

Hydrodynamic theory of wetting by active particles

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The accumulation of self-propelled particles on repulsive barriers is a widely observed feature in active matter. Despite being implicated in a broad range of biological processes, from biofilm formation to cytoskeletal movement, wetting of surfaces by active particles remains poorly understood. In this work, we study this active wetting by considering a model comprising an active lattice gas, interacting with a permeable barrier under periodic boundary conditions, for which an exact hydrodynamic description is possible. We consider a hydrodynamic scaling limit that eliminates dynamical noise while retaining microscopic interpretability, enabling a precise characterization of steady states and their transitions. We demonstrate that the accumulation of active particles has remarkable similarities to equilibrium wetting, and that active wetting transitions display all the salient characteristics of the equilibrium critical wetting transition—despite fundamental differences in underlying microscopic dynamics. However, our framework also enables the investigation of subtle but important nonequilibrium effects in active wetting, including a spontaneous ratchet effect which leads to a global steady-state current, departure of the bulk densities from their binodal values, and a modified dynamical transition pathway. Our results provide an intrinsically nonequilibrium framework in which to study active wetting, precisely demonstrating the connection to passive wetting while clarifying the nonequilibrium consequences of activity.

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I. INTRODUCTION

Active matter comprises a broad class of self-propelled particles which consume energy at a local scale to produce directed motion. Such particles are extremely common in microbiological and ecological contexts, spanning length scales from microns to meters and involving particle numbers from dozens to millions [1–3]. Examples include swimming bacteria [4–6], flocking birds [7,8], and cytoskeletal networks [2,9], as well as synthetic cases such as Janus particles [10] and robotic swarms [11].

The continuous consumption of energy drives active particles out of equilibrium, leading to a diverse range of emergent collective behaviors. For example, active particles undergo motility-induced phase separation (MIPS) into dense (liquid) and dilute (vapor) phases due to the interplay between volume exclusion and self-propulsion, even in the absence of explicitly attractive interactions [12,13]. It is expected that collective behavior in active systems should show some degree of independence of the microscopic details of particle interactions [2,14], so simple particle-level [15–17] or field-theoretical

[13,18,19] models of active matter are natural settings for studying their emergent behaviors.

A distinguishing feature of active systems, compared to their passive counterparts, is the disproportionate influence of boundaries. Nonlocal effects arising from steady-state currents mean that boundaries play a much more central role in determining the collective behavior of active matter than in the passive case. Indeed, recent studies have shown that boundaries can regulate transport, transfer energy between length scales, and create or destroy large-scale structures [20–26].

One prominent boundary-driven effect is the accumulation of active particles on repulsive barriers. Recent works have shown that the self-propulsion of active particles results in a layer of excess density on any repulsive wall or barrier [20,26–40]; this behavior is reminiscent of the wetting of surfaces in equilibrium [41–43], and so the phenomenon has become known as active wetting [44]. Experiments have implicated active wetting in a wide variety of biological processes [45,46], including biofilm formation [47–51], epithelial tissue organization [52,53], and cytoskeletal movement [54].

In equilibrium systems, adsorption (a microscopically thin film of fluid near a surface) is carefully distinguished from genuine wetting transitions [41,55–58]. In particular, equilibrium critical wetting is a continuous surface phase transition on varying a surface attraction parameter at two-phase coexistence, at which the adsorbed film thickness diverges [59]. A complete wetting transition involves the divergence of the wetting film approaching bulk coexistence on the wet side of the wetting transition. Simulations have suggested that in systems undergoing MIPS, accumulation/adsorption of active

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particles [30,33,34,36] can develop into macroscopic wetting [28,31], with transitions analogous to equilibrium complete wetting [29] and critical wetting [27,35,37,38].

Despite its significant role in active systems, and biological relevance, the physics of active wetting remains poorly understood—complex interfacial phenomena and nonlocal effects mean that even basic results of equilibrium wetting theory are not necessarily applicable to the active case [20,28,31,60–62]. In particular, the phenomenology of an active interface cannot be rationalized in terms of a single effective surface tension [13,61]. These confounding effects also make it difficult to distinguish between the distinct phenomena of adsorption and critical or complete wetting—these distinctions have not always been maintained in the active wetting literature [20,21,44]. An intrinsically nonequilibrium framework is therefore required to understand these nonequilibrium surface phenomena.

In this work, we study active wetting by analyzing a lattice model of self-propelled particles whose large-scale (hydrodynamic) behavior can be derived exactly [63–65] and is governed by deterministic equations. A consequence of the hydrodynamic limit as taken here is that the large-scale behavior of our model is not affected by fluctuation corrections (in the sense of the renormalization group). This leads to critical behavior of the mean-field type. Nevertheless, our results provide generic insights for active wetting.

The hydrodynamic framework—which has also provided important recent insights in other types of active system [66,67]—allows us to use precise analytic techniques familiar from field theories while retaining a direct connection to microscopic behavior.¹ This means we can identify subtle yet physically significant effects which may be obscured in particle-level studies [27] by the confounding role of noise.

In equilibrium systems, wetting is often investigated in a slit geometry, wherein a fluid is confined between planar walls [43,68–75], simplifying the investigation of wetting transitions (see Sec. II and Appendix A). We seek an analogous geometry to study active wetting transitions. However, it is known that active particles generically adsorb to hard walls [30,33,34,36], leading to full wetting only [27–29,31,35]. We therefore use a smooth, permeable barrier (as in [27,35]) to access partial wetting states. We take periodic boundaries to suppress any other boundary-driven phenomena and use fixed number of particles (see Sec. III A) such that the model is always at liquid-vapor coexistence (the transition we observe is therefore analogous to *critical* rather than *complete* wetting). Our deterministic equations result in a system which remains translationally symmetric in the direction(s) parallel to the wall, and so we study one-dimensional (1D) profiles—although our hydrodynamic model admits a general and exact description even in higher dimensions.

Using this framework, we demonstrate remarkable similarities between active and equilibrium wetting.² Specifically, we show that both full- and partial-wetting states arise as a consequence of activity, separated by a critical wetting transition. These observations are surprising because active wetting takes place without any explicitly attractive particle-particle or particle-wall interactions, so connections to equilibrium liquids are not at all guaranteed. Our hydrodynamic framework therefore enables us to generalize the already powerful analogy between equilibrium phase separation and active MIPS to include surface phenomena.

Our model also enables the identification of subtle but important nonequilibrium phenomena in active wetting, which arise as a consequence of activity. In particular, we demonstrate that a ratchet effect emerges on a partially wet barrier, leading to steady-state currents (as may be generically expected in nonequilibrium steady-states [20,23,28,31]). The ratchet current leads to a number of intrinsically nonequilibrium effects in active wetting, including departure of the bulk phases from their binodal densities and the emergence of a modified dynamical pathway for the wetting transition—these are explored in the companion Letter [76]. However, the scaling laws associated with critical wetting remain unchanged. Particle-based simulations also indicated that active wetting is critical in 2D, but large fluctuation effects made this hard to characterize [27]. Persistent currents were also found previously for active wetting in a different (droplet) geometry, where they affect the contact angle via a correction to Young’s equation [31]. Our results complement these previous works; in particular, the combination of a hydrodynamic approach and slit geometry leads to efficient characterization of the wetting transition, without confounding finite-size or fluctuation effects. Overall, we find active wetting states to display an equilibriumlike critical-wetting surface phase transition, but one which is also modified by the emergence of spontaneous nonequilibrium current flow.

The remainder of this paper is structured as follows: In Sec. II, we provide a short review of relevant results from the equilibrium theory of wetting. In Sec. III, we precisely define the hydrodynamic model studied here, as well as its numerical implementation. Section IV then presents the surface phase diagram and characterizes the wetting states and their transitions, as well as introducing the nonequilibrium effects evident in active wetting. Finally, Sec. V summarizes our conclusions, discussing the results presented here and in [76], and directions for continued research.

II. BACKGROUND—EQUILIBRIUM WETTING

This section reviews the equilibrium theory of wetting, in order to provide context to our work and points of comparison. We focus mainly on results derived from the Cahn theory of

¹One aspect that aids this connection is that particle density and orientation fields both appear in the hydrodynamic equations; this differs from most active field theories for isotropic fluids, which take the hydrodynamic limit in a different way that eliminates orientational dynamics but retains diffusive noise.

²The analogy with equilibrium wetting transitions requires that the barrier-free system is in a state of two-phase coexistence [27,28,31], which is the case in this work. Adsorption of active particles onto a boundary from a homogeneous fluid phase [30,33,34,36] is sometimes also called active wetting in the literature, but this is not our focus here.

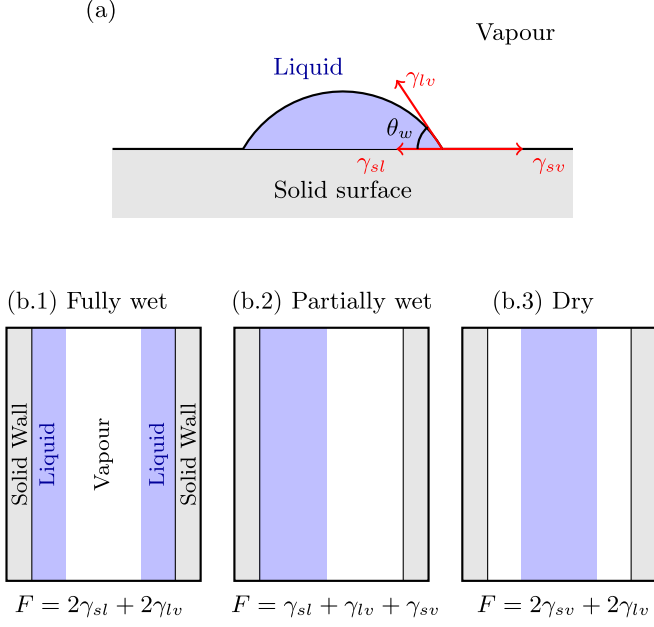


FIG. 1. (a) Diagram of equilibrium wetting on a solid surface. A liquid droplet rests on the solid in coexistence with the vapor phase, forming a contact angle θ_w . The interfacial tensions γ_{sl} , γ_{sv} , and γ_{lv} correspond to solid-liquid, solid-vapor, and liquid-vapor interfaces, respectively. (b) Schematic of wetting states in a slit geometry with a fixed mass of fluid. (b.1) Fully wet: Macroscopic liquid films coat both walls. (b.2) Partially wet: A macroscopic liquid film forms on one side. (b.3) Dry (unpinned): No bulk liquid wets the walls.

wetting [77], as this (mean-field, noise-free) model is most analogous to our exact hydrodynamic description.

