Diversity of phase transitions and phase co-existences in active fluids

Thibault Bertrand\textsuperscript{1, \dagger} and Chiu Fan Lee\textsuperscript{2, \dagger}

\textsuperscript{1}Department of Mathematics, Imperial College London, South Kensington Campus, London SW7 2AZ, United Kingdom
\textsuperscript{2}Department of Bioengineering, Imperial College London, South Kensington Campus, London SW7 2AZ, United Kingdom

(Dated: December 11, 2020)

Active matter is not only indispensable to our understanding of diverse biological processes, but also provides a fertile ground for discovering novel physics. Many emergent properties impossible for equilibrium systems have been demonstrated in active systems. These emergent features include motility-induced phase separation, long-ranged ordered (collective motion) phase in two dimensions, and order-disorder phase co-existences (banding and reverse-banding regimes). Here, we unify these diverse phase transitions and phase co-existences into a single formulation based on generic hydrodynamic equations for active fluids. We also reveal a novel co-moving co-existence phase and a putative novel critical point.

Active matter refers to many-body systems in which each volume element can generate its own mechanical stresses [1–3]. As the fluctuation-dissipation relation is broken at the microscopic level, active matter can be viewed as an extreme form of far-from-equilibrium systems. Given the relevance of active matter to nonequilibrium physics and biophysics, the subject area has been rapidly expanding and many approaches have been used to study this diverse class of non-equilibrium, many-body systems. Arguably, the most generic way to investigate the emergent properties of an active matter system is to first formulate a model based solely on the underlying symmetries and conservation laws of the system [4].

This is what was done in the case of active fluids – a class of active matter in which translation invariance holds – in the seminal work by Toner and Tu [5–8]. Motivated by the simulation study by Vicsek et al. [9], Toner and Tu provided the generic equations of motion (EOM) for polar active fluids and demonstrated the existence of the polar ordered, or collective motion, phase using a renormalization group analysis. Subsequently, a co-existence regime consisting of the ordered and disordered phases was also found, which generically separates the disordered phase and the ordered phase in typical polar active fluid models [10–15]. Numerous studies have also confirmed the Toner-Tu EOM for polar active fluids using formal coarse-graining strategies that link microscopic models of self-propelled particles and hydrodynamic level equations [16–22].

Concurrently, dense collections of active particles without aligning interactions were shown to spontaneously phase separate in the presence of purely steric interactions. This phenomenon is now known as motility-induced phase separation (MIPS) [23]; it was first predicted theoretically and simulationally [24–26], and then experimentally observed [27, 28]. Scalar field theories, typically based on the density field of the particles, have also been formulated to describe this process, e.g., the so-called Active Model B [29, 30]. In terms of symmetries and conservation laws, MIPS and polar active fluids are completely identical. It is therefore natural to view the emergence of the ordered phase and MIPS as properties of the same class of active systems. Through scattered efforts, recent studies have attempted to gain insight into the competition between Vicsek-like aligning interactions and steric repulsion in experiments [31] and in models [32–37] of active particles. Nevertheless, our understanding of the connections between the emergence of collective motion and phase separation is still crucially lacking. In this Letter, we elucidate the interplay between these phenomena; specifically, we unify these diverse types of phases and phase co-existences in a single formulation based on generic hydrodynamic EOM for active fluids. In the process, we also uncover a new co-existence regime and a putative novel critical point.

Conservation law and symmetries.—Our model EOM are based on the conservation law and symmetries in the system. Specifically, mass conservation leads to the continuity equation:

\[ \partial_t \rho + \nabla \cdot \mathbf{J} = 0 , \]

where the total flux \( \mathbf{J} = \mathbf{p} - \eta \nabla \rho \) is composed of an active flux \( \mathbf{p} \) and a gradient term (leading to a diffusive term in the continuity equation).

For the EOM of the active flux \( \mathbf{p} \), following [5–8], we impose temporal, translation, rotation, and chiral invariances to obtain:

\[ \partial_t \mathbf{p} + \lambda \mathbf{p} \cdot \nabla \mathbf{p} = \mu \nabla^2 \mathbf{p} - \kappa(\rho) \nabla \rho + \alpha(\rho) \mathbf{p} - \beta \rho^2 \mathbf{p} , \]

where we have only retained the terms crucial to our discussion here. We have also emphasized the density dependency of the “compressibility” coefficient \( \kappa \) and that of the “order-disorder” control parameter \( \alpha \) in the above equations.

Note that our EOM differ from the Toner-Tu EOM in our choice of hydrodynamic variables and the imposition of the diffusive term in the EOM of \( \rho \). The presence of
FIG. 1. Phase co-existence in an one-dimensional active fluid. The mass density $\rho$ (upper panels) and total flux $J$ (lower panels) profiles of the four distinct types of phase co-existences: (a) a dilute (d) and disordered (D) phase, denoted as dD-cD; (b) two possible dO-cO co-existences where condensed and dilute phases are transported in the same direction (blue) or in opposite directions (red); in both cases shown here, the condensed ordered phase is moving to the right as the magnitude of the flux in the condensed phase, $|J_c|$, is higher than that in the dilute phase $|J_d|$; (c) two examples of dD-cO co-existence, where in both cases, the condensed ordered phase is moving to the right; and (d) dO-cD co-existence, the condensed disordered phase is here moving to the left due to differential adsorption at the two interfaces. The profiles shown here correspond to the stationary states of the hydrodynamic EOM. In the lower panels, the symbols correspond to the symbols shown in Fig. 2. Movies of these cases can be found in [38].

