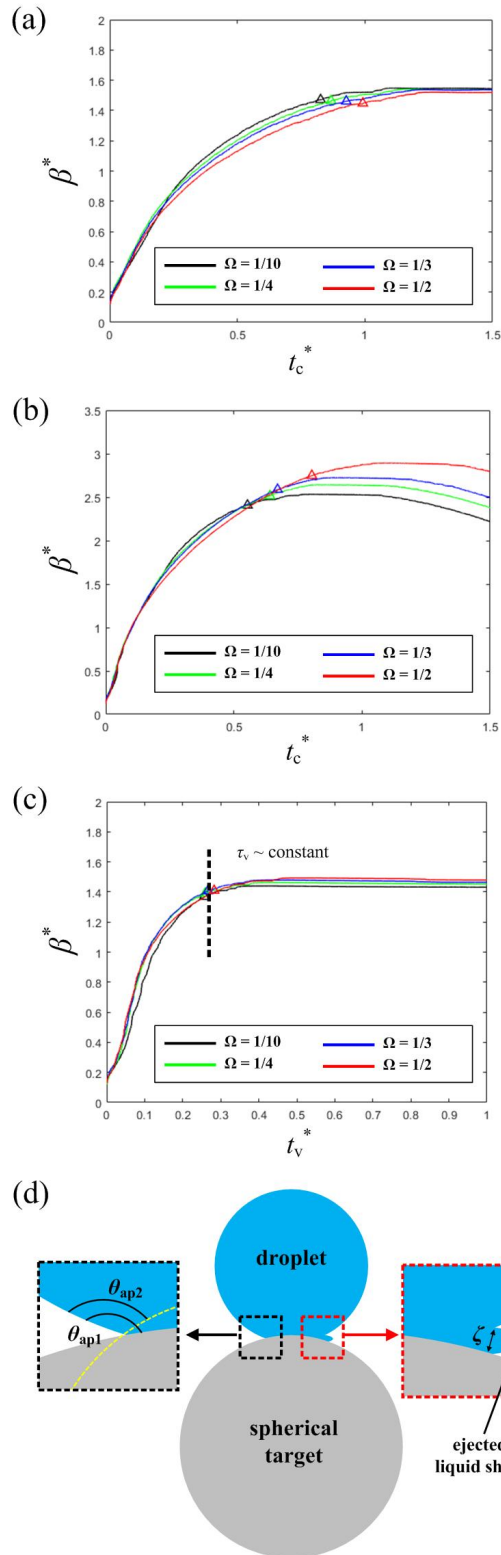
Fig. 11. Effect of the impact parameters on τ_v for the IV regime. (a) Effect of We (We = 30–110, $\theta_{eqi} = 20^\circ$ and Oh = 0.5246). (b) Effect of θ_{eqi} (We = 70, $\theta_{eqi} = 20^\circ$ – 160° and Oh = 0.5246). (c) Effect of Oh (We = 70, $\theta_{eqi} = 90^\circ$ and Oh = 0.0525–0.7869).