Wetting describes the equilibrium behavior of a fluid with coexisting liquid and vapor phases when it comes into contact with an attractive solid wall [41–43,57,77]. The liquid typically forms droplets on the wall, characterized by a contact angle θ_w [see Fig. 1(a)]. This contact angle can be calculated from the interfacial tensions between the liquid, vapor, and solid phases, and is given by Young’s equation [41]

$$\cos(\theta_w) = \frac{\gamma_{sv} - \gamma_{sl}}{\gamma_{lv}}, \quad (1)$$

where γ_{lv} , γ_{sv} , and γ_{sl} denote the liquid–vapor, solid–vapor, and solid–liquid interfacial tensions, respectively. Young’s equation can be derived either mechanically, by considering the balance of forces at the contact line [41], or using statistical physics, via the variational minimization of a free energy [78]. These perspectives are equivalent because the system is at equilibrium.

When $\theta_w \in (0, \pi)$, the liquid forms a droplet on the wall; this is called partial wetting. A wetting transition occurs when the contact angle θ_w vanishes, and the liquid spreads across the wall to form a macroscopic film; this state with $\theta_w = 0$ is called full wetting. A drying transition occurs as θ_w approaches π and the droplet detaches from the wall.

Wetting can alternatively be studied in the context of a fluid in a slit geometry at constant total mass (i.e., a canonical ensemble) [68–71,74,75], where the fluid is confined between two parallel planar walls, separated by a macroscopic distance

L . The system is assumed to be translationally symmetric in the direction(s) parallel to the walls; the total fluid mass, and other control parameters, are chosen such that both a bulk liquid and bulk vapor phase coexist. This allows for three possible wetting states (see Fig. 1(b) and [43,68,74] for a more detailed discussion):

(i) *Fully wet*: The fluid forms a macroscopic liquid layer on both walls with a vapor bulk between them, such that the total energy per unit (transverse) area is $2\gamma_{sl} + 2\gamma_{lv}$.

(ii) *Partially wet*: The fluid forms a macroscopic liquid layer on one wall while the other remains dry,³ such that the total energy per unit area is $\gamma_{sl} + \gamma_{lv} + \gamma_{sv}$.

(iii) *Dry (unpinned)*: Both walls are dry, and there is an unpinned liquid bulk between the walls such that the total energy per unit area is $2\gamma_{sv} + 2\gamma_{lv}$.

Comparing the energies, the transition from partial to full wetting in the slit geometry occurs at $\gamma_{sv} = \gamma_{sl} + \gamma_{lv}$, which corresponds to $\theta_w = 0$, as for the droplet. The wetting transitions in the slit and droplet geometry therefore correspond to the same physical transition. The same applies to the drying transition between states (ii) and (iii). We emphasize that a wet wall is covered by a macroscopic liquid layer of thickness $\zeta = \mathcal{O}(L)$. For attractive walls, even a dry wall is typically coated by a microscopic film of fluid, $\zeta = \mathcal{O}(1)$ [41].

Wetting transitions occur as some control parameter is varied (for example, the strength of an attractive interaction between fluid and wall); we denote this quantity here by ϵ . It is useful to define a spreading factor $S = \gamma_{sv} - \gamma_{sl} - \gamma_{lv}$ [41], which is equal to

$$S = \gamma_{lv}[\cos(\theta_w) - 1] \quad (2)$$

in the partially wet state, with $S < 0$. A wetting transition means that $S \rightarrow 0^-$ as a control parameter ϵ tends to a critical value ϵ^* . These transitions come in two types: first-order wetting involves $S \rightarrow 0$ with a finite slope, $dS/d\epsilon \neq 0$; in this case, the fluid film thickness ζ (on the dry side) approaches a finite limit as $S \rightarrow 0$ and then jumps to a macroscopic value $\zeta = \mathcal{O}(L)$ only in the fully wet state. Critical⁴ transitions have $S \rightarrow 0$ and $dS/d\epsilon \rightarrow 0$ together; this signals a continuous divergence of the thickness ζ , see below.

Cahn [41,77] showed that the equilibrium wetting transition of a fluid on a planar wall, or in the slit geometry, can be described at the hydrodynamic level by a one-dimensional field theory for the fluid density ρ , with free energy

$$F[\rho] = \int_0^\infty \frac{\kappa}{2} (\nabla \rho)^2 + W(\rho) + f_w(\rho) \delta(x) dx, \quad (3)$$

where $W(\rho)$ is a double-welled potential describing bulk phase separation. The free energy due to interaction with the

³ Although this dry wall has no liquid bulk attached, a thin adsorbed film remains present within the partially wet regime. This leads to the relevant moist surface tension γ_{sv} [41].

⁴ Note that these wetting transitions are critical as they involve divergence of the film width as a surface parameter is varied, at two-phase coexistence. This is distinct from the complete wetting transition which occurs approaching coexistence from a single phase, e.g., [29].

wall (at $x = 0$) is contained in the function $f_w(\rho)$, which can be expanded in a Taylor series as

$$f_w(\rho) = c_0 - c_1\rho + \frac{1}{2}c_2\rho^2 + \mathcal{O}(\rho^3). \quad (4)$$

Here c_1 roughly corresponds to the attractive interactions between the wall and the fluid, and c_2 to the effect of exclusion of fluid by the wall, resulting in the loss of attractive fluid-fluid interactions (see [41]). Taking the control parameter $\epsilon = c_1$ (so the wetting transition occurs on varying the attraction strength), $\epsilon < \epsilon^*$ corresponds to partial wetting (weaker attraction) and $\epsilon > \epsilon^*$ to full wetting.

The wetting transition in a slit geometry, where impenetrable walls are placed at $x = 0, L$, can then be characterized using a global asymmetry order parameter [27]:

$$\mathcal{A}[\rho] = \frac{\int_0^L |\rho(x) - \rho(L-x)| dx}{2L(\bar{\rho} - \rho_v)}. \quad (5)$$

Here $\mathcal{A} = 0$ characterizes full wetting, $\mathcal{A} > 0$ indicates partial wetting, and the normalization is such that $\mathcal{A} \in [0, 1]$.

Minimizing the free energy (3) results in an equilibrium theory of wetting with mean-field character (because the hydrodynamic approach neglects density fluctuations). This theory is summarized in Appendix A; both critical and first-order wetting are possible in the Cahn theory, with critical wetting observed if c_2 is sufficiently large. Physically, this corresponds to the scenario where the effect of the wall in preventing attractive interactions between nearby fluid particles is large.

In later sections, we will compare the active wetting transition to the equilibrium critical wetting transition. To facilitate that discussion, we identify here the main distinguishing features of critical wetting at equilibrium (as opposed to first-order wetting), which can be derived from the Cahn theory in a slit with constant total mass. First, the film thickness on the dry wall diverges continuously at the wetting transition, as

$$\zeta \sim \ln \frac{1}{\epsilon^* - \epsilon}, \quad \epsilon < \epsilon^*. \quad (6)$$

Second, the asymmetry parameter vanishes continuously as full wetting is approached (exhibiting a pitchfork bifurcation with a characteristic exponent of $1/2$):

$$\mathcal{A} = \begin{cases} \alpha(\epsilon^* - \epsilon)^{1/2}, & \epsilon < \epsilon^* \\ 0, & \epsilon \geq \epsilon^*, \end{cases} \quad (7)$$

where α is a proportionality constant. Third, there are no metastable states so one does not expect hysteresis as ϵ passes through ϵ^* .

The theory of wetting has been extended to describe numerous surface phenomena beyond those outlined above. See [42] for a modern review. For example, away from liquid-vapor coexistence in a grand canonical ensemble, prewetting or complete wetting transitions may still be observed (see [55–57, 79, 80]) or slits may fill entirely with liquid in a phenomenon known as capillary condensation [81, 82]. These effects are not relevant for our model as we work at liquid-vapor coexistence in the canonical ensemble. We include in Appendix A 2 a discussion of the distinct surface phenomena which are possible in equilibrium, and in Appendix A 3 a discussion of the different partially wet morphologies which

are possible in a slit geometry with two spatial dimensions. Noise, dimensionality, and the nature and range of particle-particle and particle-wall interactions may all subtly change the order of wetting and drying transitions—see, for example, [43, 83, 84].

III. MODEL AND METHODS

A. Active lattice gas model

1. Exact hydrodynamic description

We study the active lattice gas introduced in [64], which involves particles that move on an $N \times N$ square lattice, with at most one particle per site, modified by the presence of an external potential V . We embed the lattice into a domain of size $L \times L$, so the lattice spacing is L/N . The hydrodynamic limit is $N \rightarrow \infty$ at fixed L . Each particle i is endowed with a position $\mathbf{x}_i$ and an orientation vector $\mathbf{e}(\theta_i) = (\cos(\theta_i), \sin(\theta_i))^T$. It hops from $\mathbf{x}_i$ to a vacant neighboring site $\mathbf{x}_i + \mathbf{u}$ with rate

$$w_i(\mathbf{u}) = N^2 \left\{ D_E + \frac{v_0}{2} \mathbf{u} \cdot \mathbf{e}(\theta_i) + \frac{V(\mathbf{x}_i) - V(\mathbf{x}_i + \mathbf{u})}{2} \right\}. \quad (8)$$

(Note that $|\mathbf{u}| = L/N$ is small in the hydrodynamic limit.) Each angle θ_i undergoes Brownian motion on the unit circle with orientational diffusion constant D_O , and spatial diffusion of particle positions $\mathbf{x}_i$ occurs with diffusion constant D_E (as above). The only interactions between particles are excluded volume interactions. On taking the hydrodynamic limit, it is important that particles' movement rates are chosen to have a nontrivial N -dependence, so both diffusive and self-propelled contributions remain comparable as $N \rightarrow \infty$. This feature enables the exact calculation; it involves a limit of fast diffusion that suppresses fluctuations so the resulting system is of mean-field type, similar to the Cahn theory. We note that, due to this suppression of noise-induced fluctuations in the hydrodynamic limit, the bubbly phases seen in some active systems [13, 19, 85] are not possible here, even though bubbles may remain present in the limit of a large system size for some parameter ranges [86] (see also Sec. V B). Our dynamics also exclude inertial effects; while inertia is known to influence wetting behavior in active systems [32], this is not considered here.

