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Capillary interfacial tension in active phase separation

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In passive fluid-fluid phase separation, a single interfacial tension sets both the capillary fluctuations of the interface and the rate of Ostwald ripening. We show that these phenomena are governed by two different tensions in active systems, and compute the capillary tension $\sigma_{cw}$ which sets the relaxation rate of interfacial fluctuations in accordance with capillary wave theory. We discover that strong enough activity can cause negative $\sigma_{cw}$. In this regime, depending on the global composition, the system self-organizes, either into a microphase-separated state in which coalescence is highly inhibited, or into an ‘active foam’ state. Our results are obtained for Active Model B+, a minimal continuum model which, although generic, admits significant analytical progress.

Active particles extract energy from the environment and dissipate it to self-propel [1, 2]. One of their notable self-organizing features is phase separation into dense (liquid) and dilute (vapor) regions, happening even for purely repulsive particles [3–5]. Although generically a far-from-equilibrium effect, active phase separation was first described via an approximate mapping onto equilibrium liquid-vapor phase separation [3, 4]. This led to early speculation that time reversal symmetry might be restored macroscopically in steady state [3, 6–11]. Indeed, activity is an irrelevant perturbation near the liquid-vapor critical point, albeit without leading to emergent reversibility [12].

Recently it has become clear, however, that bulk phase separation in active systems displays strongly non-equilibrium features. Bubbly phase separation [13] was evidenced in simulations of repulsive self-propelled particles [14, 15]: here large liquid droplets contain a population of mesoscopic vapor bubbles that are continuously created in the bulk, coarsen, and are ejected into the exterior vapor, thereby creating a circulating phase-space current in the steady state. Microphase separation of vapor bubbles [15, 16] has been further observed numerically, whereas a similar phase of finite dense clusters is often found in experiments with self-propelled colloids [17, 18] and bacteria [19]. Recently, even more intriguing forms of phase separation have been reported in an active system composed of nematodes [20].

Much understanding of the physics of active phase separation has been gained from continuum field theories. In the simplest setting [13, 21, 22], these only retain the evolution of the density field $\phi$, while hydrodynamic [23, 24] or polar [25] fields can be added if the phenomenology requires. Their construction, via conservation laws and symmetry arguments, follows a path first introduced with Model B for passive phase separation [26–28]. Yet, these field theories differ from Model B because locally broken time-reversal symmetry implies that new non-linear terms are allowed. The ensuing minimal self-consistent theory, Active Model B+ (AMB+) [11, 13] includes all terms that break detailed balance up to order $O(\nabla^4 \phi^2)$ in a gradient expansion [11, 13] and is defined by

$$\partial_t \phi = -\nabla \cdot \left( J + \sqrt{2DM} \Lambda \right)$$

$$J/M = -\nabla \mu + \zeta (\nabla^2 \phi) \nabla \phi$$

$$\mu_\lambda [\phi] = \frac{\delta F}{\delta \phi} + \lambda |\nabla \phi|^2$$

where $F = \int dw \left[ f(\phi) + \frac{\kappa(\phi)}{2} |\nabla \phi|^2 \right]$, $f(\phi)$ is a double-well local free energy, and $\Lambda$ is a vector of zero-mean Gaussian white noises with unit variance. Model B is recovered at vanishing activity ($\lambda = \zeta = 0$), unit mobility ($M = 1$) and constant noise level $D$ [26].

It is known that at low activity (small $\lambda, \zeta$), AMB+ undergoes conventional bulk phase separation. At higher activity, Ostwald ripening [29], the classical diffusive pathway to macroscopic phase separation, can be thrown into reverse [13]. This explains the emergence of bubbly phase separation and microphase-separated vapor bubbles. (These phases arise when $\zeta, \lambda > 0$; for $\zeta, \lambda < 0$ the identities of liquid and vapor phases are interchanged.) More specific mechanisms, due to hydrodynamics [24, 30] or chemotaxis [31, 32], have also been proposed to piecewise explain some of these phases. AMB+ does not refute such specific explanations but offers a minimal framework within which to address generic features of active phase equilibria. Its simplicity admits both significant analytical progress, and efficient numerics.

For active systems showing bulk liquid-vapor phase separation it has been debated, on the basis of numerical and analytical studies, how to define the liquid-vapor interfacial tension [33–39]. One of the key results of this Letter is to confirm that no unique definition is possible. Inspired by work on equilibrium interfaces [28], we derive an effective equation for the interface height, and calculate the capillary tension $\sigma_{cw}$ which sets the spectrum of capillary waves and the relaxation times of height fluctuations. We find $\sigma_{cw}$ differs from $\sigma$, the tension introduced
in [13] to describe the Ostwald process. Whereas $\sigma < 0$ in the reverse Ostwald regime, this does not ensure capillary instability, which instead requires $\sigma_{cw} < 0$. When this holds, depending on the global density, we find two new types of interfacial instability (Fig. 1), driven by an interfacial instability similar to that of Mullins and Sekerka [40]. We find both a microphase-separated droplet state, where coalescence among droplets is highly inhibited, and an ‘active foam’ state (Fig. 4).