Diversity of phase co-existence. — Phase separation occurs in systems with a conserved quantity. Here, the conserved quantity is the total mass and so a phase co-existence consists of one condensed density phase (denoted by c) and one dilute density phase (denoted by d). At the same time, the Toner-Tu model allows for two distinct spatially homogeneous phases: the disordered (D) and the ordered (O) phases, characterized by whether the non-conserved order parameter $|\langle p \rangle|$ is zero or not. We therefore generically expect four possible phase co-existences: (i) dD-cD (i.e., a dilute disordered phase co-existing with a condensed disordered phase), (ii) dO-cD, (iii) dD-cO, and (iv) dO-cO (see Fig. 1). Indeed, three out of these four co-existences have already been demonstrated: (i) corresponds to MIPS, (ii) corresponds to the banding regime, and (iii) corresponds to the recently uncovered reverse-banding regime [41–43]. To the best of our knowledge, type (iv) co-existence has never been demonstrated; here, we first predict analytically and then confirm its existence numerically (see Fig. 1) in a particular model. To that end, we will first describe how a generic phase diagram can be constructed approximately following a linear stability analysis.

Linear stability and phase separation. — In thermal phase separation, a linear stability analysis of the dynamical equation of a phase separating system (e.g., Cahn-Hilliard equations) can reveal the spinodal decomposition region of the phase diagram of the system [44]. Furthermore, a signature of phase separation is that the most unstable mode from the linear instability analysis corresponds to the $k \to 0$ mode where $k$ is the wavenumber. We will use these criteria as our guiding principles in constructing an approximate phase diagram for a particular hydrodynamic model. Specifically, we will perform a linear stability analysis on the EOM and focus on the $k \to 0$ limit. Furthermore, since in the disordered phase, the instability has no direction dependency; while in the ordered phase, the most unstable direction is longitudinal to the direction of the collective motion, the initial perturbation in our stability analysis will always be along the direction of the ordered state [18, 19].

As an example, in the disordered case ($\alpha < 0$), we expand around the homogeneous disordered state with $\rho = \rho_0 + \delta \rho \exp[st - ikx]$, $p = \delta p \exp[st - ikx]$, where we have arbitrarily chosen the $x$ direction to be the direction of interest. Expanding $\kappa$ and $\alpha$ in (2) as follows:

$$\kappa(\rho) = \sum_{i=0}^{\infty} \kappa_i \delta \rho^i \quad \text{and} \quad \alpha(\rho) = \sum_{i=0}^{\infty} \alpha_i \delta \rho^i,$$

(3)

to linear order, the EOM read:

$$\partial_t \delta \rho = -\partial_x \delta p + \eta \partial_x^2 \delta \rho,$$

(4)

$$\partial_t \delta p = \mu \partial_x^2 \delta p - \kappa_0 \partial_x \delta \rho - |\alpha_0| \delta p.$$
Solving for $s$ and focusing on the hydrodynamic limit ($k \to 0$), we have

$$s = \begin{cases} 
-|\alpha_0| + \left(\frac{\kappa_0}{|\alpha_0|} + \mu\right) k^2 + \mathcal{O}(k^3), \\
- \left(\frac{\kappa_0}{|\alpha_0|} + \eta\right) k^2 + \mathcal{O}(k^3).
\end{cases} \quad (5)$$

The first eigenvalue ($-|\alpha_0|$) corresponds to the fast relaxation when the active flux deviates from the mean field value $p_0 = 0$ in the absence of spatial variations. The second eigenvalue quantifies when the instability sets in, which happens whenever $\kappa_0 + \eta|\alpha_0|$ becomes negative. Since the system is in the disordered regime, within this instability region, the system exhibits dD-cD coexistence as shown in Fig. 1(a).

The analysis in the ordered regime ($\alpha > 0$) follows the exact same procedure; the full details of the linear stability analysis can be found in [38]. Here, we just recall that when $\kappa - \eta \alpha > 0$, the homogeneous disordered phase will generically be separated from the homogeneous ordered phase by phase co-existence regions. We will now present a particular model that illustrates the diversity of phase transitions and phase co-existences possible in active fluids.

A model with all four phase co-existences.—The linear stability analysis can be applied straightforwardly once $\kappa(\rho)$ and $\alpha(\rho)$ in (2) are explicitly defined. Here, we consider the following model:

$$\alpha(\rho) = -A + 18\rho - 10/3\rho^2 \quad (6)$$
$$\kappa(\rho) = 140 - 145\rho + 30\rho^2, \quad (7)$$

with $\eta = 2$, $\lambda = 1$, $\beta = 0.5$ and $\mu = 1$. A microscopic-level (particle based) system that realizes this model could for instance be a polar active fluid with contact inhibition of alignment (e.g., as discussed in [42]) such that its equation of state dictates that it can also phase separates within the homogeneous ordered phase.