The system is described in the hydrodynamic limit by a probability density $f(\mathbf{x}, \theta, t)$ for particles with spatial position $\mathbf{x}$ and orientation θ at time t . The spatial density ρ and polarization $\mathbf{p}$ of the particles are

$$\rho(\mathbf{x}, t) = \int_0^{2\pi} f(\mathbf{x}, \theta, t) d\theta, \quad (9)$$

$$\mathbf{p}(\mathbf{x}, t) = \int_0^{2\pi} f(\mathbf{x}, \theta, t) \mathbf{e}(\theta) d\theta. \quad (10)$$

(Note that $\rho(\mathbf{x}, t) \in [0, 1]$; this quantity represents the average occupancy of sites in the vicinity of position $\mathbf{x}$. The number density of particles diverges proportional to N^2 in the hydrodynamic limit.) The deterministic evolution of f is given by [64]

$$\begin{aligned} \partial_t f = & \nabla \cdot \int \sigma_f(\mathbf{x}, \theta, \theta') \left[D_E \nabla \frac{\delta \mathcal{S}}{\delta f(\mathbf{x}, \theta')} + v_0 \mathbf{e}_{\theta'} \right] d\theta' \\ & + D_O \partial_\theta \left[f(\mathbf{x}, \theta) \partial_\theta \frac{\delta \mathcal{S}}{\delta f(\mathbf{x}, \theta)} \right], \end{aligned} \quad (11)$$

where the free energy is

$$S[f] = \int d\mathbf{x} \{ [1 - \rho(\mathbf{x})] \log[1 - \rho(\mathbf{x})] + \rho(\mathbf{x}) V(x) \} + \int d\mathbf{x} d\theta f(\mathbf{x}, \theta) \log(2\pi f(\mathbf{x}, \theta)) \quad (12)$$

and the quantity

$$\sigma_f(\mathbf{x}, \theta, \theta') \equiv s(\rho(\mathbf{x})) f(\mathbf{x}, \theta) f(\mathbf{x}, \theta') + d_s(\rho(\mathbf{x})) f(\mathbf{x}, \theta) \delta(\theta - \theta') \quad (13)$$

plays the role of a mobility. We define $d_s(\rho)$ as the diffusion constant in a symmetric exclusion process at density ρ (see below), and use the shorthand

$$s(\rho) := \frac{1 - d_s(\rho) - \rho}{\rho}. \quad (14)$$

This formulation of the equation of motion allows the terms involving $\delta S / \delta f$ to be identified with equilibriumlike processes, which act to minimize S . By contrast, the term proportional to v_0 is the nonequilibrium self-propulsion force.

The potential V enters the dynamics as an extra equilibriumlike force, adding an additional contribution to the free energy (12) compared to an otherwise equivalent expression in [64]. The added potential results in an extra contribution of $V(\mathbf{x})$ to $\frac{\delta S}{\delta f(\mathbf{x}, \theta)}$; the additional term is independent of θ' , which reflects that the force from the barrier does not depend on the particle orientation.

Writing (11) explicitly, the deterministic dynamics of f is given by

$$\frac{\partial f}{\partial t} = -D_E \nabla \cdot \mathbf{J}_{\text{diff}} - v_0 \nabla \cdot \mathbf{J}_{\text{activity}} + D_O \partial_\theta^2 f, \quad (15)$$

where

$$\mathbf{J}_{\text{diff}} = -d_s(\rho) \nabla f - f \mathcal{D}(\rho) \nabla \rho - f(1 - \rho) \nabla V, \quad (16)$$

$$\mathbf{J}_{\text{activity}} = (\mathcal{D}(\rho) - 1) f \mathbf{p} + \mathbf{e}(\theta) d_s(\rho) f. \quad (17)$$

Here we take $\mathcal{D}(\rho) := [1 - d_s(\rho)] / \rho$ for ease of writing. The fact that $\int_0^{2\pi} \sigma_f(\rho, \theta, \theta') d\theta' = f(\theta)$ explains why V appears in such a simple way in Eq. (16). It was shown in [87] that the self-diffusion coefficient can be approximated as

$$d_s(\rho) = (1 - \rho) \left(1 - \alpha \rho + \frac{\alpha(2\alpha - 1)}{2\alpha + 1} \rho^2 \right), \quad (18)$$

where $\alpha = \frac{\pi}{2} - 1$. This formula for d_s is exact in both the low- and high-density limits and gives accurate approximations for all densities. We will use it below whenever an explicit form for $d_s(\rho)$ is required.

We identify two length scales associated with the particle dynamics:

$$l_p = \frac{v_0}{D_O}, \quad l_D = \sqrt{\frac{D_E}{D_O}}, \quad (19)$$

where l_p corresponds to the typical distance a particle moves due to self-propulsion before changing direction, and l_D corresponds to the typical distance a particle moves due to diffusion in the reorientation time.

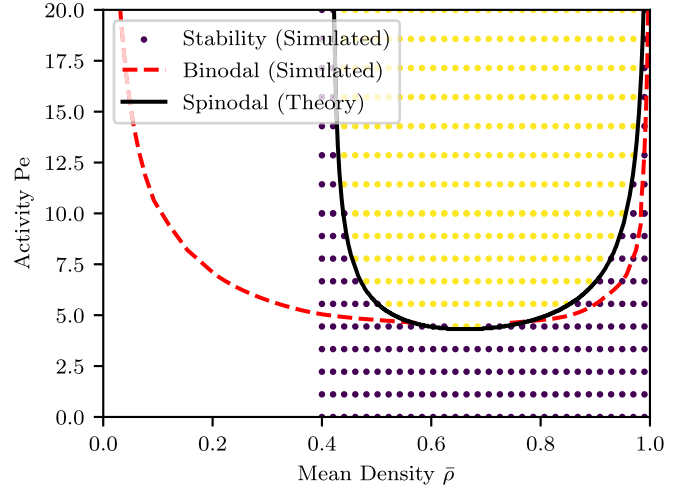


FIG. 2. Phase diagram of the active lattice gas. The binodal (coexistence curve) is obtained by solving (15)–(17) numerically, and characterizing the phase-separated steady states. The spinodal is obtained from [64], and it is compared with numerical stability results for (15)–(17) starting from weakly perturbed homogeneous states (colored points, no data shown for $\bar{\rho} < 0.4$). The critical point where the spinodal and binodal curves intersect is $(\bar{\rho}^*, \text{Pe}^*) = (0.659, 4.33)$. As is usual for active phase separation [13], these binodals (dashed curves) do not depend on the value $\bar{\rho}$, so long as it lies between them. See Appendix B 2 for further details.

It is convenient to rescale time as $t \rightarrow D_O t$ and space as $x \rightarrow \frac{x}{l_D}$, such that both are made dimensionless. The hydrodynamic equation (15) becomes

$$\frac{\partial f}{\partial t} = -\nabla \cdot \mathbf{J}_{\text{diff}} - \text{Pe} \nabla \cdot \mathbf{J}_{\text{activity}} + \partial_\theta^2 f, \quad (20)$$

where

$$\text{Pe} = \frac{l_p}{l_D} = \frac{v_0}{\sqrt{D_E D_O}} \quad (21)$$

is the Péclet number, describing the dimensionless strength of self-propulsion. We work with the rescaled space and time dynamics (20) throughout the following.

The active lattice gas exhibits MIPS, arising from the interplay between the self-propulsion of the particles and the (purely repulsive) excluded volume interactions between them [12,14,64]. The spinodal curve was computed in [64]. The binodal can be computed by simulating the hydrodynamic equations (15)–(17)—see Appendix B 2. The resulting phase diagram is shown in Fig. 2. Generically, we expect the width of liquid-vapor interfaces to scale as $\xi \sim l_D f_\xi(\text{Pe})$ and diverge as $\text{Pe} \rightarrow \text{Pe}^*(\bar{\rho})$ [64].

2. System geometry and wall (barrier) potential

Following [27,35], we study a two-dimensional system with periodic boundary conditions, and a penetrable wall (or barrier) located at $x = 0$; see Fig. 3. In the limit of large barrier height, this becomes the slit geometry of Fig. 1(b). It is known that active particles generically adsorb to hard walls [30,33,34,36], leading to full (complete) wetting in the context of two-phase coexistence [27–29,31,35,38]. Partial wetting states have been observed in active systems, either with soft

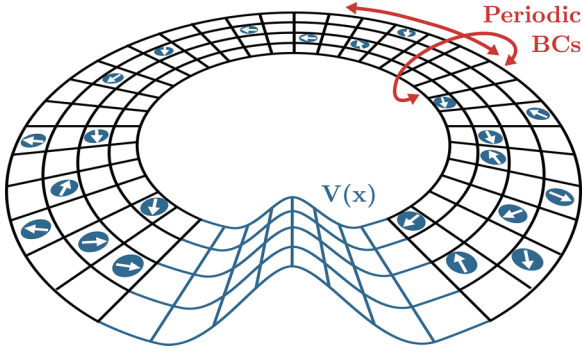


FIG. 3. Schematic illustration of the system geometry. The system occupies a two-dimensional periodic domain, with a smooth permeable barrier centered at $x = 0$. The barrier is represented by an external potential $V(x)$ of finite width w_b and height ϵ , which locally repels particles while allowing finite probability of crossing. Active particles move with self-propulsion v_0 along the $\mathbf{e}(\theta)$ direction and interact only via excluded volume. The orientation-dependent probability density $f(\mathbf{x}, \theta, t)$ is translationally invariant in the transverse direction(s), so the wetting behavior is fully characterized by the one-dimensional profile shown here.

walls [31] or using penetrable barriers [27,35]. We choose a smooth, penetrable barrier to access a transition between full- and partial-wetting states. Periodic boundary conditions are used to avoid introducing any additional confining boundaries. Experimental realizations of this geometry are discussed in Sec. V B. We emphasize that our study is at fixed total fluid mass and at liquid-vapor coexistence (in contrast to, e.g., [29]), and so the phenomena of prewetting and capillary condensation [42] are not relevant here.

Since the governing PDEs are deterministic and the initial condition is chosen to be translationally invariant along the direction parallel to the walls, the probability density $f(\mathbf{x}, \theta, t)$ only depends on the component x of the position $\mathbf{x}$ perpendicular to the barrier. Hence, we consider 1D density profiles $f(x, \theta, t)$ in the following, although we emphasize that the underlying lattice model is 2D and can easily be generalized to higher dimensions [63,64], see Appendix A 3. (Indeed, it is not possible to have orientational diffusion at fixed propulsion speed v_0 in one spatial dimension.) See Sec. II above for a discussion of equilibrium wetting with one non-translationally-invariant direction. We choose a system with domain size L in the x direction.