As is standard [13, 26] we now set $M = 1$, assume constant $D, K$, and select $f(\phi) = a(-\phi^2/2 + \phi^4/4)$ with $a > 0$. (Our results can be extended to any double-well $f$ and any $K(\phi) > 0$.) We present results for $\zeta > 0$, meaning that reversed Ostwald ripening happens only for vapor bubbles. The corresponding results for $\zeta < 0$ follow from the invariance of our model under $(\phi, \lambda, \zeta) \rightarrow -(\phi, \lambda, \zeta)$. We denote by $\phi_1$ and $\phi_2$ the coexisting vapor and liquid densities in the mean-field limit, $D = 0$; note that $\phi_{1,2} = \pm 1$ in the passive case only. More generally they are found by changing variables from $\phi$ and $f$ to $\psi$ and $g$ (denoted ‘pseudo-variables’ [13]), which solve $K\partial^2\psi/\partial \phi^2 = (\zeta - 2\lambda)\partial\psi/\partial \phi$ and $\partial g/\partial \psi = \partial f/\partial \phi \equiv \mu$, whence $\psi = K(\exp[(\zeta - 2\lambda)\phi/K] - 1)/(\zeta - 2\lambda)$. In terms of them, the equilibrium conditions $\mu_1 = \mu_2$ and $\lambda_1 = (\mu \psi - g)_1 = (\mu \psi - g)_2$ still hold [13, 41]. All our analytic results are valid in dimensions $d \geq 2$, while our numerics were done in $d = 2$ with periodic boundary conditions and system size $L_x \times L_y$, using an efficient pseudo-spectral algorithm with Euler updating [42].

![Figure 1. Mean-field phase diagram of AMB+ for $\zeta > 0$, showing sign regimes of interfacial tensions $\sigma_{cw}$ and $\sigma_{cw}$. When $\sigma_{cw} > 0$, the interface is stable and unstable otherwise. Orange circles and blue squares respectively denote the results of direct simulations of AMB+ where the instability of the interface is or is not observed. This shows the accuracy of our analytical prediction of the critical line $\sigma_{cw} = 0$. Right: interfacial instability ($\zeta = 1.5, \lambda = 2$).](image)

We start, following [29], by deriving the effective dynamics for the interface height $\tilde{h}(x, t)$ above a $(d - 1)$ plane, with in-plane and vertical coordinates $(x, y) = r$. We assume the absence of overhangs. On a rapid time-scale, diffusion relaxes $\phi(\mathbf{r}, t)$ to a quasistatic profile that depend only on the distance to the interface which, for small amplitude, long-wavelength perturbations, can be measured along the y-direction:

$$\phi(\mathbf{r}, t) = \varphi(y - \tilde{h}(x, t)).$$ (4)

It will turn out that $\tilde{h}$ solves a non-local equation in space, and we thus work in terms of its Fourier transform $h(\mathbf{q}, t)$. We proceed by plugging (4) into (1) and inverting the Laplace operator. We multiply $\nabla^{-2} \partial_t \varphi$ by $\partial_y \psi$, integrate across the interface, Fourier transform along the $x$-direction, and expand in powers of $h$. Denoting $q = |\mathbf{q}|$, we obtain [42]

$$\partial_t h = -\frac{1}{\tau(q)} h + \chi + O(h^2)$$ (5)

$$\frac{1}{\tau(q)} = \frac{2\sigma_{cw}(q)}{A(q)}$$ (6)

where

$$\sigma_{cw}(q) = \sigma + \frac{3\zeta}{4} \int dy_1 dy_2 \frac{(y_1 - y_2) \varphi'(y_1) \varphi'(y_2)}{|y_1 - y_2|}$$ (7)

$$\sigma = K \int dy \varphi'(y) \varphi(y)$$ (8)

and $\chi$ is a Gaussian noise of zero mean and correlations $\langle \chi(q_1, t_1) \chi(q_2, t_2) \rangle = C_\chi(q_1) \delta(q_1 + q_2) \delta(t_1 - t_2)$, with

$$C_\chi(q) = 4(2\pi)^d - 1 \frac{DB(q)}{A^2(q)} q.$$ (9)

In (6,9), $A(q) \equiv \int dy_1 dy_2 \varphi'(y_1) \varphi'(y_2) \exp(-q|y_1 - y_2|)$ and $B(q) \equiv \int dy_1 dy_2 \varphi'(y_1) \varphi'(y_2) \exp(-q|y_1 - y_2|)$. Note that (5) omits nonlinear terms which are derived in [42]; these are essentially the same as studied in models of conserved surface roughening [43, 44].