For large enough values of $A$, $\alpha$ remains negative, and so the system remains in the disordered phase. In this case, instability occurs if $(\kappa - \eta \alpha) < 0$, and we expect dD-cD co-existence in this parameter range. On the other hand, as $A$ decreases, the range of densities $\rho$ for which the system is in the ordered phase gets wider, and, importantly, is separated from the disordered phase by two instability regions: a dD-cO co-existence to the left and a dO-cD co-existence to the right. Simultaneously, $(\kappa - \eta \alpha)$ remains negative around $\rho \sim 2.5$. We therefore expect an interesting interplay of distinct phase separations.

In Fig. 2, the instability regions resulting from our linear stability analysis are shown as the shaded area, while the homogeneous disordered (D) and ordered (O) regions are shown in white. We equate the instability region to be within the phase separating region, but to which of the four possible types of phase co-existences?

Since the conserved quantity here is the total mass, $\rho$ can be redistributed as long as the overall density remains the same. Therefore, we can characterize the phases as follows: given any starting point on the phase diagram within the instability region (shaded area in the phase diagram), we extend a horizontal line from that point; the first homogeneous phase encountered to the right (respectively, to the left) will describe the nature of the condensed (respectively, dilute) phase (D or O).

Using the above construction, we see that this particular model contains all four variations of the phase co-existences (Fig. 2). As aforementioned, we report for the first time the existence of a phase co-existence in which both the dilute and condensed phases are ordered. A priori, these two co-moving phases can move either in the same direction or in opposite directions. By directly solving the hydrodynamic EOM numerically [38], this is evidenced in the stationary profiles shown in Fig. 1(b). Besides this particular model, we note that further diversity of phase diagrams are rendered possible by varying the specific definitions of $\alpha(\rho)$ and $\kappa(\rho)$; we discuss other
A multi-critical point.—Besides uncovering the novel dO-cO co-existence regime, our analysis also reveals a potentially novel critical point. To illustrate this (Fig. 3), we consider a polar active fluid system in which there are two generic parameters $K_1$ and $K_2$ that control the phase behavior of the system. Specifically, the system undergoes dD-cD phase separation at high $K_1$ while at small $K_1$ the system is in the ordered phase. In addition, we assume that the second parameter $K_2$ controls the threshold level $K_1$ at which the distinct phase separations happen (Fig. 3). In other words, instead of having additional phase separation due to a negative $\kappa$ inside the homogeneous ordered phase as in the previous example, we have here a dD-cD phase co-existence in the homogeneous disordered phase instead.

Now, dD-cD phase separation at criticality belongs to the Ising universality class [46] (but see also [47, 48]). In terms of our hydrodynamic EOM, this critical point corresponds to having $\alpha > 0, \kappa_0 = \kappa_1 = 0$ in (3). On the other hand, the order-disorder critical point that accompanies critical dD-cO and dO-cD phase separations belongs putatively to a novel universality class ($\kappa > 0, \alpha_0 = \alpha_1 = 0$ [42]. Therefore, by fine tuning $\kappa_0, \kappa_1, \alpha_0$ and $\alpha_1$ to zero (indicated by the yellow ball in Fig. 3), these two distinct critical points coincide and the resulting multi-critical point is likely to correspond to yet a distinct universality class for the following reasons: At the linear level around this critical point, the EOM of the active flux $p$ is completely decoupled from that of the density field $\rho$. Specifically, the linear EOM are

$$\partial_t \delta \rho + \nabla \cdot p = \eta \nabla^2 \delta \rho , \quad \partial_t p = \mu \nabla^2 p + f$$

where $f$ is a Gaussian noise term with a non-zero standard deviation. The fact that $\delta \rho$ does not feature in the linear EOM of $p$ is distinct from all known active fluids at the order-disorder critical transition [42, 49–55].

Using the linear theory above, we can also identify some interesting novel features of this critical point. To do so, we first perform the following re-scalings:

$$\begin{align*}
    r &\mapsto e^\xi r , \quad t \mapsto e^{\xi t} \\
    \delta \rho &\mapsto e^{\chi_\rho } \delta \rho , \quad p \mapsto e^{\chi_p } p
\end{align*}$$

We can then conclude that the following choice of scaling exponents leave the linear EOM invariant [38]:

$$z = 2 , \quad \chi_\rho = \frac{2 - d}{2} , \quad \chi_p = \frac{4 - d}{2} .$$

Applying these linear exponents to the generic nonlinear EOM then indicates that (i) the upper critical dimension $d_c$ is 8 and (ii) the first nonlinear term that becomes irrelevant right below $d_c$ is $\nabla \delta \rho^2$ in the EOM of $p$ [38].

Summary & Outlook.—Starting from generic hydrodynamic EOM of polar active fluids, we have unified existing phase transitions and phase co-existences into a single formulation. In particular, we showed that there are interesting cases in [38].