The barrier potential is taken as in [27]

$$V(x) = \begin{cases} 0 & \text{if } |x| \geq w_b \\ \epsilon \left[\cos\left(\frac{\pi x}{w_b}\right) + 1 \right] & \text{if } |x| < w_b, \end{cases} \quad (22)$$

so the strength of the barrier potential is a dimensionless parameter ϵ ; its half-width w_b is fixed throughout as $w_b = l_D/2$. The final (dimensionless) control parameter, alongside the barrier strength ϵ and the activity level Pe , is the (conserved) total density

$$\bar{\rho} \equiv \frac{1}{L} \int_0^L \rho(x) dx. \quad (23)$$

In summary, our system has four control parameters. Two intensive quantities, the Péclet number Pe and mean density

$\bar{\rho}$, control particle behavior. The geometry is controlled by the barrier height ϵ and the (reduced, dimensionless) system size $\ell_s = L/l_D$. The latter is generally chosen such that the slit width is much larger than the scale of liquid-vapor interfaces (although its exact value depends on the context—see Secs. III B and IV), which is the thermodynamic limit.

B. Details of numerical implementation

To solve the hydrodynamic equation (15), we discretize the orientation-dependent density $f(x, \theta, t)$ onto an $N_x \times N_\theta$ grid, and integrate (15)–(17) using an explicit first-order upwind update scheme. Further details are given in Appendix B, along with computational parameters. Writing the evolution equation as

$$\partial_t f = -\partial_x [M_x U_x] - \partial_\theta [M_\theta U_\theta] \quad (24)$$

allows separate treatment of the spatial and orientational dynamics; we choose mobilities $M_x = f d_s(\rho)$ and $M_\theta = f$ for numerical stability (in particular, the presence of $\log(f)$ in the resulting velocity fields strongly protects the dynamics against negative densities). A Courant-Friedrichs-Lewy (CFL) condition controls the adaptive time step and protects against advective instabilities. We verified that this method conserves the total density to high accuracy, preserves the parity symmetry of steady states, and reproduces the expected linear stability properties of the homogeneous active fluid (i.e., the spinodal matches the theoretical prediction given in [64]).

Determination of the binodal and spinodal lines. For the bulk (barrier-free) system, the spinodal is obtained from the linear stability analysis of [64]. The onset of instability is numerically verified by perturbing a homogeneous state with the maximally unstable⁵ eigenfunction and measuring the growth rate of this instability. The numerical results are in good agreement with the theoretical spinodal (see Fig. 2).

The binodal is computed numerically (a theoretical prediction is not currently available); starting from a weakly perturbed homogeneous state inside the spinodal, the system is evolved until a liquid-vapor interface forms. This profile is then embedded in a larger system with reflecting boundaries, yielding well-converged estimates of the coexisting densities ρ_v and ρ_l . These values depend only on Pe , consistent with standard MIPS phenomenology. These results are presented in Fig. 2.

Classification of wetting states. Because direct simulation to asymptotically long times is computationally costly, we classify steady states using measurements taken at a finite final time t_f . States are first distinguished as homogeneous or phase separated via the variance of the density away from $\bar{\rho}$, excluding the barrier region. For phase-separated states, wetting of the barrier is detected from the behavior of $\rho(x)$ near $x = 0$; finally, full versus partial wetting is identified using the asymmetry order parameter $\mathcal{A}$. The wetting transition line (distinct from the classification of individual states) is obtained by testing the linear stability of the fully wet state: whenever small

⁵That is the eigenfunction whose eigenvalue in the linearized dynamics is largest, even if this is negative (which we expect to be stable).

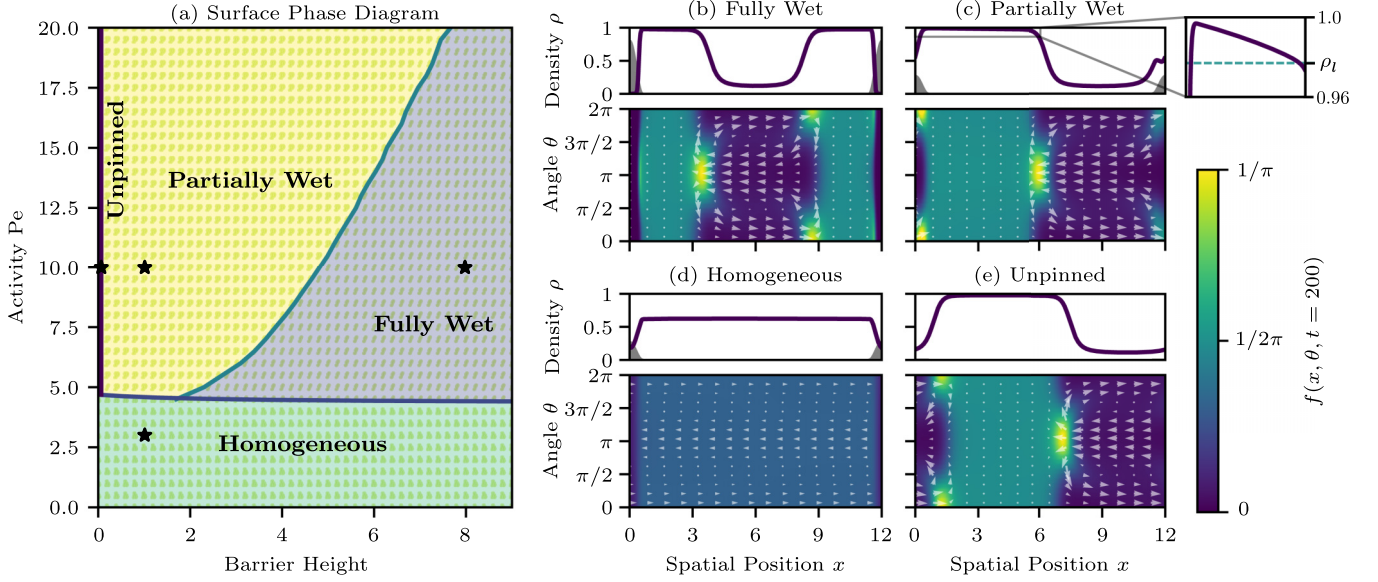


FIG. 4. Wetting of a barrier by the active lattice gas. (a) Surface phase diagram in the (ϵ, Pe) plane. Each point represents a numerical solution (simulation) of the hydrodynamic equation, whose final state is classified with the routine described in Sec. III B, which also explains the estimation of the phase boundaries. The blue line is the onset of MIPS, the green line is the wetting transition, and the dark purple line likewise represents unpinning of the liquid at $\epsilon = 0$. The structure of this surface phase diagram does not depend on $\bar{\rho}$, as long as it remains inside the spinodal. See Sec. III B for details of numerics. (b)–(e) Snapshots of the steady-state densities $f(x, \theta)$ and $\rho(x)$ for the starred points shown in the phase diagram, which illustrate the four phases. The upper plots show the density profile $\rho(x)$ and the potential $V(x)$; the lower plots show $f(x, \theta)$ (color) and the flux $\mathcal{J}$ of f (arrows scaled proportionally to $|\mathcal{J}|^{(1/4)}$). The magnified section of (c) shows a linear density gradient throughout the liquid bulk, in the partially wet state; it shows $x \in [0, 6]$.

asymmetries about the barrier grow in time, the system lies in the partially wet regime. Comparison with longer simulations confirms that this procedure reliably identifies the long-time behavior except within a narrow neighborhood of the transition lines, and that these classifications are unchanged when the simulations are run with a larger system size or using different values of the computational parameters. Further details can be found in Appendix B.

IV. RESULTS

We access the wetting transition by varying the self-propulsion Pe and the barrier height ϵ . We fix the total density at its critical value $\bar{\rho}^* = 0.659$, identified via linear stability analysis (see Sec. III B and [64]); we mostly take Pe above its critical value, in which case results depend weakly on $\bar{\rho}$, provided it remains inside the spinodal (recall Fig. 2). Repeating the analysis at different values of $\bar{\rho}$ inside the spinodal leads only to small movements of the transition lines, instead of changes to the structure of the phase diagram. We choose different values of the reduced system size ℓ_s according to the context, but these are always large enough that we recover bulk phase separation, and so the interfacial width is much smaller than the slit width.

A. Full and partial wetting

Figure 4(a) shows the surface phase diagram for a range of ϵ, Pe . To obtain this, we initialize the system in a

homogeneous state

$$f(x, \theta, 0) = \frac{\bar{\rho}^*}{2\pi} + \delta f, \quad (25)$$

where $\delta f \sim 10^{-4}$ is a small perturbation, uncorrelated in space. The system is then evolved to a final time $t_f = 200$, and states and transition lines are classified according to the routine described in Sec. III B, based on an extrapolation to the large time limit. We identify four distinct steady states, analogous to the states found in equilibrium wetting. These are labeled fully wet, partially wet, unpinned, and homogeneous. They are illustrated in Figs. 4(b)–4(e).

The fully wet state is symmetric, with liquid layers of equal width on either side of the barrier. The partially wet state has a macroscopic liquid layer on one side of the barrier; the other side is coated by a thin film of finite width ζ , see also Secs. II and IV B. (Here, the term “thin” indicates that the layer width is comparable to the interfacial width, i.e., microscopic; it is much smaller than the system size.) The unpinned (dry) state is defined by the liquid bulk detaching from the barrier entirely; in contrast to systems with noise [27], this drying transition is only observed when the barrier is removed completely, i.e., on the $\epsilon = 0$ line.⁶ This may be expected, since any repulsive barrier causes accumulation above the vapor density; it is then always favorable to have this density accumulation inside a liquid phase or near a liquid-vapor interface (such that it overlaps with the exponential tails

⁶We note that this line is labeled “unpinned” to illustrate this point, although the physics at $\epsilon = 0$ is that of standard MIPS.