The effective height equations (5-9) are the fundamental analytic results of this Letter. For wavelengths much larger than an interfacial width $\xi \sim (K/\alpha)^{1/2}$, we can replace $\sigma_{cw}(q), A(q)$ and $B(q)$ with their limiting values as $q \to 0$. These, with a slight abuse of notation, are denoted as $\sigma_{cw}, A$ and $B$. Explicitly, the resulting capillary-wave tension $\sigma_{cw}$ obeys

$$\sigma_{cw} = \sigma - \frac{3\zeta}{2} \int dy \left[ \varphi(y) - \frac{\psi_1 + \psi_2}{2} \right] \varphi^2(y)$$ (10)

where $\psi_{1,2} = \psi(\phi_{1,2})$ are the pseudo-densities at the binodals. As expected, $\sigma_{cw}$ reduces to the standard interfacial tension, $\sigma_{eq} = K \int dy \varphi'(y)$, in the equilibrium limit $\lambda, \zeta \to 0$ [45]. This governs not only the capillary fluctuation spectrum, but the Laplace pressure and the rate of Ostwald ripening [28, 29]. Switching on activity breaks this degeneracy. Indeed the tension determining the rate of Ostwald ripening of a bubble was given in [13] as $\sigma = \sigma_{cw} - \zeta \int dy [\psi - \psi(0)] \varphi^2(y)$, where $\psi(0)$ is the value of the pseudo-density at the droplet center. Therefore $\sigma \neq \sigma_{cw}$ in general.
To obtain explicit predictions from (5-9), we need to evaluate $\sigma_{cw}$, $A$ and $B$. This requires knowledge of the interfacial shape $\varphi(y)$. At equilibrium, this is well-known [45]: $\varphi_{eq}(y) = \pm \tanh(y/\xi_{eq})$ with $\xi_{eq} = \sqrt{2K/a}$ and $\sigma_{eq} = \sqrt{8Ka}/9$. (Note that $A = B = 4$ in this case.) Also, whenever $2\lambda = \zeta$ it is readily shown that $\varphi = \varphi_{eq}$ so that $\sigma_{cw} = \sigma_{eq}$, although the Ostwald tensions are $\sigma = \sigma_{eq}(1 \mp \zeta/K)$ for bubble growth (−) and liquid droplet growth (+) respectively [13]. We do not have closed-form results for $\sigma_{cw}$ at general $\lambda, \zeta$; however, a change of variable to $w(\varphi) = \varphi^2$ in the integrals defining $\sigma_{cw}, A, B$ allows use of a simple numerical procedure as first used in [21] and detailed in [42] to find the low $q$ behavior. To examine $q \neq 0$ below we instead extract the interface profile from simulations at $D = 0$.

Fig. 1 shows a phase diagram of AMB+ for $\zeta > 0$ at mean-field level, delineating the zones of negative $\sigma$ and $\sigma_{cw}$. (There are none at negative $\lambda$.) For small activity, or for $\lambda \zeta < 0$, $\sigma_{cw} > 0$, even where $\sigma < 0$; here vapor bubbles undergoing reversed Ostwald ripening have stable interfaces. But at high activity a new regime emerges where $\sigma_{cw} < 0$ implying that such interfaces (and also flat ones) become locally unstable.

We first consider the regime where $\sigma_{cw} > 0$, where our theory predicts this capillary tension to govern, via (6), the relaxation times of the interface $\tau(q)$. To check this, we performed simulations of AMB+ for $D = 0$ starting from a phase separated state with the interface perturbed via a simple mode; Fig. 2 confirms that $h(q,t) = h(q,0) \exp(-t/\tau(q))$ as predicted by (5-9), for either sign of the Ostwald tension $\sigma$. Our theory also predicts the stationary structure factor of the interface

$$S(q) = \lim_{t \to \infty} \langle |h(q,t)|^2 \rangle:$$

$$S(q) = \frac{(2\pi)^{d-1} D B(q)}{\sigma_{cw}(q) q^2 A(q)} \rightarrow_{q \xi^{-1} \ll 1} \frac{(2\pi)^{d-1} D_{eff}}{\sigma_{cw} q^2} \quad (11)$$

where $D_{eff} = D(\psi_2 - \psi_1)/(\phi_2 - \phi_1)$ is an effective capillary temperature. This result generalizes capillary wave theory. Its equilibrium analog, $S(q) \propto D/\sigma_{eq} q^2$ [46], is often justified using equipartition arguments but, even in equilibrium, higher order gradient terms give sub-leading corrections at finite $q$ [47, 48]. Hence the hallmark of activity is not the similar corrections entering in (11); instead activity impacts the interfacial fluctuations by renormalizing the temperature $D \to D_{eff}$ and, separately, replacing $\sigma_{eq}$ with $\sigma_{cw}$. Even though (11) also neglects the additional nonlinearities omitted from (5), it is quite accurate at small $D$ (Fig. 2). The use of capillary wave theory in phase-separated active systems was previously advocated heuristically [34, 36, 37] but until now, only qualitative estimates were provided for the coefficient $D_{eff}/\sigma_{cw}$ in (11).