Instability vs. phase separation.—The instability region obtained from a linear stability analysis does not correspond exactly to the whole phase separation region. Indeed, as in thermal phase separation, the instability region in fact corresponds to the spinodal decomposition region, which is always flanked by the so-called nucleation and growth regions on either side [44, 45]. This is no different here: the actual phase separation boundaries encapsulate the instability regions (Fig. 1). Of course, while the phase separation boundaries (i.e., the binodal lines) for thermal systems in equilibrium can be obtained by analyzing the free energy, e.g., by using the Maxwell tangent method, no free energy exists in our non-equilibrium systems and so the phase separation boundaries will instead be given by the appropriate boundary conditions obtained from the stable steady-state solution of the actual hydrodynamic EOM. This is exactly what we did to obtain the profiles shown in Fig. 1. Specifically, the locations of the binodal lines correspond to the density values of the stationary regions of the condensed and dilute phases (see [38] for further details).
generically four distinct types of phase separations, and illustrated them with a particular model. In doing so, we exhibited a novel co-existence regime: the co-existence of a dilute ordered phase and a condensed ordered phase. We also revealed a putative novel critical behavior. Our work highlights the richness of generic polar active fluid models. The phase behavior can be further enriched by considering variations in other parameters in the EOM. For instance, patterns other than phase separation has been observed when the coefficient \( \mu \) in the EOM of the momentum density field (2) becomes negative \[56\]. We believe that elucidating these diverse phase behaviors will be a fruitful research direction in the future.