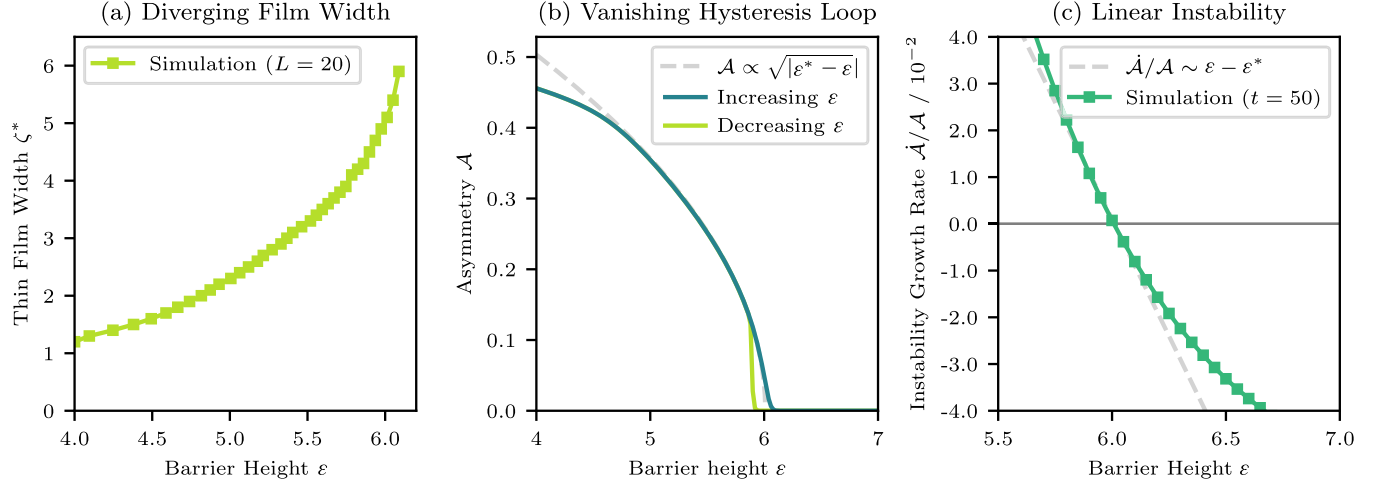


FIG. 5. The wetting transition between fully and partially wet states is critical. All plots are at $Pe = 10$. (a) The width of the liquid film on the dry wall, estimated by measuring the distance between the inflection points of ρ at the liquid-boundary and liquid-vapor interfaces. This increases as the transition is approached. Results are shown for a system of size $\ell_s = 20$, so maximum width of (symmetric) liquid layers is approximately 6.3 (using the lever rule). For an infinite system, we expect ζ to diverge at $\epsilon = \epsilon^*$. (b) Asymmetry $\mathcal{A}$ as a function of ϵ , scanning both up and down to check for hysteresis. The total scan time is $T_{\text{tot}} = 2 \times 10^5$; we use a system size $\ell_s = 5$, smaller than in (a), such that the scan is quasistatic and the behavior of $\mathcal{A}$ is not too sharp (see Appendix A). The dashed grey line is a fit to (7). We estimate $\epsilon^* \approx 6.01$. The hysteresis loop is observed to shrink as the scan time is increased. (c) Linear stability of the fully wet state with the same parameters. We plot the exponential growth rate of the asymmetry $\dot{\mathcal{A}}/\mathcal{A}$ as a function of barrier height ϵ . Positive values correspond to partial wetting (the fully wet state is unstable), and we estimate $\epsilon^* \approx 6.01$ from a linear fit. The growth rate is measured at $t = 50$, when asymmetry growth is approximately linear and $\mathcal{A}$ remains small.

of this interface). Any barrier therefore breaks translational symmetry, which is sufficient to localize the liquid slab in our noise-free dynamics. The homogeneous state occurs for $Pe < Pe^* = 4.33$, where the system does not phase separate and wetting phenomenology is not expected.

Full wetting occurs for barriers larger than the critical barrier height $\epsilon^*(Pe)$, above which the wall is effectively hard [27,29,35]. Additionally, Fig. 4(a) shows this critical barrier height $\epsilon^*(Pe)$ is *increasing* with the activity Pe . This is also unsurprising; at higher activity, particles exert a stronger driving force and can more easily overcome the barrier. The barrier thus becomes effectively softer and more permeable, favoring partial wetting (recall Sec. III A 2).

We check for finite size effects by repeating this procedure with a larger ℓ_s at isolated points on the surface phase diagram. We do not observe any structural changes to the transition lines or significant movements of their positions, and so we conclude that the surface phase diagram is well converged to its large-system limit. The boundary between full and partial wetting in Fig. 4(a) is obtained numerically by testing the linear stability of the fully wet state (see Sec. III B). We do not observe any metastability at this transition (see Sec. IV B below), so this procedure allows accurate determination of the phase boundary based on short simulations.

The steady-state results for $f(x, \theta, t)$ are consistent with expectations for active systems [25], in that the orientations of the active particles are isotropic in bulk phases but develop anisotropic structure at the liquid-vapor interface and at the barrier. One observes significant probability currents in the steady state; these are divergence-free (consistent with $\partial f / \partial t = 0$) and arise due to the particles' self-repulsion.

B. The wetting transition is critical

We now examine in detail the wetting transition, which separates fully and partially wet states. We find that this transition bears all the characteristics of an equilibrium critical wetting transition (recall Sec. II). The control parameter for active wetting is the barrier height ϵ . Figure 5 shows results for $Pe = 10$. We verified that the required qualitative features are robust across a range of $Pe > Pe^*$ and for other values of $\bar{\rho}$ inside the spinodal. The features are:

(a) *Continuous divergence of the liquid film width*. Recall from (6) that for critical wetting, the dry wall in a partially wet state carries a thin liquid film whose width ζ diverges continuously at the wetting transition. Evidence for such a divergence is shown in Fig. 5(a), which may be contrasted with a first-order wetting scenario, where ζ would have a finite limit (and finite gradient) as the transition is approached. The nature of the divergence is logarithmic in equilibrium [41,77] but we are not aware of any theoretical predictions for active systems.

(b) *Square-root scaling of the asymmetry*. The fully wet state is symmetric under reflection about $x = 0$ so $\mathcal{A} = 0$, while the partially wet state has $\mathcal{A} > 0$. Recall from (7) that for critical equilibrium wetting, the asymmetry is predicted to vanish as $(\epsilon^* - \epsilon)^{1/2}$ on approaching the transition. This behavior is illustrated in Fig. 5(b); the exponent 1/2 appears robust even out of equilibrium, as may be expected generically for pitchfork bifurcations.

(c) *Vanishing hysteresis loop*. Figure 5(b) also contrasts the behavior on increasing and reducing ϵ across the transition. There is no evidence for metastable states (neither partially nor fully wet). This is the expected behavior for critical wetting. At a first-order transition, we would expect our noiseless

hydrodynamic equations to instead support a finite hysteresis loop, even if ϵ is varied quasistatically.

(d) *Linear instability of the fully wet state.* Another characteristic feature of critical wetting, which provides further confirmation of the absence of metastability, is that the fully wet state becomes linearly unstable to density perturbations exactly at ϵ^* . This means that

$$\frac{d}{dt} \ln(A(t)) = \lambda(\epsilon^* - \epsilon) + \mathcal{O}((\epsilon^* - \epsilon)^2). \quad (26)$$

This behavior is shown numerically in Fig. 5(c): We apply a small perturbation to the fully wet state and measure the asymmetry as a function of time. As expected, we see linear instability of the fully wet state below the critical barrier height ϵ^* .

The transition between fully and partially wet states therefore bears all the hallmark characteristics of a critical wetting transition, and so we conclude that this active wetting transition is critical, at least in a two-dimensional mean-field-type model—see also [27]. In the Cahn theory of equilibrium wetting, critical wetting occurs when the effect of the barrier in preventing attractive interactions between fluid particles is significant (i.e., c_2 is large, recall Sec. II and Appendix A). In the active wetting scenario, these attractions are indeed strongly suppressed near the barrier; the large polarization $\mathbf{p}$ near the barrier means that active particles near the barrier are locally aligned. Since effective attractive interactions in MIPS arise due to collisions between oppositely aligned particles [12], the barrier inhibits the attractive fluid-fluid interactions dramatically—in the large Pe limit, we expect this inhibitive effect to become absolute. This offers a physical explanation for the observation that the wetting transition at two-phase coexistence is critical in the active system.

C. Ratchet currents in partial wetting

Although the steady states are qualitatively similar to their passive counterparts, nonequilibrium consequences of activity are also visible. In particular, we show in [76] that any asymmetrically wet barrier (i.e., $\rho(x) \neq \rho(-x)$ for $x \in [-w_b, w_b]$) must have a steady-state (uniform) current $J^{(\rho)}$ flowing through it. This ratchet current flows through the barrier from the dry (low-density) to the wet (high-density) side (i.e., in the opposite direction to diffusion).

Although ratchet currents are expected for any system of (nonmomentum-conserving) active particles in a potential that violates reflection symmetry [25,88], the current here is a result of spontaneously broken mirror symmetry about a symmetric barrier. This spontaneous ratchet current therefore lies beyond the usual ratchet principle [25]. Since this current arises from a localized driving force on the barrier, while friction acts against all particles, we expect the current to scale as $J^{(\rho)} \sim \ell_s^{-1}$, and also to vanish continuously as the (critical) wetting transition is approached. This behavior is confirmed in Fig. 6.

This ratchet current has two consequences. First, a linear density gradient is present throughout the bulk liquid,⁷ leading

⁷A linear gradient should also develop over the vapor bulk for the same reason; however, this is less evident in the numerical steady

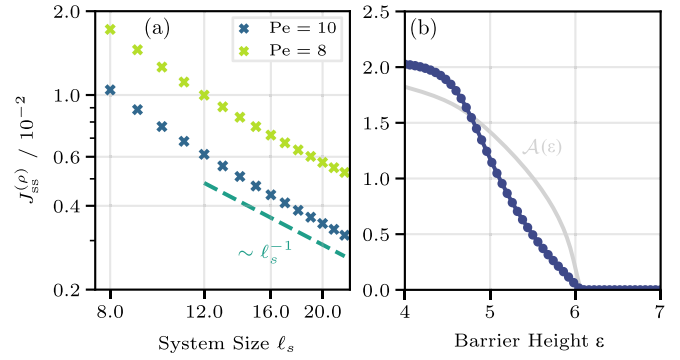


FIG. 6. Steady-state ratchet current in the partial-wetting scenario. (a) System-size dependence, asymptotically going to $1/\ell_s$. (b) Continuous vanishing of the steady-state current as the wetting transition is approached [same parameters as Fig. 5(b)—asymmetry also shown].

to an $\mathcal{O}(1)$ deviation of the density from its binodal value ρ_l that persists far from any interfaces [see Fig. 4(c), magnified section]. We verified that the density drop across the liquid is asymptotically constant as $\ell_s \rightarrow \infty$ (as expected, see discussion in [76]), and so this effect is apparent even in the limit of a large system.⁸

Second, the presence of a current through the barrier leads to a modified dynamical pathway for the full-to-partial wetting transition. Specifically, this transition involves the pinching off of one liquid layer, leading to a detached slab of liquid on one side of the barrier, which evaporates gradually until a partially wet state is reached. Such a transition would be forbidden in a passive fluid due to the creation of new interfaces; the equilibrium pathway (which proceeds due to interactions between the liquid-vapor interface and the wet barrier) is much slower. A more detailed account of this modified dynamical pathway is presented in [76], which also provides a simplified mechanistic derivation in terms of the global effects of an interfacially localized, active pump.