We now turn to study the region where $\sigma_{cw} < 0$. Here a drastically new non-equilibrium phenomenology arises. Although the vapor–liquid interface is unstable to height fluctuations, we will see that the system remains phase separated. For, unlike in equilibrium where demixing itself cannot be sustained at negative tension, the active interface does not undergo diffusive collapse but remains stable against normal perturbations $\phi(x,y) = \varphi(y) + \partial_x w(y)$. To see this, we linearize the mean-field dynamics around $\varphi$ to obtain $\partial_t w = \mathcal{L}w$, where $\mathcal{L} = \partial_y [f''(\varphi) - K\partial_y^2 + (2\lambda - \zeta)\varphi \partial_y] \partial_y$ and compute the spectrum of $\mathcal{L}$. In the equilibrium case, and thus also when $\zeta = 2\lambda$, this is continuous for an infinite system and it touches 0 [49, 50], resulting in algebraic decay of $w$ in time. Extending this analysis to general $\zeta, \lambda$ lies beyond our scope. Instead we compute numerically the spectrum of $\mathcal{L}$ [42], concluding that the flat interface remains stable to normal perturbations irrespective of the
Figure 4. (Left) phase diagram when $\sigma_{cw} < 0$ as a function of the global density $\phi_0 = -1, -0.4, 0.4, 1.2$ at $D = 0.05, L_x = 256, L_y = 512$ and $\lambda = 1.75, \zeta = 2$, for which $\phi_1 = -0.9, \phi_2 = 1.08$. At high and low $\phi_0$, the system is homogeneous (liquid or vapor states). Within the binodals, when the liquid is the majority phase, the system shows microphase-separated vapor bubbles whose coalescence is highly inhibited. At lower $\phi_0$, the system forms a continuously evolving active foam state. (Middle and Right): area distribution of vapor regions for the active foam state ($\phi_0 = -0.4$) and in the microphase-separated state for the noise values in the legend ($\phi_0 = 0.2$).

The stable case, $\sigma_{cw} > 0$, the stationary states seen by varying the global density $\phi_0 = f \phi d V$ are reported in Fig. 4 and Movie 2. When $\phi_0$ lies outside the mean-field binodals $\phi_{1,2}$, the system remains homogeneous. Within them, at large $\phi_0$ such that the liquid is the majority phase, we find a microphase-separated state where coalescence of crowded bubbles is highly inhibited. The bubble size distribution $P(A)$ is strongly peaked, the more peaked the lower the noise, suggesting that the average bubble size $\langle A \rangle$ is finite when $D \to 0$. Our results are converged in time for $D > 0.1$; at lower noise the system gets trapped into metastable states, evolving only because of rare fluctuations of the bubbles interface [42]. Understanding analytically the stability of circular bubbles would be key to compute $\langle A \rangle$ and $P(A)$. Clearly, though, the average size is not set by the most unstable mode of the flat interface, as the steady state is attained by a series of secondary instabilities (Movie 1 and 3). This phenomenology is at odds with the bubble phase at $\sigma_{cw} > 0$ [13], where a dynamical balance between nucleation, coalescence and reversed Ostwald causes $\langle A \rangle \to \infty$ when $D \to 0$. The difference between these two microphase separated states is also apparent in the dynamical evolution starting from bulk phase separation (Movie 3).

When the liquid is the minority phase, bubbles cannot avoid touching and coalescing. One might expect that the system attains a micro-phase separated state of liquid droplets (for $\zeta > 0$); this is not the case because, as is clear from our mechanistic argument above, the interface bends toward the vapor side. Instead, we find a previously unknown form of phase separation, which we call the ‘active foam’ state. Thin filaments of liquid are dispersed in the vapor phase, which continuously break up and reconnect. These filaments are bent on the most unstable length-scale of the flat interface, as if there were an underlying ‘virtual packing’ of polyhedral bubbles separated by thin films – but with many of these films absent. The area distribution of the vapor regions (Fig. 4b) is now peaked at at size that corresponds to the merging of two bubbles, but a power-law tail $A^{-2}$ emerges, only
cut off by the system-size. Note that the boundaries in $\phi_0$ between the different phases of Fig. 4 are qualitative: while the vapor density is almost independent of $\phi_0$ the liquid density varies markedly [42].

The techniques introduced here could be used to elucidate the capillary tension in particle-based active models, by applying them to various field-theoretical descriptions obtained by explicit coarse-graining [13, 41, 51]. The roughening properties of the interface also merit further study: the anomalous scaling found in particle-based simulations was interpreted to be in the Edwards-Wilkinson universality class [36, 37]. Our linear theory, instead, gives the critical exponents $\sigma \approx 3$ and $\chi \approx (z-d)/2$, but the impact of non-linearities should be studied by renormalisation.

Finally, although the reversal of the “Ostwald tension” $\sigma$ was previously understood [13], it is remarkable that (a) the capillary tension can likewise become negative, and that (b) this leads to new types of phase separation including active foam states. Our approach is based on AMB+, whose generic, leading-order form is agnostic as to the microscopic mechanisms underlying activity (and even phase separation). This means that the microscopic ingredients needed for our new phases remain to be identified. For the exact same reason, we expect them to be widely present, not only in motility-induced phase separation [4], but in other phase-separating systems with locally broken detailed balance.