---

1 Electronic address: b.bertrand@imperial.ac.uk
2 Electronic address: clee@imperial.ac.uk

[1] Sriram Ramaswamy, “The mechanics and statistics of active matter,” Annual Review of Condensed Matter Physics 1, 323–345 (2010).
[2] M. C. Marchetti, J. F. Joanny, S. Ramaswamy, T. B. Liverpool, J. Prost, Madan Rao, and R. Aditi Simha, “Hydrodynamics of soft active matter,” Rev. Mod. Phys. 85, 1143–1189 (2013).
[3] Clemens Bechinger, Roberto Di Leonardo, Hartmut Löwen, Charles Reichhardt, Giorgio Volpe, and Giovanni Volpe, “Active particles in complex and crowded environments,” Rev. Mod. Phys. 88, 045006 (2016).
[4] P. M. Chaikin and T. C. Lubensky, Principles of Condensed Matter Physics (Cambridge University Press, 1995).
[5] John Toner and Yuhai Tu, “Long-range order in a two-dimensional dynamical XY model: How birds fly together,” Phys. Rev. Lett. 75, 4326–4329 (1995).
[6] John Toner and Yuhai Tu, “Flocks, herds, and schools: A quantitative theory of flocking,” Phys. Rev. E 58, 4828–4858 (1998).
[7] John Toner, Yuhai Tu, and Sriram Ramaswamy, “Hydrodynamics and phases of flocks,” Annals of Physics 318, 170 – 244 (2005).
[8] John Toner, “Reanalysis of the hydrodynamic theory of fluid, polar-ordered flocks,” Phys. Rev. E 86, 031918 (2012).
[9] Tamás Vicsek, András Czirók, Éshel Ben-Jacob, Iain Cohen, and Ofer Shochet, “Novel type of phase transition in a system of self-driven particles,” Phys. Rev. Lett. 75, 1226–1229 (1995).
[10] Hugues Chaté, Francesco Ginelli, Guillaume Grégoire, and Franck Raynaud, “Collective motion of self-propelled particles interacting without cohesion,” Phys. Rev. E 77, 046113 (2008).
[11] Guillaume Grégoire and Hugues Chaté, “Onset of collective and cohesive motion,” Phys. Rev. Lett. 92, 025702 (2004).
[12] Florian Thüroff, Christoph A. Weber, and Erwin Frey, “Numerical treatment of the boltzmann equation for self-propelled particle systems,” Phys. Rev. X 4, 041030 (2014).
[13] A. P. Solon and J. Tailleur, “Revisiting the flocking transition using active spins,” Phys. Rev. Lett. 111, 078101 (2013).
[14] A. P. Solon and J. Tailleur, “Flocking with discrete symmetry: The two-dimensional active ising model,” Phys. Rev. E 92, 042119 (2015).
[15] Arvind Gopinath, Michael F. Hagan, M. Cristina Marchetti, and Aparna Baskaran, “Dynamical self-regulation in self-propelled particle flows,” Phys. Rev. E 85, 061903 (2012).
[16] Aparna Baskaran and M. Cristina Marchetti, “Hydrodynamics of self-propelled hard rods,” Phys. Rev. E 77, 011920 (2008).
[17] Aparna Baskaran and M. Cristina Marchetti, “Enhanced diffusion and ordering of self-propelled rods,” Phys. Rev. Lett. 101, 268101 (2008).
[18] Eric Bertin, Michel Droz, and Guillaume Grégoire, “Boltzmann and hydrodynamic description for self-propelled particles,” Phys. Rev. E 74, 022101 (2006).
[19] Eric Bertin, Michel Droz, and Guillaume Grégoire, “Hydrodynamic equations for self-propelled particles: microscopic derivation and stability analysis,” Journal of Physics A: Mathematical and Theoretical 42, 445001 (2009).
[20] Anton Peshkov, Igor S. Aranson, Eric Bertin, Hugues Chaté, and Francesco Ginelli, “Nonlinear field equations for aligning self-propelled rods,” Phys. Rev. Lett. 109, 268701 (2012).
[21] A. Peshkov, E. Bertin, F. Ginelli, and H. Chaté, “Boltzmann-ginzburg-landau approach for continuous descriptions of generic vicsek-like models,” The European Physical Journal Special Topics 223, 1315–1344 (2014).
[22] Eric Bertin, Aparna Baskaran, Hugues Chaté, and M. Cristina Marchetti, “Comparison between smoluchowski and boltzmann approaches for self-propelled rods,” Phys. Rev. E 92, 042141 (2015).
[23] Michael E. Cates and Julien Tailleur, “Motility-induced phase separation,” Annual Review of Condensed Matter Physics 6, 219–244 (2015).
[24] J. Tailleur and M. E. Cates, “Statistical mechanics of interacting run-and-tumble bacteria,” Phys. Rev. Lett. 100, 218103 (2008).
[25] Yaouen Fily and M. Cristina Marchetti, “Athermal phase separation of self-propelled particles with no alignment,” Phys. Rev. Lett. 108, 235702 (2012).
[26] Gabriel S. Redner, Michael F. Hagan, and Aparna Baskaran, “Structure and dynamics of a phase-separating active colloidal fluid,” Phys. Rev. Lett. 110, 055701 (2013).
[27] Jeremie Palacci, Stefano Sacanna, Asher Preska Steinberg, David J. Pine, and Paul M. Chaikin, “Living crystals of light-activated colloidal surfers,” Science 339, 936–940 (2013).
[28] Ivo Buttinoni, Julian Bialke, Felix Kümmel, Hartmut Löwen, Clemens Bechinger, and Thomas Speck, “Dynamical clustering and phase separation in suspensions of self-propelled colloidal particles,” Phys. Rev. Lett. 110, 238301 (2013).
[29] Raphael Wittkowski, Adriano Tiribocchi, Joakim Stenhammar, Rosalind J. Allen, Davide Marenduzzo, and Michael E. Cates, “Scalar φ∗ field theory for active-particle phase separation,” Nature Communications 5, 4351 (2014).
[30] Ellen Tjhung, Cesare Nardini, and Michael E. Cates, “Cluster phases and bubbly phase separation in active fluids: Reversal of the ostwald process,” Phys. Rev. X 8,
31] Marjolein N. van der Linden, Lachlan C. Alexander, Dirk G. A. L. Aarts, and Olivier Dauchot, “Interrupted motility induced phase separation in aligning active colloids,” Phys. Rev. Lett. 123, 098001 (2019).
32] Fernando Peruani, Tobias Klaus, Andreas Deutsch, and Anja Voss-Boehme, “Traffic jams, gliders, and bands in the quest for collective motion of self-propelled particles,” Phys. Rev. Lett. 106, 128101 (2011).
33] F. D. C. Farrell, M. C. Marchetti, D. Marenduzzo, and J. Tailleur, “Pattern formation in self-propelled particles with density-dependent motility,” Phys. Rev. Lett. 108, 248101 (2012).
34] Aitor Martín-Gómez, Demian Levis, Albert Díaz-Guilera, and Ignacio Pagonabarraga, “Collective motion of active brownian particles with polar alignment,” Soft Matter 14, 2610–2618 (2018).
35] E. Sesé-Sansa, I. Pagonabarraga, and D. Levis, “Velocity alignment promotes motility-induced phase separation,” EPL (Europhysics Letters) 124, 30004 (2018).
36] L. Caprini, U. Marini Bettolo Marconi, and A. Puglisi, “Spontaneous velocity alignment in motility-induced phase separation,” Phys. Rev. Lett. 124, 078001 (2020).
37] Ashreya Jayaram, Andreas Fischer, and Thomas Speck, “From scalar to polar active matter: Connecting simulations with mean-field theory,” Phys. Rev. E 101, 022602 (2020).
38] Supplemental Material.
39] Jörn Dunkel, Sebastian Heidenreich, Markus Bär, and Raymond E. Goldstein, “Minimal continuum theories of structure formation in dense active fluids,” New Journal of Physics 15, 045016 (2013).
40] Vasco M. Worlitzer, Gil Ariel, Avraham Be’er, Holger Stark, Markus Bär, and Sebastian Heidenreich, “Motility-induced clustering and meso-scale turbulence in active polar fluids,” (2020), arXiv:2011.12219 [cond-mat.soft].
41] Simon K. Schnyder, John J. Molina, Yuki Tanaka, and Ryoichi Yamamoto, “Collective motion of cells crawling on a substrate: roles of cell shape and contact inhibition,” Scientific Reports 7, 5163 (2017).
42] D. Nesbitt, G. Pruessner, and C. F. Lee, “Uncovering novel phase transitions in dense dry polar active fluids using a lattice boltzmann method,” (2019), arXiv:1902.00530v2.
43] Delphine Geyer, David Martin, Julien Tailleur, and Denis Bartolo, “Freezing a flock: Motility-induced phase separation in polar active liquids,” Phys. Rev. X 9, 031043 (2019).
44] Christoph A. Weber, David Zwicker, Frank Jülicher, and Chiu Fan Lee, “Physics of active emulsions,” Reports on Progress in Physics 82, 064601 (2019), arXiv:1806.09552.
45] Jean-Louis Barrat and Jean-Pierre Hansen, Basic Concepts for Simple and Complex Liquids (Cambridge University Press, 2003).
46] Benjamin Partridge and Chiu Fan Lee, “Critical motility-induced phase separation belongs to the ising universality class,” Phys. Rev. Lett. 123, 068002 (2019).
47] Jonathan Tammo Siebert, Florian Dittrich, Friederike Schmid, Kurt Binder, Thomas Speck, and Peter Virnau, “Critical behavior of active Brownian particles,” Physical Review E 98, 030601 (2018).
48] Fernando Caballero, Cesare Nardini, and Michael E Cates, “From bulk to microphase separation in scalar active matter: a perturbative renormalization group analysis,” Journal of Statistical Mechanics: Theory and Experiment 2018, 123208 (2018).
49] Francesco Ginelli and Hugues Chaté, “Relevance of Metric-Free Interactions in Flocking Phenomena,” Physical Review Letters 105, 168103 (2010).
50] Anton Peshkov, Sandrine Ngo, Eric Bertin, Hugues Chaté, and Francesco Ginelli, “Continuous Theory of Active Matter Systems with Metric-Free Interactions,” Physical Review Letters 109, 098101 (2012).
51] Leiming Chen, John Toner, and Chiu Fan Lee, “Critical phenomenon of the order-disorder transition in incompressible active fluids,” New Journal of Physics 17, 042002 (2015).
52] Robert Großmann, Fernando Peruani, and Markus Bär, “Superdiffusion, large-scale synchronization, and topological defects,” Physical Review E 93, 040102 (2016).
53] B. Mahault, X.-c. Jiang, E. Bertin, Y.-q. Ma, A. Patelli, X.-q. Shi, and H. Chaté, “Self-Propelled Particles with Velocity Reversals and Ferromagnetic Alignment: Active Matter Class with Second-Order Transition to Quasi-Long-Range Polar Order,” Physical Review Letters 120, 258002 (2018).
54] A. Cairoli and C. F. Lee, “Hydrodynamics of active lévy matter,” (2019), arXiv:1903.07565v3.
55] A. Cairoli and C. F. Lee, “Active lévy matter: Anomalous diffusion, hydrodynamics and linear stability,” (2019), arXiv:1904.08326v2.
56] Henricus H Wensink, Jörn Dunkel, Sebastian Heidenreich, Knut Drescher, Raymond E Goldstein, Hartmut Löwen, and Julia M Yeomans, “Meso-scale turbulence in living fluids,” Proceedings of the National Academy of Sciences 109, 14308–14313 (2012).
I. LINEAR STABILITY ANALYSIS