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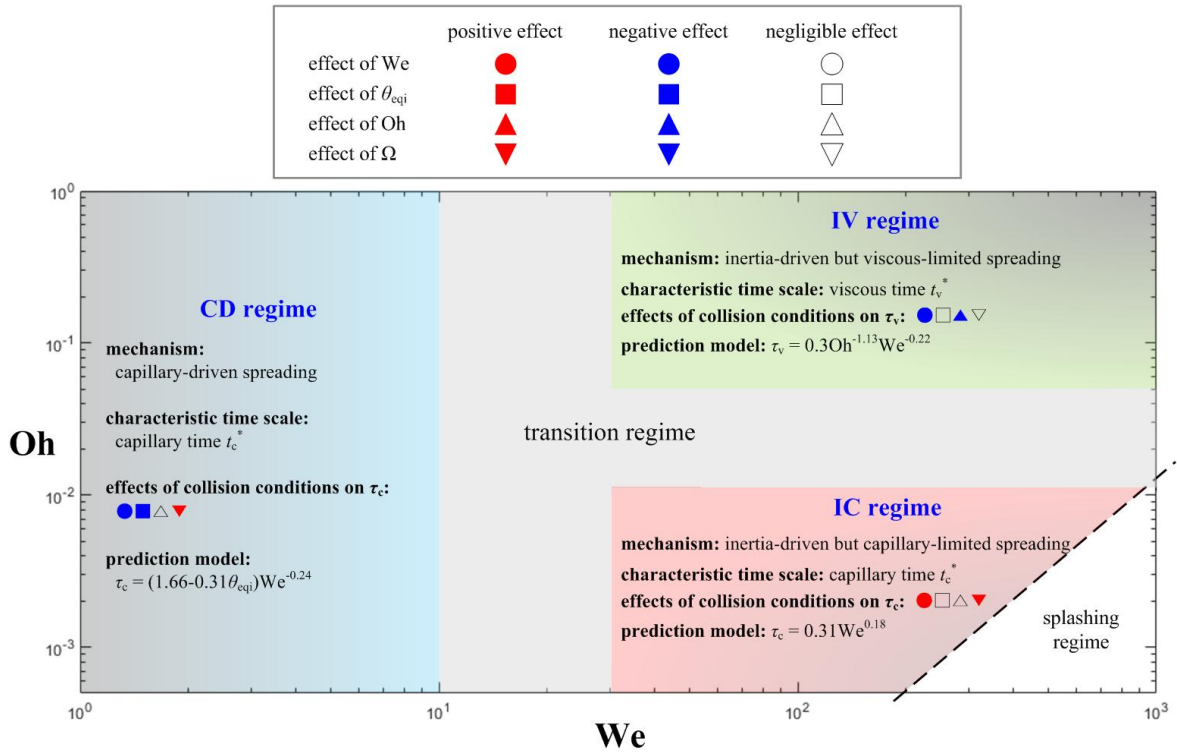
Fig. 12. A total of 250 τ_c data ($We = 4\text{--}110$, $\theta_{eqi} = 20^\circ\text{--}160^\circ$ and $Oh = 0.0525\text{--}0.7869$). The proposed prediction model, i.e., Eq.(8), for the IV regime is indicated by the black solid line. The green dashed line depicts the threshold for the IV regime.



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Fig. 13. (a) Effect of Ω on τ_c for the CD regime ($We = 4$, $\theta_{eqi} = 90^\circ$, and $Oh = 0.0026$). (b) Effect of Ω on τ_c for the IC regime ($We = 50$, $\theta_{eqi} = 90^\circ$ and $Oh = 0.0026$). (c) Effect of Ω on τ_v for the IV regime ($We = 50$, $\theta_{eqi} = 90^\circ$ and $Oh = 0.5246$). (d) Schematics of the initial apparent contact angles (on the left side) and the lamella characteristics (on the right side).