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[42] See Supplemental Material at [URL will be inserted by publisher].
Obtain the leading non-linear terms, of order $\mathcal{O}(h^2)$, and consider the effect of noise in Appendix B2. We also note that we can approximate $\theta(x, \theta) > \phi_{th}$, where $\phi_{th} = (\phi_1 + \phi_2)/2$. We then place the interface at

$$h(x) = y + \frac{\phi(x, y) - \phi_{th}}{\phi(x, y) - \phi(x, y + \Delta x)}.$$  

Different values of $\phi_{th}$ do not change the results as far as $\phi_{th}$ is sufficiently far from the binodals.

2. Measure of the bubble-size distribution

To measure the distribution of bubbles area, we first solved the density field in a binary matrix using the threshold $\phi_{th}$. Then we applied a breadth-first search algorithm; the outcome is an $L_x \times L_y$ matrix where each pixel is labeled accordingly to the cluster it belongs to. Building the probability distribution function (PDF) of the vapor regions is then straightforward. The analysis was performed on a smoothened density field obtained by running a few time-steps of the dynamics without noise.

Appendix B: Effective interface equation in the mean-field approximation

In this Appendix we detail the derivation of the effective interface equation (5) of the main text. We start from the mean-field problem ($D = 0$) in Appendix B1 and consider the effect of noise in Appendix B2. We also obtain the leading non-linear terms, of order $\mathcal{O}(h^2)$, that correct eq. (5).

The full, non-linear, effective equation for $h$ that we obtain is

$$\partial_t h = -\frac{2\sigma_w(q)q^2}{A(q)}h + \frac{C(q)}{A(q)} \left[ -q^2 F \left[ \nabla_x \hat{h}^2 \right] \right. + \mathcal{F} \left[ \nabla_x \cdot \left( \nabla^2_x \hat{h} \nabla_x \hat{h} \right) \right] + \chi + \mathcal{O}(\xi q^3 h^3) \quad \text{(B1)}$$

where $q = |q|$, $\mathcal{F}[\cdot] = \int d\mathbf{x} e^{-i\mathbf{q} \cdot \mathbf{x}}$, is the Fourier transform operator along the $x$ direction (the same convention for the Fourier transform is used throughout what follows), and

$$C(q) = \int dy_1 dy_2 \psi(y_1)\varphi^2(y_2) e^{-q|y_1 - y_2|}.$$  

The other quantities appearing in (B1) were defined in the main text and are reported here for convenience:

$$\sigma_w(q) = \sigma_{\lambda} + \frac{3\zeta}{4} \int dy_1 dy_2 \frac{(y_1 - y_2)}{|y_1 - y_2|} \psi(y_1) \varphi^2(y_2) \quad \text{(B3)}$$

$$A(q) = \int dy_1 dy_2 \psi(y_1) \varphi(y_2) \exp(-q|y_1 - y_2|) \quad \text{(B4)}$$

$$B(q) = \int dy_1 dy_2 \psi(y_1) \varphi(y_2) \exp(-q|y_1 - y_2|).$$  

Finally, the noise $\chi$ is Gaussian and has correlations $\langle \chi(q_1, t_1) \chi(q_2, t_2) \rangle = C_\chi(q_1) \delta(q_1 + q_2) \delta(t_1 - t_2)$, where

$$C_\chi(q) = 4(2\pi)^{d-1} DB(q) A^2(q) q.$$  

1. Effective interface equation for $D=0$

Assuming

$$\phi(r, t) = \varphi(\mathbf{y} - \hat{h}(\mathbf{x}, t)) \quad \text{(B7)}$$

in the AMB+ dynamics, we obtain

$$\nabla^2 \partial_t \varphi = f'(\varphi) - K \nabla^2 \varphi + \lambda |\nabla \varphi|^2 - \zeta \nabla^2 \nabla \cdot [\nabla^2 \varphi] \quad \text{(B8)}$$

where $\hat{g}(\mathbf{x}, y) = \nabla^2 \hat{g}(\mathbf{x}, y)$ means that $\hat{g}$ solves $\nabla^2 \hat{g} = \hat{s}$. It is easy to show that the Fourier transform of $\hat{g}$ along $\mathbf{x}$ is given by

$$g(q, y) = -\frac{1}{2q} \int dy_1 e^{-q|y - y_1| s(q, y_1)}.$$  

The leading non-linear terms, of order $\mathcal{O}(h^2)$, that correct eq. (5).
Let us first consider the equilibrium case $\lambda = \zeta = 0$, hence generalizing the approach of [1] to arbitrary $q$-values. Applying the chain rule to (B8) gives

$$-
abla^2 \left[ \phi'(u) \partial_t \hat{h} \right] = f'(-\phi) - K \phi''(1 + |\nabla_x \hat{h}|^2) + K \phi' \nabla_x^2 \hat{h} \quad \text{(B10)}$$