In this section, we proceed to the linear stability analysis of our hydrodynamic EOM in one dimension

\begin{align}
\partial_t \rho + \partial_x p &= \eta \partial_x^2 \rho \\
\partial_t p + \lambda \partial_x p &= \mu \partial_x^2 p - \kappa(\rho) \partial_x \rho + \alpha(\rho) p - \beta p^3
\end{align}

A. Instability of the homogeneous disordered state

First, we consider the stability conditions for the homogeneous disordered state, i.e. a state with zero average momentum density. We linearize Eqs. (1-2) around the homogeneous disordered state with \( \rho = \rho_0 + \delta \rho \), \( p = \delta p \), \( \alpha(\rho) = -|\alpha_0| \) and \( \kappa(\rho) = \kappa_0 \). To linear order, we obtain

\begin{align}
\partial_t \delta \rho &= -\partial_x \delta p + \eta \partial_x^2 \delta \rho \\
\partial_t \delta p &= \mu \partial_x^2 \delta p - \kappa_0 \partial_x \delta \rho - |\alpha_0| \delta p
\end{align}

with the fluctuation terms written as

\begin{align}
\delta \rho &= \delta \rho_0 \exp[st - ikx] \\
\delta p &= \delta p_0 \exp[st - ikx]
\end{align}

Reinjecting the fluctuation terms in the linearized equations, one obtains

\begin{align}
s \delta \rho_0 &= ik \delta p_0 - \eta k^2 \delta \rho_0 \\
s \delta p_0 &= -k^2 \mu \delta p_0 + ik \kappa_0 \delta \rho_0 - |\alpha_0| \delta p_0
\end{align}

which we can rewrite as the following eigenvalue problem

\[ A \delta u_0 = s \delta u_0 , \]

with \( \delta u_0 = (\delta \rho_0, \delta p_0)^\top \) and

\[ A = \begin{bmatrix} -k^2 \eta & ik \\ ik \kappa_0 & -k^2 \mu - |\alpha_0| \end{bmatrix} \]

We know that the stability conditions are given by the sign of Re\([s]\); the eigenvalues of this 2 \( \times \) 2 matrix are given by

\[ s = \frac{\text{Tr} A}{2} \pm \frac{1}{2} \sqrt{(\text{Tr} A)^2 - 4 \det A} \]

i.e.