V. DISCUSSION

A. Summary of results

We have studied active wetting by considering the hydrodynamic behavior of an active lattice gas in the presence of a smooth permeable barrier under periodic boundary conditions. This geometry is analogous to a slit geometry widely used to study equilibrium wetting [43,68,74,75]. We take as control parameters the barrier height ϵ and self-propulsion Pe .

The surface phase diagram reveals wetting states whose density profiles are analogous to the equilibrium wetting states

states due to the longer tails of the liquid-vapor interfaces into the vapor bulk.

⁸We emphasize that this density deviation is distinct from the corrections to the density of a thin adsorbed film that arise in equilibrium wetting due to interactions between the surface-liquid and liquid-vapor interfaces. The active density shift is a bulk effect: the linear density gradient extends throughout the entire bulk layers as a nonlocal consequence of the ratchet current.

known as dry, partially wet, and fully wet. As our model is hydrodynamic (i.e., large-scale) and is studied under conditions of liquid-vapor coexistence, we conclude that these are true macroscopic wetting states, instead of microscopic accumulations of active matter or active adsorption. We observe a drying transition at $\epsilon = 0$ and a wetting transition at finite $\epsilon = \epsilon^*(\text{Pe})$. The wetting transition displays all the salient characteristics of equilibrium critical wetting [41,42,77], including a diverging film width on the dry side, vanishing hysteresis, and a pitchfork bifurcation in the asymmetry order parameter $\mathcal{A}$.

Our hydrodynamic framework also facilitates an investigation of intrinsically nonequilibrium physical phenomena in active wetting, which arise as a consequence of activity. We demonstrate two pronounced nonequilibrium effects. First, the permeable barrier acts as a ratchet, pumping mass from the dry to the wet side of the wall. This leads to a circulating current in the partially wet steady-state (recall Sec. IV A) and departure of the bulk densities from their binodal values. Second, we show that this ratchet current gives rise to a modified dynamical pathway between fully and partially wet states, which is inaccessible for passive fluids, as discussed in greater detail in [76]. Overall, these findings establish a precise correspondence between active and equilibrium wetting, while allowing us to identify subtle yet significant nonequilibrium phenomena which arise solely due to activity.

Our results complement previous findings from numerical investigations of active wetting. We use the same geometry as in [27] and [35], i.e., a slit geometry with periodic boundaries and a permeable wall; however, our deterministic equations allow us to analyze scaling in the wetting transition without finite-size effects (albeit in a model of mean-field-type). We find that wetting in two dimensions is critical, consistent with [27,29]. Finally, we emphasize that the connections we find between active and equilibrium wetting arise when the active system is in a state of two-phase coexistence. This is distinct from previous work studying accumulation and adsorption of particles on boundaries in single-phase active fluids [30,36].

B. Outlook

There are several promising directions for future work. The introduction of translational symmetry breaking in the direction(s) parallel to the wall would allow investigation of other partially wet geometries (i.e., droplet and bridge) within the same model. The presence of droplets on the barrier in the partial wetting state could result in a circulating flux pattern on the macroscopic scale, even with closed boundary conditions. Additionally, extension of this work to two dimensions would allow a comparison with the results of [31], in which the authors identified local steady-state currents around the contact line as a robust feature of active partial wetting which modifies contact angles and droplet shapes. Our deterministic hydrodynamic model would allow the theoretical predictions contained in this work to be explored in greater detail, without finite-size effects, but subject to mean-field scaling.

Another extension would be to introduce dynamical noise to the system, which can be done in our lattice-based approach in a controlled manner via a leading-order correction to the hydrodynamic equations [66,89] or by using a noisy

coarse-grained field theory such as active model B+ [19]. The inclusion of dynamical noise and fluctuations can have dramatic consequences in equilibrium [43] and nonequilibrium [27] wetting phenomenology, including changing the stability of steady-states and the order and scaling of transitions. In active systems, noise leads to a wide variety of exotic phase-separated states [13] (these are not seen here due to the hydrodynamic limit, see also [63,64]). Investigating the consequences of noise in active wetting, and the associated bubbly phase separated states, could therefore help confirm the generality and robustness of the findings presented here.

Finally, it is important to consider how the effects presented here might be observed or harnessed experimentally. A promising experimental direction could be to recreate active wetting in a slit geometry using light-activated colloidal particles [22,90,91], in which activity can be spatially confined to the central region of a slit without the need for confinement in the perpendicular direction. This approach (which allows for quasi-one-dimensional periodic boundary conditions, for example, by global confinement in a toroidal region, into which a force field of barrier type could also be introduced optically [22]) may enable direct measurement of asymmetries, currents, and wetting states, and could provide direct experimental evidence of the nonequilibrium effects described here and in [76].

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DATA AVAILABILITY

The data that support the findings of this article are openly available [92].

APPENDIX A: EQUILIBRIUM THEORY OF CRITICAL WETTING

1. Cahn theory of wetting

We briefly review some of the results of the Cahn theory of wetting, as discussed in [41,77]. These are useful points of comparison with the active wetting model, as summarized in the main text. Minimization of the free energy (3) results in a condition on the fluid density at the wall $\rho_s \equiv \rho(0)$:

$$-\frac{df_w}{d\rho_s} = \sqrt{2\kappa W(\rho_s)}. \quad (\text{A1})$$

Here, without loss of generality, we have assumed that the bulk free-energy density $W(\rho)$ has minima $W(\rho_l) = W(\rho_v) = 0$. This may give one or multiple solutions for the surface density ρ_s ; stable solutions are shown by the black points in Fig. 7. If the solution $\rho_s \in [\rho_v, \rho_l)$, the density profiles approach ρ_v as $x \rightarrow \infty$ (in the bulk) to form a partially wet or dry wall, or approach ρ_l to form a fully wet wall. If $\rho_s \geq \rho_l$, decay to the liquid density in the bulk is always energetically favorable, and the wall is always fully wet.

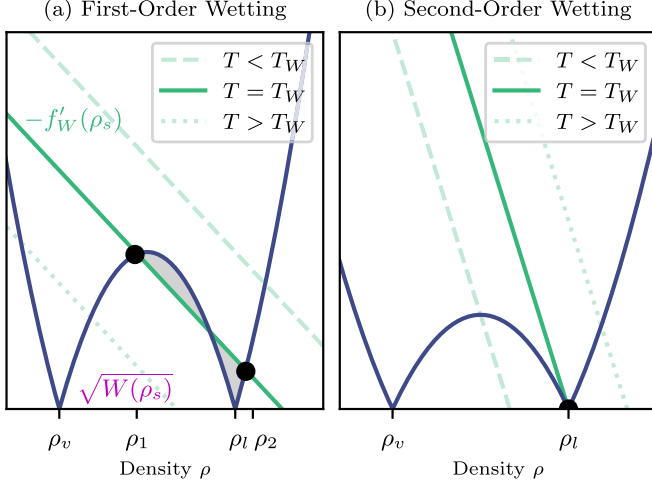


FIG. 7. Illustration of first-order and critical (second-order) wetting in the Cahn theory. (a) In a first-order wetting transition, c_2 is small, and multiple solutions for ρ_s exist; the discontinuous transition between the solutions is determined by a Maxwell construction. (b) In a critical transition, only one solution exists for ρ_s , and the wetting transition occurs when this passes continuously through ρ_l .

In a first-order wetting transition, two stable solutions exist for ρ_s , with $\rho_1 \in (\rho_v, \rho_l)$ and $\rho_2 > \rho_l$. The first solution ρ_1 corresponds to partial wetting, and the second solution ρ_2 to full wetting [see Fig. 7(a)]. The transition point between these is determined by a Maxwell construction [41]. The surface density changes discontinuously, and metastable solutions for ρ_s exist.

Critical (second-order) wetting occurs when there is only one solution for ρ_s near the wetting transition, i.e., if $c_2 > \lim_{\rho \rightarrow \rho_l} \frac{\partial}{\partial \rho} [\sqrt{2\kappa W(\rho)}]$. In this case, ρ_s passes continuously through ρ_l and no metastable states exist.

Just below a critical wetting transition, we have $\rho_s = \rho_l - \delta$. By minimizing the free energy of the density profile away from the wall, we find $\delta \propto e^{-\zeta/\xi}$, where ζ is the thickness of the wetting film and ξ is the liquid-vapor interfacial width [41]. Taking temperature as a control parameter, the free energy of this film is then given by (for a detailed discussion, see [41])

$$F_{\text{film}}(\zeta) = F_0 - \alpha(T)e^{-\zeta/\xi} + \beta(T)e^{-2\zeta/\xi}, \quad (\text{A2})$$

where $\alpha \propto (T_w - T) + \mathcal{O}((T_w - T)^2)$ and $\beta = \text{constant} + \mathcal{O}(T_w - T)$. The effect of temperature T is to control the shape of $W(\rho)$; at a higher temperature, the coexistence densities become closer and the energetic barrier between these densities becomes smaller. The optimal width of the liquid film ζ^* , which minimizes the free energy (A2), is thus given by

$$\zeta^* = -\xi \ln(T_w - T) + \text{const.} \quad (\text{A3})$$

The width of a liquid film near a critical-wetting transition therefore diverges logarithmically. In contrast, the film width at a first-order wetting transition will jump discontinuously between microscopic ($\sim \xi$) and macroscopic values. A similar analysis demonstrates that the divergence of the liquid film width is logarithmic approaching a critical transition for any

other control parameter (such as the wall attractiveness c_1 , as in the main text).

In a slit geometry, it is favorable for two initially equal liquid films of width ζ_0 to remain equal in size (so the system will remain fully wet) if and only if

$$\left. \frac{d^2 F_{\text{film}}}{d\zeta^2} \right|_{\zeta_0} \leq 0, \quad (\text{A4})$$

which, for $\zeta_0 \rightarrow \infty$, occurs precisely when $T \geq T_w$. Below this critical temperature, it is favorable for mass to redistribute from one film to another (via a translational instability) until one film is microscopic ($\sim \xi$) and the other is macroscopic; this is the partial wetting scenario. A fully wet state is therefore unstable below this critical transition temperature in a slit geometry, and a partially wet state is unstable above it. Thus, there are no metastable states in the critical wetting scenario, as mentioned in the main text. Once again, this analysis is general for any other choice of order parameter.

Asymmetry order parameter

The Cahn theory concerns the behavior of a fluid near a single, attractive boundary. We now derive the behavior of a fixed total mass of fluid in a slit geometry (i.e., in the presence of two barriers) from the Cahn theory for comparison with the active case.