where $\nabla_x$ is the gradient with respect to $x$. We then multiply by $\phi'$ and integrate over $u = y - \hat{h}(x,t)$ across the interface to get

$$- \int du \phi'(u) \nabla^2 \left[ \phi'(u) \partial_t \hat{h} \right] = \Delta f + \sigma_{eq} \nabla_x^2 \hat{h} \quad \text{(B11)}$$

where $\Delta f = f(\phi_2) - f(\phi_1)$, we have assumed that $\phi(y \rightarrow \infty) = \phi_2$ and $\phi(y \rightarrow -\infty) = \phi_1$, and that $\phi'$ vanishes in the bulk. Fourier transforming along the $x$ direction and using (B9) gives

$$\partial_t \hat{h} = -\frac{2\sigma_{eq} q^3}{A_{eq}(q)} \hat{h} \quad \text{(B12)}$$

where $A_{eq}(q) = \int dy_1 dy_2 \phi'(y_1)\phi'(y_2) \exp(-q|y_1 - y_2|)$. Observe that the term coming from $\Delta f$ in (B11) is proportional to $q^2 \delta(q)$ and thus vanishes. Eq. (B12) is the deterministic part of the effective interface equation for Model B.

We now consider $\lambda, \zeta \neq 0$. From (B8), the analog to (B10) now reads

$$- \nabla^2 \left[ \phi'(u) \partial_t \hat{h} \right] = \mu_\lambda + \mu_\zeta \quad \text{(B13)}$$

where

$$\mu_\lambda = f'(-\phi) + (1 + |\nabla_x \hat{h}|^2) (\lambda \phi'' - K \phi''') + K \phi' \nabla_x^2 \hat{h}$$

$$\mu_\zeta = -\zeta \nabla^2 \left[ \nabla_x \cdot (\phi'' |\nabla_x \hat{h}|^2 - \phi' \nabla_x^2 \hat{h} + \phi''') \right] \left[ \nabla_x \cdot (\phi'' |\nabla_x \hat{h}|^2 - \phi' \nabla_x^2 \hat{h} + \phi''') \right]$$

In order to progress we need to introduce the pseudo-variables [2] $\psi, g$ as the solution to

$$\frac{\partial^2 \psi}{\partial \phi^2} = \zeta - 2\lambda \frac{\partial \psi}{\partial \phi} \quad \text{and} \quad \frac{\partial g}{\partial \psi} = \frac{\partial f}{\partial \phi} \quad \text{(B15)}$$

such that $\psi \rightarrow \phi$ and $g(\phi) \rightarrow f(\phi)$ in the passive limit ($\lambda \rightarrow 0$ and $\zeta \rightarrow 0$). These quantities play the same technical role as the one played by the density $\phi$ and by the local free energy $f$ in equilibrium systems for computing the binodals [2, 3] or for computing the surface tension $\sigma$ [2], although they lack the analogous physical interpretation.

We then multiply (B13) by $\psi'$, integrate across the interface and apply the Fourier transform along $x$. For the left hand side of (B13) we obtain

$$\frac{A(q)}{2q} \partial_t h(q,t) \quad \text{(B16)}$$

where $A(q)$ is given in the main text and in (B4). Concerning the right hand side of (B13), the first term in $\mu_\lambda$ becomes

$$\delta(q) \int_{-\infty}^{\infty} du \psi'(u) f'(\phi) = g(\psi_2) - g(\psi_1) = 0$$

where we used the definition of $g$. To evaluate the second term in $\mu_\lambda$ we observe that

$$\int_{-\infty}^{\infty} du \psi'(u) \psi'(u) \psi'(u) = \frac{\zeta}{2} \int_{-\infty}^{\infty} du \psi'^2(u) \psi'(u) \quad \text{(B17)}$$

where we have used (B15). The contribution in (B17) will be cancelled by an opposite one coming from $\mu_\zeta$. The third term in $\mu_\lambda$ gives

$$-q^2 \sigma_\lambda h(q,t) \quad \text{(B18)}$$

where $\sigma_\lambda$ is defined in (B3).

We now consider $\mu_\zeta$ in (B14). Expanding in powers of $\hat{h}$, we have

$$\nabla_x \cdot (\frac{1}{\zeta} \mu_\zeta) = -\frac{1}{2} \partial_y^2 \phi'(\phi) + \frac{1}{2} \partial_y^2 \phi'(\phi) \nabla_x^2 \hat{h} \quad \text{(B19)}$$

$$- \phi'' \nabla_x \cdot [\nabla_x^2 \hat{h} \nabla_x \hat{h}] - \partial_y^2 \phi'(\phi) \nabla_x^2 \hat{h}^2 + O(\zeta q^3 h^3)$$

We then invert the Laplacian using (B9) and use

$$\partial_u e^{-\frac{q}{|y-u|}} = q \sgn(y-u) e^{-\frac{|y-u|}{q}} \quad \text{(B20)}$$

where $\sgn$ is the sign function. Applying the same procedure as before and adding up the result with (B16), (B17), (B18) we obtain the deterministic part of (B1).