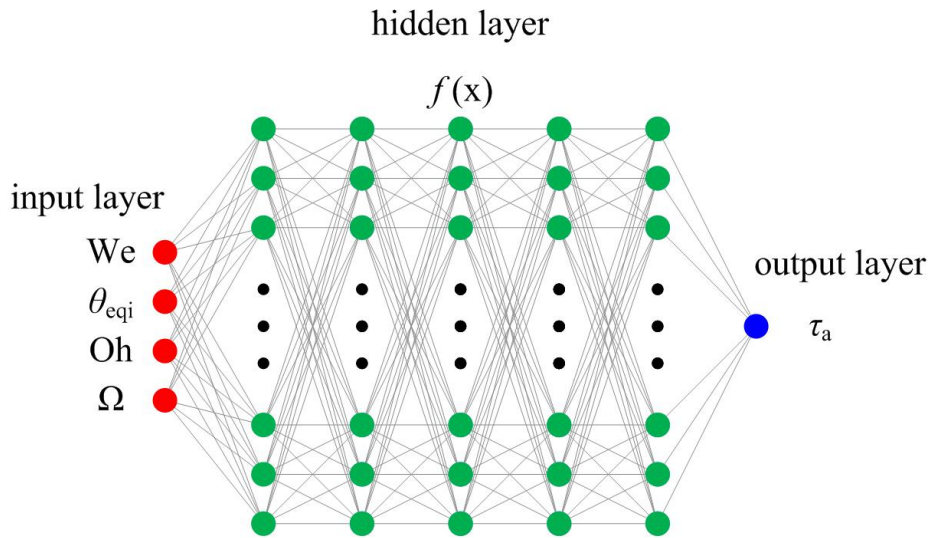
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Fig. 14. Refined regime-map of the droplet spreading mechanism. Three different spreading regimes, i.e., CD, IC, and IV regimes with the transition area are depicted. The physical mechanism driving the spreading process, the characteristic time scale, the effects from 4 impact parameters (We, θ_{eqi} , Oh, and Ω) and the prediction models are also presented for each regime.

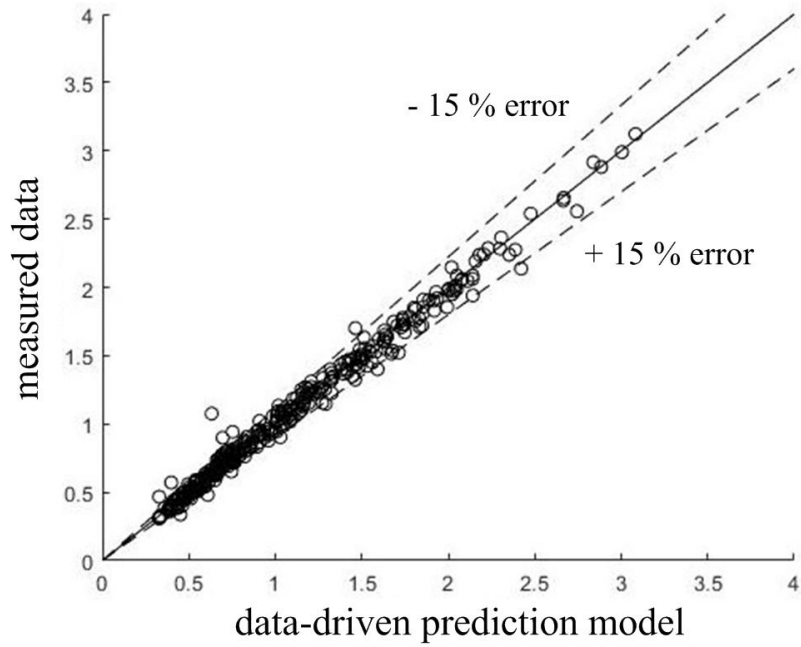
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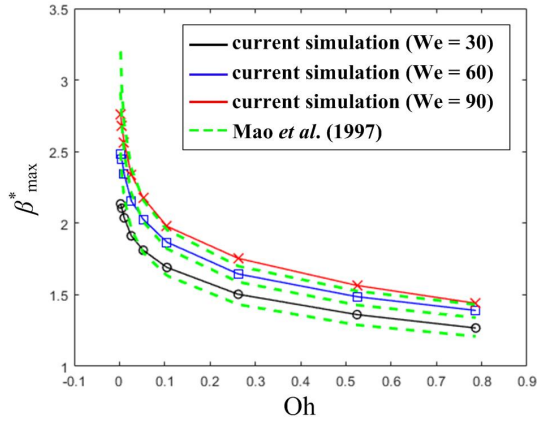
Fig. 15. Schematic diagram of the multilayer perceptron (MLP) used in the current study.