Consider wetting in a slit geometry of total width $L_{\text{tot}} > 2L$, with two liquid films of thickness $L - \delta, L + \delta$ on either wall (and without loss of generality $\delta \in [0, L]$). We assume that $L \gg \xi$, such that the liquid-vapor interfaces are microscopic. The free energy due to the wetting layers becomes

$$F_{\text{film}}(\delta) = -2\alpha(T)e^{-L/\xi} \cosh(\delta/\xi) + 2\beta(T)e^{-2L/\xi} \cosh(2\delta/\xi). \quad (\text{A5})$$

To find the optimal δ , we set $\frac{\partial F_{\text{film}}}{\partial \delta} = 0$, yielding the condition

$$\cosh(\delta/\xi) = \frac{\alpha}{4\beta} e^{L/\xi}. \quad (\text{A6})$$

As the liquid film thickness varies continuously approaching a critical wetting transition, we expect δ to be small just below criticality. So, expanding $\cosh(\delta/\xi) \simeq 1 + \frac{1}{2}(\frac{\delta}{\xi})^2$:

$$\delta + \mathcal{O}(\delta^2) = \xi \sqrt{\frac{\alpha}{2\beta} e^{L/\xi} - 2}. \quad (\text{A7})$$

The asymmetry parameter can be expressed as $\mathcal{A} = \frac{L^+ - L^-}{L^+ + L^-} = \frac{\delta}{L} \in [0, 1]$ [this is equivalent to (7)]. Define $\bar{\alpha} = \alpha - 4\beta e^{(-L/\xi)}$. We may assume that this is proportional to a modified critical temperature near criticality, i.e., $\bar{\alpha} = \mu(\bar{T}_w - T) + \mathcal{O}(|\bar{T}_w - T|^2)$. The modified critical temperature accounts for a finite system size correction: $\bar{T}_w = T_w + \mathcal{O}(e^{-L/\xi})$. With these definitions, we have the following behavior for $\mathcal{A}$ near a critical wetting transition, in a finite system:

$$\mathcal{A} = \begin{cases} 0 & T \geq \bar{T}_w \\ \left[\frac{\xi}{L} \sqrt{\frac{\mu}{2\beta}} e^{\frac{L}{2\xi}} \right] (\bar{T}_w - T)^{1/2} & T < \bar{T}_w. \end{cases} \quad (\text{A8})$$

One sees that this asymmetry exhibits a pitchfork bifurcation near a critical wetting transition, with characteristic exponent

TABLE I. Summary of distinct surface phenomena in the Ising wetting problem (following Binder and Landau, [55]). Critical wetting is a two-field surface phase transition at coexistence; complete wetting is a one-field divergence within the wet phase; adsorption is a generic accumulation. The surface excess $m_s = \sum_{\text{layers } n} (m_\infty - m_n)$ measures the total deviation of the magnetization from its bulk value.

Phenomenon	Fixed	Tuned	Features
Adsorption			Finite surface excess
Critical wetting	$H = 0$	$H_1 \rightarrow H_{1c}$	$m_s \sim -\ln H_1 - H_{1c} $
Complete wetting	$H_1 < H_{1c}$	$H \rightarrow 0^+$	$m_s \sim -\ln(H/J)$
Layering	$T < T_R, H = 0$	H_1	Discrete filling

1/2. In the large-system limit $L \rightarrow \infty$, then $\mathcal{A} \rightarrow \mathcal{O}(1)$ for all $T < T_w$; this effect is signaled in Eq. (A8) by the divergence of the prefactor in this limit. In practice, this means that observation of power-law scaling in Eq. (A8) requires relatively small system sizes ($L/\xi \sim 1$), and T close to $\tilde{T}_w$.

2. Categorization of surface transitions

To contextualise the wetting and surface accumulation behaviors observed in our model, we provide a brief review of the corresponding wetting transitions and surface accumulation phenomena in equilibrium. In particular, we distinguish between critical wetting, complete wetting, adsorption, layering, and prewetting. While the active wetting literature has not always strictly maintained these distinctions, this framework is necessary for interpreting our results and comparing them with their equilibrium counterparts. For more comprehensive reviews, see, e.g., [41,55–58,80].

In equilibrium (for example, [55]), wetting may be studied using lattice models such as a semi-infinite Ising system. This model is parameterized by a bulk field H (setting the bulk chemical potential), a bulk coupling constant J , and the temperature T . The surface introduces a planar edge layer governed by a surface field H_1 (controlling the short-range wall-fluid interaction), and a modified surface coupling J_s , where the ratio J_s/J quantifies how the wall alters local fluid-fluid interactions.

The study of wetting focuses on the behavior of the interface between the phase adsorbed at the wall and the bulk phase. Depending on the thermodynamic path taken, several distinct surface phenomena emerge (summarized in Table I) [59,93]:

(1) Adsorption: A generic, finite accumulation of fluid at any attracting surface, occurring without a phase transition.

(2) Critical and First-Order Wetting Transition: Surface phase transitions that occur exactly at bulk coexistence ($H = 0$). Tuning the surface field H_1 toward a critical value can cause the initially microscopic adsorbed film thickness to diverge to a macroscopic value, either continuously (critical wetting) or discontinuously (first-order wetting).

(3) Complete wetting transition: This describes the behavior *after* the surface parameter is already wet, but before coexistence. The (unbounded) film grows continuously as the bulk field H approaches coexistence ($H \rightarrow 0^+$).

(4) Layering transitions: At temperatures below the roughening transition ($T < T_R$), the discrete lattice geometry (or packing geometry) becomes relevant. Wetting then proceeds via a series of discrete layer-by-layer filling transitions rather than continuous growth.

(5) Prewetting transition: When the wetting transition at coexistence is first-order, the discontinuity in the film thickness can extend off-coexistence into the bulk region ($H \neq 0$). This extension is known as a prewetting transition, where the film thickness jumps from thin to thick before eventually diverging at bulk coexistence. The prewetting line ultimately terminates at a surface critical point.

The topology of the surface phase diagram depends on the ratio J_s/J . Importantly, critical wetting can only occur for $J_s/J \leq 1$. For $J_s/J > 1$, a surface critical point appears at temperature T_{sc} (distinct from the bulk critical temperature T_c), associated with prewetting. (See Fig. 8 for an illustration of the surface phase diagram.)

The same categorizations also refer meaningfully to wetting in off-lattice models, and in models with long-range interactions. Although exact results are typically unavailable for models containing long-range fluid-fluid or fluid-wall interactions (e.g. Van der Waals forces), as well as fluctuations, DFT simulations still allow these cases to be studied with precision, using the same classification of surface phase transitions [43].

The parameters in our model have a clear qualitative analogy to this equilibrium wetting model. The barrier height ϵ plays the role of H_1 , while Pe dictates the effective fluid-fluid interactions, analogous to J . By fixing the global density $\bar{\rho}$ between the binodal and spinodal lines, we work strictly at bulk MIPS coexistence, analogous to $H = 0$. Furthermore, the limits taken to derive our hydrodynamic description naturally preclude discrete layering transitions. Finally, as noted in

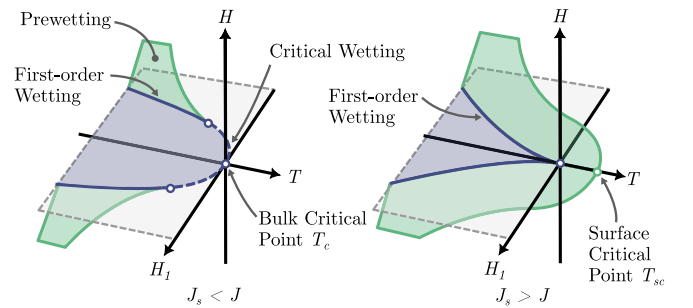


FIG. 8. Surface phase diagram for the wetting and prewetting transitions for $J_s < J$ (left), $J_s > J$ (right) after [55,80]. First-order transitions are shown with bold lines; critical transitions are shown with dashed lines; multicritical points are shown with open circles. The blue region is the fully wet area of the phase diagram, while the green region indicates prewetting (which only appears attached to a first-order wetting transition).

Sec. IV B, the orientational alignment of active particles near a wall tends to inhibit their effective attractive interactions. Consequently, our active model is analogous to the equilibrium scenario where $J_s < J$, illustrated in Fig. 8 (left). As the only microscopic interparticle interactions are short-ranged steric repulsion and the wall potential is truncated, there are no long-range forces in our model.

3. Slabs and droplets in partial wetting

In the partial-wetting regime of a slit geometry in two dimensions [i.e., $\cos(\theta_w) \in (-1, 1)$] the liquid bulk can adopt three qualitatively distinct morphologies [28,70,84]:

(1) Slab: A laterally uniform liquid bulk coating one wall (with a microscopically thin film on the other dry wall), as in Fig. 1(b.2).

(2) Droplet(s): One or more circular segments of liquid bulk on the wall [as in Fig. 1(a)], with a thin film on the remaining area of wall.

(3) Bridge: A bridge of bulk liquid connecting the two walls of the slit, whose interfaces are arcs of circles.

For thermodynamic ground states of passive systems, liquid-vapor interfaces are circular arcs and contact angles obey Young's equation. The observed morphology depends on four parameters; the slit aspect ratio L_y/L_x , two distinct interfacial tensions (where we may set $\gamma_{LV} \equiv 1$) that set the contact angle θ , and the liquid volume fraction ϕ .

The existence of some morphologies depends on geometric constraints. For example, a liquid droplet is only possible if it physically fits inside the slit. Among the existing morphologies, the ground state is the one that minimizes the free energy, which (at bulk coexistence) corresponds to minimizing the lengths of interfaces multiplied by their relevant surface tensions. (Any metastable state which is not the ground state will eventually transition to the ground state via a rare fluctuation.) Using the above considerations, a straightforward numerical calculation reveals that droplet states are preferred for small liquid volume fractions, while slab states are preferred if the aspect ratio is small, and bridge states are preferred if the slit aspect ratio is large.

Extending these arguments to active matter systems is difficult, in general. For example, the presence of steady-state currents means that droplet interface shapes and contact angles cannot be easily predicted [31]. This means that even the geometrical constraints are not straightforward to enforce. Moreover, the absence of a global free energy means that metastable and globally stable ground states cannot be distinguished in hydrodynamic theories, although stability to weak perturbations remains clear. Slab [27], droplet [35], and bridge [31] states have all been observed in active wetting systems in a slit geometry.

The one-dimensional calculations of active wetting in this paper describe only the slab states. We certainly expect that these should be the dominant states for small aspect ratios, in which droplets and bridges both tend to be suppressed by geometrical constraints. More generally, we at least expect these slab states to be metastable. We have verified this numerically for several points in the phase diagram of Fig. 4(a): we initialized partially wet slab states in a two-dimensional realization of our system and added small, spatially uncorre-

lated perturbation to f . These perturbations decay under the hydrodynamic time evolution, confirming the stability of the slab. Future work will investigate the behavior of droplet and bridge states.