2. Effect of $D \neq 0$

We now consider $D \neq 0$. Our goal is twofold: first, we derive the noise $\chi$ that enters in the effective interface equation (B1); second, we show that the Ito term that might correct (B1) actually vanishes. We start from this latter point by rewriting (B7) as

$$\dot{\phi}(r,t) = \phi \left( y - \frac{1}{(2\pi)^{d-1}} \int dq \ h(q,t)e^{iq\cdot x} \right) \quad \text{(B21)}$$

The time derivative of $\phi$ gives

$$\partial_t \phi = -\phi' \partial_t \hat{h} + \frac{1}{(2\pi)^{d-1}} \phi'' \int dq_1 dq_2 e^{(iq_1 + q_2) \cdot x} 2Dq_1 \frac{B(q_1)}{A^2(q_1)} \delta(q_1 + q_2) \quad \text{(B22)}$$

and hence

$$\partial_t \phi = -\phi' \partial_t \hat{h} + \frac{2D}{(2\pi)^{2(d-1)}} \phi'' \int dq \ q B(q) A^2(q) \quad \text{(B23)}$$

where last term in (B23) is the Ito contribution. It is then easy to show that this term gives a contribution proportional to $q\delta(q)$ in $\partial_t \hat{h}$, and hence vanishes.
where it equals the equilibrium case. For generic values of \( \lambda, \zeta \), however, \( w(\phi) = \varphi^2 \) can be obtained with a very simple numerical procedure using a technique introduced in [4] and then applied to AMB+ in [2]. We report it here for completeness.

First, as shown in [2], the binodals are easily obtained numerically solving \( \mu = f'(\phi_1) = f'(\phi_2) \) and \( \mu \psi_1 - g(\phi_1) = \mu \psi_2 - g(\phi_2) \). It is then easy to show that \( w \) solves

\[
K w' = (2\lambda - \zeta)w + 2(f' - \mu)
\]

which is solved by

\[
w(x) = e^{-\frac{\zeta - 2\lambda}{K}x} \left[ c + \frac{2}{K} \int_1^x e^{\frac{\zeta - 2\lambda}{K}y} (f'(x) - \mu) dy \right]
\]

The knowledge of \( \phi_{1,2} \) allows to fix the integration constant \( c \) and the numerical evaluation of \( w \) via (C2) is straightforward.

**Appendix D: Stability against normal perturbations**

We study here the linear stability to normal perturbations of the flat interface \( \varphi \) at mean-field level \( (D = 0) \). In this Appendix, for simplicity, we restrict to the case of one-dimensional interfaces. Due to mass conservation, the perturbed interface can be written as

\[
\phi(x, y, t) = \varphi(y) + \partial_y \epsilon(y, t)
\]

where \( \varphi \) solves

\[
\partial_y^2 \left[ f''(\varphi) - K \partial_y^2 \varphi + (\lambda - \zeta/2) \varphi^2 \right] = 0.
\]

Hence \( \epsilon \) satisfies

\[
\partial_t \epsilon = \mathcal{L} \epsilon + O(\epsilon^2)
\]

where the linear operator \( \mathcal{L} \) is

\[
\mathcal{L} = \partial_y \left[ f''(\varphi) - K \partial_y^2 + (2\lambda - \zeta/2) \varphi \partial_y \right] \partial_y.
\]

We are thus led to study the spectrum of \( \mathcal{L} \). In the equilibrium case, and thus also when \( \zeta = 2\lambda \), it was shown in [5, 6] analytically that this is continuous for an infinite system and it touches 0, resulting in algebraic decay of \( \epsilon \) in time. This result relies on the fact that \( \mathcal{L} \) is self-adjoint when \( \zeta = 2\lambda \). Extending this analysis to general \( \zeta, \lambda \) lies beyond our scope. We compute numerically the spectrum of \( \mathcal{L} \), concluding that the flat interface remains stable to normal perturbations irrespective of the sign of \( \sigma \) and \( \sigma_{cw} \). Again the spectrum of \( \mathcal{L} \) approaches 0 in the large-system limit, indicating algebraic decay of \( \epsilon \).

For generic \( \lambda \) and \( \zeta \), we studied numerically the spectrum of \( \mathcal{L} \) for finite systems. By Fourier transforming along \( y \), and for the choice of a double well local free energy \( f \), we consider the kernel \( \mathcal{L}(q_1, q_2) \) of \( \mathcal{L} \) defined

\[
\mathcal{L}(q_1, q_2) = \int \frac{dq}{\mathcal{A}(q)} \int_0^L dx \int_0^L dy \left( \int du \psi'(u) \xi(x, u + h(x, t), t) \right) \delta(q + q_1 - q_2)
\]

which is also Gaussian. It is now straightforward to show that the correlation of \( \chi \) is given by (B6). This concludes the derivation of the effective interface equation (B1).