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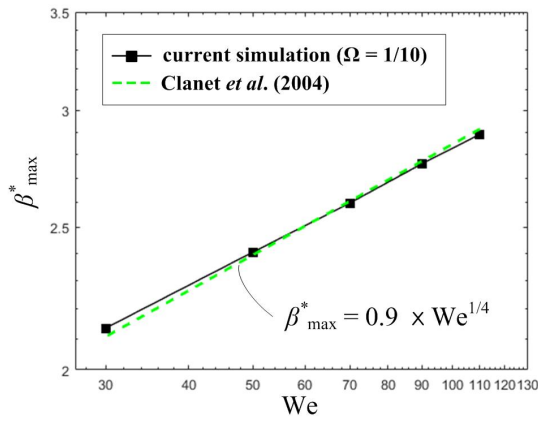


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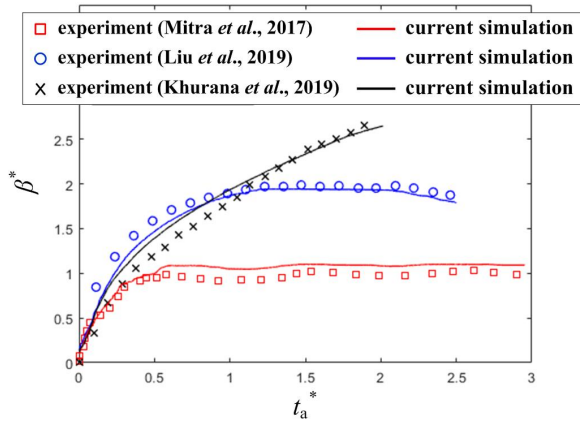
Fig. 16. Comparison of the spreading time, τ_a , between predicted data by using the data-driven (MLP) model and true (measured) data obtained by our simulations. The deviation range of $\pm 15\%$ is marked by the two dashed-black lines.



(a)



(b)

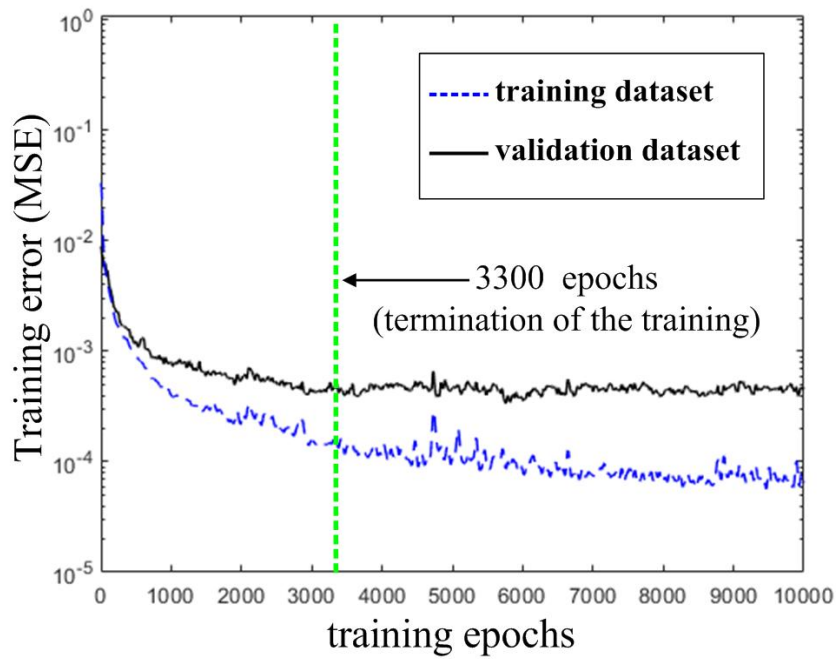


(c)