APPENDIX B: NUMERICAL IMPLEMENTATION

1. Numerical timestepping routine

To simulate the dynamics of the model, as well as generate the steady states, we use an explicit time-stepping routine with one spatial and one angular dimension. We first rewrite the equations of motion in the form

$$\partial_t f = -\partial_x[M^x U^x] - \partial_\theta[M^\theta U^\theta], \quad (\text{B1})$$

where M^x, M^θ are the scalar mobilities and U^x, U^θ are the velocity fields. The mobilities are given by $M^x = f d_s(\rho)$, $M^\theta = f$, and the velocity fields by

$$U^x = -D_E \left[\partial_x(\log f) + \frac{D(\rho)}{d_s(\rho)} \partial_x \rho + \frac{(1-\rho)}{d_s(\rho)} \partial_x V \right] + v_0 \left[\cos(\theta) + \frac{p_x s(\rho)}{d_s(\rho)} \right] \quad (\text{B2})$$

and

$$U^\theta = -D_O \partial_\theta [\log f], \quad (\text{B3})$$

where $p_x \equiv \int_0^{2\pi} f(x, \theta) \cos(\theta) d\theta$ is the x component of the polarization. We then follow the same first-order update scheme as Sec. 4 of [64] and Appendix E of [67], but modified with the slightly different angular and spatial dynamics (B2) and (B3). We discretize the system into an $N_x \times N_\theta$ grid and use an upwind condition to calculate the fluxes, updating the discretized field f accordingly. A CFL condition is used for an adaptive timestep⁹ (as in [64,67])

$$\Delta t_{\max} = \min \left(\frac{\Delta x}{a \max |U^x|}, \frac{\Delta \theta}{b \max |U^\theta|}, \tau_{\max} \right). \quad (\text{B4})$$

We find that parameters $a = b = 6$, $\tau_{\max} = 10^{-4}$ are sufficient to ensure numerical stability and accuracy, while remaining practical to compute. We typically ran simulations with $N_x = 200$ and $N_\theta = 30$, although we verified that results are accurate when compared with finer grids. We always use an even N_θ to ensure the simulation is parity-symmetric.

Various checks were performed to verify that this integration scheme functions as expected; for example, the mean density is conserved up to a computational error that is typically of order 10^{-10} , the averages of higher angular moments vanish similarly in the steady state, and the stability boundary matches the theoretical prediction of [64] (see Fig. 2). Symmetry-breaking events were confirmed to be random in direction, as expected—for example, in a trial of 70 partially wet states, the liquid layer formed on the left of the barrier in 37 simulations. The dynamics of full wetting was qualitatively

⁹Although the CFL condition is necessary to prevent instabilities and inaccuracy arising from advective behavior, it does not address issues related to the density leaving its valid range of $\rho \in [0, 1]$. For this reason, $a, b > 1$ and finite $\tau_{\max}$ are necessary to ensure stability.

compared to the particle-level simulations of [38] and [27], confirming that the hydrodynamic equations reproduce the appropriate physical behavior.

2. Generation of binodal and spinodal

For the bulk fluid (without boundaries), the presence of two-phase coexistence is characterized by the spinodal (which is the boundary of linear stability of the homogeneous fluid) and the binodal (which determines the densities of the coexisting liquid and vapor). To compute the spinodal, we follow the method outlined in Sec. 3 of [64]. As the spinodal is defined in the absence of a potential, our modified equations have no effect on the linear stability analysis. For the theoretical spinodal line shown in Fig. 2, we truncate the (infinite) instability matrix of a homogeneous state to a 10×10 matrix, which is sufficient for good convergence as discussed in [64].

We confirm this theoretical prediction by finding the spinodal numerically. For this, we initialize the system with a homogeneous density field and then add a small perturbation corresponding to the eigenfunction with the largest eigenvalue (if this eigenvalue is positive, we expect instability). We then classify the phases according to the mean square deviation of the density field from the homogeneous state, using the same classification scheme as in [64]. Although the spinodal is formally defined in the limit of large system size, our simulations are run in a system size of $L_x = 20$ with a grid of 200×30 points. We find good agreement between theory and numerics, as shown in Fig. 2. No significant deviations are visible if a larger system size or smaller lattice spacing are used. The theoretical spinodal also provides a check for the validity of our numerical routine and computational parameters, as we can see that the physics of stability are preserved appropriately.

A theoretical expression for the binodal is not available, so we must rely on numerical methods to generate it. We use the following routine:

(1) Initialize a simulation from a homogeneous state plus a maximal instability perturbation as above. The parameters used are $L = 20$, $N_x = 400$, $N_\theta = 30$, $\rho = \rho^* = 0.659$, and v_0 is a specified parameter. If $v_0 > 15$, we instead use $N_x = 800$ so the (sharper) interfaces remain well-resolved.

(2) Evolve this system to time $t_0 = 200$, after which a single liquid slab has formed.

(3) Initialize a new simulation with the same parameters apart from L , but using reflecting boundary conditions (i.e., $\frac{df(x,\theta)}{dx}|_{x=0,L} = 0$), in a system of size $L = 40$. Take one-half of the previous simulation as a new initial condition, such that the profile contains vapor on one side and liquid on the other, with a single interface.

(4) Evolve to a time $t_1 = 200$. Identify the liquid and vapor densities as the maximum and minimum densities in this final state.

(5) Repeat for eight uniformly spaced values of $v_0 \in (v_{0,\text{crit}}, 6.0]$ and another ten values in the range $v_0 \in (6.0, 30.0]$.

The numerical binodal is then generated by interpolating between the discrete data points found via steps 1–5 above. The effective size over which the interface is resolved is $L = 40$, and so we expect convergence to the bulk values

given a typical interfacial scale $l_D = 1.0$ and approximately exponential tails of the liquid-vapor interface. Simulations were run for isolated points on the binodal with a larger system size and finer grid spacing; no significant departure was seen from the interpolated binodal, and so we assume that the above parameters are sufficient for good convergence for the range of Pe shown. For example, the binodals for $v_0 = 10.0$ were estimated at $(\rho_v, \rho_l) = (0.1086, 0.9855)$ and $(\rho_v, \rho_l) = (0.1115, 0.9872)$ for systems of size $L = 40$, $L = 50$, respectively, a difference of about 0.2%. Similar checks were performed for several values of v_0 , and all quoted values did not change by more than 1% when the system size was increased to $L = 50$. Therefore, we are confident that the binodal is sufficiently converged for the purposes of the figure.

3. Classifying wetting states

We now discuss numerical methods for analysis of wetting transitions in the presence of a barrier. In simulations run for a sufficiently long time, we observe four distinct phases, described in detail below. However, obtaining the long-time limiting behavior is numerically expensive, so we instead developed a numerical method for classifying what will be the long-time behavior from finite simulations up to time t_f . We verified the effectiveness of this procedure by performing long simulations at a few points on the surface phase diagram: it is generally accurate, with errors in classification only appearing very close to the transition lines (usually at a distance less than the resolution of the surface phase diagram).

We choose as model parameters $(D_E, D_O, \bar{\rho}, L) = (1, 1, \bar{\rho}_{\text{crit}}, 12)$ and computational parameters $N_x = 500$, $N_\theta = 20$, $\tau_{\text{max}} = 10^{-4}$. Repeating the analysis at single points with larger N_x, N_θ or smaller τ_{max} did not lead to any significant difference in temporal behavior or spatial structure. Our choices ensure that $L \gg l_D, w_b \gg \Delta x$, i.e., the system is much larger than the typical scale of phase interfaces or the barrier, both of which are well resolved on the grid.

As an outline of the classification algorithm: States are classified first according to whether they are phase separated, then by whether the barrier is wet, and finally by whether the barrier is fully or partially wet. Finally, the wetting transition line is calculated. We now describe each of these in turn.

a. Phase separated versus homogeneous states

The following variance of density away from its mean value is measured:

$$\sigma_\rho^2 = \int_0^{L/2-w_b} (\rho(x) - \bar{\rho})^2 dx + \int_{L/2+w_b}^L (\rho(x) - \bar{\rho})^2 dx, \quad (\text{B5})$$

where we exclude the region $x \in [L/2 - w_b, L/2 + w_b]$ to remove the effect of volume exclusion on the barrier. States are classified as homogeneous if $\sigma_\rho^2 < 0.01$. The blue phase-separation line shown in Fig. 4(a) is a smooth line fit to the data and of form $Pe_{\text{PS}} = \frac{a}{b+ce} + d$.

b. Unpinned (dry) versus wet states

As a barrier does not exist for $\epsilon = 0$, all states with $\epsilon = 0$ are trivially dry (unpinned). To find whether the states at finite ϵ are pinned to the barrier, their density profiles are examined. If the barrier is within l_D of a liquid-vapor interface or inside the liquid bulk, we consider the state to be wet. With this definition, all solutions at finite ϵ are classified as wet in the steady state. Experiments at smaller ϵ than the minimum value reported in Fig. 4(a) remain wet, indicating that any finite barrier height leads to wetting.

c. Partially wet versus fully wet states

Finally, we settle the distinction between partially and fully wet: States are classified as fully wet if $\mathcal{A} < 10^{-3}$ (i.e., the wetting films are approximately the same width); otherwise, they are classified as partially wet.

d. Full-to-partial wetting transition

To generate the wetting line shown in Fig. 4(a), we make use of the linear instability of the fully wet state. We measure the value of $\dot{\mathcal{A}}/\mathcal{A}$ at the time such that $\mathcal{A} = 10^{-3}$. A linear equation is then fit to these data points of form $\dot{\mathcal{A}}/\mathcal{A} = a(\epsilon - \epsilon^*(\text{Pe}))$, using the three smallest values of $\dot{\mathcal{A}}/\mathcal{A}$ at a given Pe (three values are taken to ensure the fit is within the linear instability regime). This gives an estimate for $\epsilon^*(\text{Pe})$, shown with the green wetting line in Fig. 4(a). The advantage of this method, compared to directly measuring the asymmetries and allocating states accordingly, is that we do not need to wait an exponentially long time for states near the transition to reach the partially wet state. This therefore allows an accurate calculation of the wetting line with a reasonable computational cost. Comparing the wetting line generated with this method (i.e., measuring instabilities) with the above classification scheme for the steady states (i.e., by measuring $\mathcal{A}$ directly), we see that only a small discrepancy is visible, as expected.

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