**Appendix C: Profile \( \varphi \) of the flat interface**

As reported in the main text, the flat interfacial profile \( \varphi(y) \) can be found analytically only for the case \( 2\lambda = \zeta \),
from the relation \((\mathcal{L}c)(q_1) = \int dq_2\mathcal{L}(q_1,q_2)c(q_2)\) for any test function \(c\). Explicitly:

\[
\mathcal{L}(q_1,q_2) = (-Kq_1^4 + Aq_2^2)\delta(q_1 - q_2) - 3Aq_1q_2\mathcal{F}_y[\varphi^2](q_1 - q_2) + (2\lambda - \zeta)q_1q_2\mathcal{F}_y[\varphi](q_1 - q_2)
\]

where \(\mathcal{F}_y[\cdot]\) denotes the Fourier transform operator along \(y\). We discretised \(\mathcal{L}(q_1,q_2)\) on a grid with discretization step \(\Delta x = 1\) and total length \(L_y\), so that \(q_i = 2\pi n_i/L_y\), \(n_i = 1, ..., N\), \(N\Delta x = L_y\). We then computed numerically the eigenvalues \(\alpha_i\) of \(\mathcal{L}\) for several values of \(\lambda, \zeta\). Some of our results are reported in Fig. 1, showing that the qualitative picture is the same as at equilibrium: the spectrum of \(\mathcal{L}\) is expected to be continuous and to touch 0 for an infinite system. It should be observed that these conclusions apply irrespectively of the sign of both \(\sigma\) and \(\sigma_{cw}\); in both cases, \(\varphi\) is stable against normal perturbations.

Appendix E: Instability against height perturbations \((\sigma_{cw} < 0)\)

In the main text, we have discussed the analogy between the instability of the flat interface taking place when \(\sigma_{cw} < 0\) and the Mullins-Sekerka instability [8]. In Fig. 2 we further support this mechanistic picture, plotting the quasi-static current close to the perturbed interface. We consider three sets of parameter values corresponding to normal Ostwald ripening \((\sigma > 0)\), reversed Ostwald ripening but stable interface \((\sigma < 0, \sigma_{cw} > 0)\), and unstable interface \((\sigma < 0, \sigma_{cw} < 0)\). The current on the vapor side is always stabilizing while it is stabilizing in the liquid side only if \(\sigma > 0\). However, \(\sigma < 0\) is not sufficient to drive the instability. For this, the destabilizing current on the liquid side needs to be stronger than the one on the vapor side. This confirms that the instability arises only when the current on the liquid side overwhelms the one on the vapor side. 

![Figure 2](image)

Figure 2. Current \(\mathbf{J}\) close to (left) the stable interface and normal Ostwald ripening \((\zeta = 2\lambda = 0.5)\), (middle) stable interface and reversed Ostwald ripening for bubbles \((\zeta = 2\lambda = 2.2)\) and (right) unstable interface \((\zeta = \lambda = 2)\). Simulations are at mean-field \((D = 0)\). The liquid is shown in dark gray and the magnitude of the current in colors. Only a small part of the system is shown. We measured \(\bar{J}_V - \bar{J}_L = 0.04\) for \(2\lambda = \zeta = 0.5\), \(\bar{J}_V - \bar{J}_L = 0.08\) for \(2\lambda = \zeta = 2.2\) and \(\bar{J}_V - \bar{J}_L = -0.02\) for \(2\lambda = \zeta = 2\). This confirms that the instability arises only when the current on the liquid side overwhelms the one on the vapor side.

![Figure 3](image)

Figure 3. (Left) Damping rate \(1/\tau(q)\) vs \(q\) at \(\zeta = 1\) crossing the stability line, located at \(\lambda \approx 1.29\). The most unstable mode (the minimum of \(1/\tau(q)\)) goes to \(q = 0\) as one approaches the critical \(\lambda\). (Right) Plot of \(\tau(q)/q^3\), showing that the change of sign of the damping rate happens at the estimated critical value of \(\lambda\).
Appendix F: Liquid and vapor densities in the microphase separated and active foam states

In Fig. 4, we report the PDF of the density as a function of the global density $\phi_0$. The vapor density is found, to a good accuracy, independent of $\phi_0$. Instead, the liquid density varies rather significantly with $\phi_0$. This is expected because of two reasons: the liquid density with which a finite size vapor bubble is in equilibrium differs from the binodal [2] and the presence of multiple droplets further change such value. Obtaining the dependence of the liquid density on $\phi_0$ is an open problem.

Appendix G: Evolution towards the microphase separated state

In Fig. 5 we plot the evolution, starting from an homogeneous state or a fully phase separated state, of the average size of bubbles and their number while converging to the microphase separated state. As shown, the convergence slows down when decreasing the noise value. This is because the initially formed bubbles are stable to small perturbations of their interface and evolution to the steady state is possible only by rare events at low noise.

Appendix H: Movies

- Movie 1 : Interfacial instability starting from a fully phase separated initial condition with noise added in the bulk. Parameters: $D = 0, \zeta = 2.25, \lambda = 1.8$. System size $L_x = L_y = 128$. Total density $\phi_0 = 0.2$, leading to a microphase separation in the

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