\[ s = -\frac{|\alpha_0| + k^2(\mu + \eta)}{2} \pm \frac{1}{2} \sqrt{\left[|\alpha_0| + k^2(\mu + \eta)\right]^2 - 4 \left[|\alpha_0| \eta k^2 + \kappa_0 k^2 + \eta \mu k^4\right]} \]
We know that \( s \) is real as long as \( |\alpha_0|^2 + k^2(\mu + \eta)^2 - 4k^2\kappa_0 > 0 \iff |\alpha_0| \geq 0 \) when \( k \to 0 \); which is true. In the hydrodynamic limit \( (k \to 0) \), we can expand the square root and we obtain
\[
s = -\frac{|\alpha_0| + k^2(\mu + \eta)}{2} \pm \frac{1}{2} \left[ |\alpha_0|^2 - 2k^2(\mu + \eta)|\alpha_0| + (\mu + \eta)^2k^4 - 4|\alpha_0|\eta k^2 - 4\kappa_0k^2 - 4\eta\mu k^4 \right]^{1/2}
\] (13)
\[
\approx -\frac{|\alpha_0| + k^2(\mu + \eta)}{2} \pm \frac{1}{2} |\alpha_0| \left[ 1 - \frac{k^2(\mu + \eta)}{|\alpha_0|} - \frac{2\eta k^2}{|\alpha_0|} - \frac{2k^2\kappa_0}{|\alpha_0|^2} \right]
\] (14)
\[
\approx -\frac{1}{2} \left[ |\alpha_0| \pm |\alpha_0| \right] \pm \frac{\kappa_0}{|\alpha_0|^2}k^2 - \frac{1}{2} \left[ (\mu + \eta)k^2 \pm (\mu + \eta)k^2 \right] \pm \eta k^2 + O(k^3)
\] (15)

Finally, this leads to the following two eigenvalues in the hydrodynamic limit
\[
s = \begin{cases} 
-|\alpha_0| + \left[ \frac{\kappa_0}{|\alpha_0|} + \mu \right] k^2 + O(k^3) \\
- \left[ \frac{\kappa_0}{|\alpha_0|} + \eta \right] k^2 + O(k^3)
\end{cases}
\] (16)

We interpret these solutions as follows:
- the first eigenvalue \(-|\alpha_0|\) characterizes the fast relaxation of the momentum fluctuations in the absence of spatial variations \((k \to 0)\);
- the second eigenvalue \(-k^2 [\kappa_0/|\alpha_0| + \eta]\) controls the onset of instability of the homogeneous disordered phase.

We conclude that the homogeneous disordered phase is unstable when \( \alpha(\rho) < 0 \) and \( \kappa(\rho) - \eta\alpha(\kappa) < 0 \) leading to phase separation and the emergence of the cD-dD phase co-existence (e.g. MIPS).

### B. Stability of the homogeneous ordered phase

The results of the previous section hint at the fact that collective motion is expected in the case the momentum fluctuations do not relax quickly, i.e. when \( \alpha(\rho) > 0 \). Indeed, it is interesting to note that for a stable system \( \beta > 0 \) necessarily, thus \( \alpha_0 > 0 \) leads to collective motion, while \( \alpha_0 \leq 0 \) gives a stationary state. Naturally, the next step is thus to explore the stability of a homogeneous state displaying collective motion, i.e. a homogeneous ordered state. The steady-state homogeneous solutions are given by
\[
\beta p^3 - \alpha_0 p = 0 \iff p_0 = 0 \text{ or } p_0 = \sqrt{\alpha_0/\beta}
\] (17)

Here, we thus expand about a homogeneous state displaying collective motion with
\[
\begin{align*}
\rho &= \rho_0 + \delta \rho = \rho_0 + \delta \rho_0 \exp[st - ikx], \\
p &= p_0 + \delta p = p_0 + \delta p_0 \exp[st - ikx], \\
\alpha &= \alpha'_0 + \alpha'_1 \rho = \alpha_0 + \alpha_1 \delta \rho_0 \exp[st - ikx],
\end{align*}
\] (18)

with \( p_0 = \sqrt{\alpha_0/\beta} \). To linear order, we obtain for the momentum equation
\[
\begin{align*}
\partial_t \delta p + \lambda p_0 \partial_x \delta p &= \mu \partial_x^2 \delta p - \kappa \partial_x \delta p + (\alpha_0 + \alpha_1 \delta \rho)(\rho_0 + \delta \rho) - \beta (p_0 + \delta \rho)^3 \\
&= \mu \partial_x^2 \delta p - \kappa \partial_x \delta p + \alpha_0 p_0 + \alpha_1 p_0 \delta \rho + \alpha_0 \delta \rho - \beta (p_0^3 + 3p_0^2 \delta \rho) \\
&= \mu \partial_x^2 \delta p - \kappa \partial_x \delta p + \alpha_0 p_0 + \alpha_1 p_0 \delta \rho + \alpha_0 \delta \rho - \alpha_0 p_0 - 3\alpha_0 \delta \rho \\
&= \mu \partial_x^2 \delta p - \kappa \partial_x \delta p + \alpha_1 p_0 \delta \rho + \alpha_0 \delta \rho - 3\alpha_0 \delta \rho
\end{align*}
\]