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Fig. 17. Validation comparisons between our simulation results and existing experimental results. (a) Comparison with Mao *et al.*'s [25] semi-empirical model for droplet impact with a flat surface ($30 \leq We \leq 90$, $0.0013 \leq Oh \leq 0.7869$). (b) Comparison with Clanet *et al.*'s [16] empirical scaling law for droplet impact with a flat surface ($30 \leq We \leq 110$, $Oh = 0.0026$). (c) comparisons with the experimental data of Mitra *et al.* (red squares; $We = 0.9$, $\Omega = 0.83$, $\theta_{eqi} = 85^\circ$, $Oh = 0.0024$) [2], Liu *et al.* (blue circles; $We = 19.2$, $\Omega = 0.15$, $\theta_{eqi} = 93^\circ$, $Oh = 0.0022$) [38], and Khurana *et al.* (black crosses; $We = 35.3$, $\Omega = 0.60$, $\theta_{eqi} = 85^\circ$, $Oh = 0.0022$) [39] for droplet impact with a spherical target. Reproduced with permission from I. Yoon and S. Shin [40], "Maximal spreading of droplet during collision on particle: effects of liquid viscosity and surface curvature," *Phys. Fluids* 33, 083310 (2021). Copyright 2021 AIP Publishing LLC.

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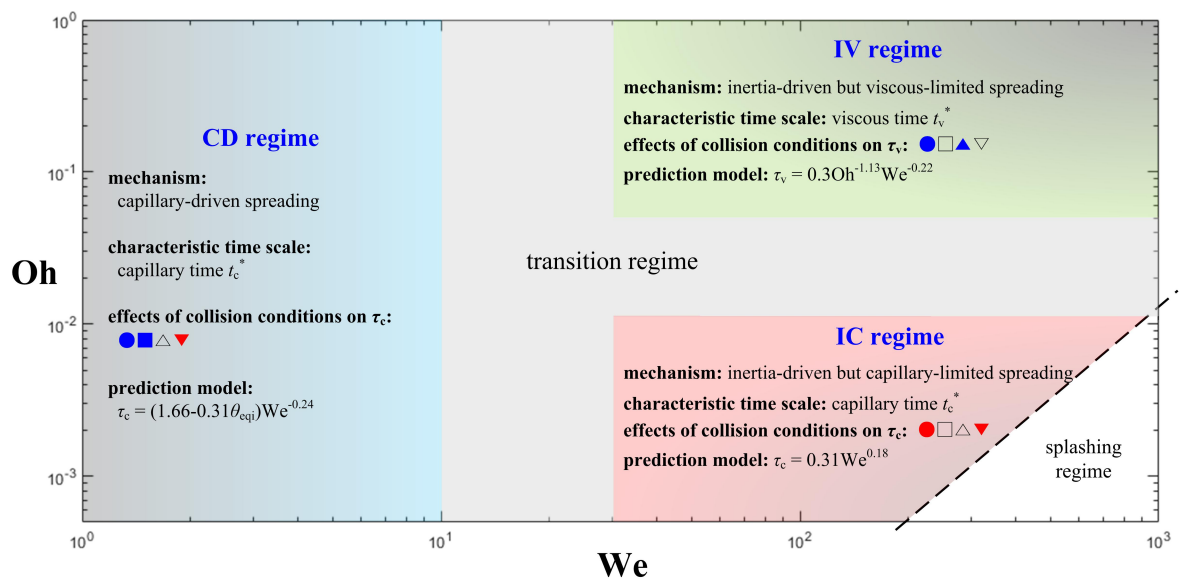
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Fig. 18. Training errors (MSE) vs. training epochs for training and validation datasets. The training is terminated at 3300 training epochs where the MSE for the validation dataset (see black solid line) is sufficiently converged. The error for the validation dataset is 4.62×10^{-4} .

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GRAPHICAL ABSTRACT

	positive effect	negative effect	negligible effect
effect of We	●	●	○
effect of θ_{eqi}	■	■	□
effect of Oh	▲	▲	△
effect of Ω	▼	▼	▽



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