The linearized equations of motion thus finally read
\[
\begin{align*}
\partial_t \delta \rho &= -\partial_x \delta p + \eta \partial_x^2 \delta p \\
\partial_t \delta p + \lambda p_0 \partial_x \delta p &= \mu \partial_x^2 \delta p - \kappa \partial_x \delta p + \alpha_1 p_0 \delta \rho - 2\alpha_0 \delta \rho
\end{align*}
\] (19) (20)

Reinjecting in these linearized equations the ansatz, we get
\[
\begin{align*}
s \delta \rho_0 &= i k \delta \rho_0 - k^2 \eta \delta \rho_0 \\
s \delta \rho_0 - i k \lambda p_0 \delta \rho_0 &= -\mu k^2 \delta p_0 + i k \kappa \delta \rho_0 + \alpha_1 p_0 \delta \rho_0 - 2\alpha_0 \delta \rho_0
\end{align*}
\] (21) (22)

Finally, we obtain the hydrodynamic limit
\[
s = \begin{cases} 
-|\alpha_0| + \left[ \frac{\kappa_0}{|\alpha_0|} + \mu \right] k^2 + O(k^3) \\
- \left[ \frac{\kappa_0}{|\alpha_0|} + \eta \right] k^2 + O(k^3)
\end{cases}
\] (16)

We interpret these solutions as follows:
- the first eigenvalue \(-|\alpha_0|\) characterizes the fast relaxation of the momentum fluctuations in the absence of spatial variations \((k \to 0)\);
- the second eigenvalue \(-k^2 [\kappa_0/|\alpha_0| + \eta]\) controls the onset of instability of the homogeneous disordered phase.

We conclude that the homogeneous disordered phase is unstable when \( \alpha(\rho) < 0 \) and \( \kappa(\rho) - \eta\alpha(\kappa) < 0 \) leading to phase separation and the emergence of the cD-dD phase co-existence (e.g. MIPS).
Once again, this can be written in matrix form as the following eigenvalue problem
\[
A \delta u_0 = s \delta u_0 ,
\]
with \( \delta u_0 = (\delta \rho_0, \delta p_0)^T \) and
\[
A = \begin{bmatrix}
-k^2 \eta & ik \\
-ik \kappa + \alpha_1 p_0 & -2 \alpha_0 + ik \lambda p_0 - \mu k^2
\end{bmatrix}
\] (24)
We know that the stability conditions are given by the sign of \( \text{Re}[s] \); the eigenvalues of this 2 \times 2 matrix are given by
\[
s = \frac{\text{Tr}A}{2} \pm \frac{1}{2} \sqrt{(\text{Tr}A)^2 - 4 \text{det}A}
\]
i.e.
\[
s = \frac{1}{2} \left[ -2 \alpha_0 + ik \lambda p_0 - k^2(\mu + \eta) \right] \\
\pm \frac{1}{2} \left[ \left( -2 \alpha_0 + ik \lambda p_0 - k^2(\mu + \eta) \right)^2 - 4(k^2 + 2 \alpha_0 \kappa \lambda k^2 + 2 \alpha_0 \eta \mu k^4 - ik \alpha_1 p_0 - ik^3 \lambda \eta p_0) \right]^{1/2}
\] (26)
Assuming that the discriminant is positive, we can expand \( s \) to lowest order in \( k \) to obtain the following two eigenvalues in the hydrodynamic limit (\( k \to 0 \))
\[
s_+ = -\frac{[8 \alpha_0^2 \eta - \alpha_1^2 p_0^2 + 4 \alpha_0^2 \kappa + 2 \alpha_0 \alpha_1 \lambda p_0^2] k^2}{8 \alpha_0^3} + i \frac{\alpha_1 p_0}{2 \alpha_0} + O(k^3)
\] (27)
and
\[
s_- = -2 \alpha_0 + \frac{[4 \alpha_0^2 (\kappa - \alpha_0 \mu) - \alpha_1^2 p_0^2 + 2 \alpha_0 \alpha_1 \lambda p_0^2 \lambda]}{8 \alpha_0^3} + i \frac{k \alpha_0}{2} \left[ 2 \kappa - \frac{\alpha_1}{\alpha_0} \right] + O(k^3)
\] (28)
Finally, at the lowest order in \( k \), the real part of \( s \) is thus given by
\[
\text{Re}[s] = \begin{cases}
-2 \alpha_0 + O(k^2) \\
\frac{k^2}{8 \alpha_0^3} \left[ \frac{\alpha_1^2}{\beta} - 2 \alpha_0 \left[ 2 \kappa + 4 \eta \alpha_0 + \frac{\alpha_1 \lambda}{\beta} \right] \right] + O(k^3)
\end{cases}
\] (29)
We interpret these solutions as follows:

- the first eigenvalue \(-2 \alpha_0\) characterizes the fast relaxation of the momentum fluctuations around the mean field value \( p_0 = \sqrt{\alpha_0 / \beta} \) in the absence of spatial variations (\( k \to 0 \));

- the second eigenvalue controls the onset of instability of the homogeneous ordered phase.

As long as \( \alpha_0 > 0 \), we have fast relaxation of the momentum fluctuations around the mean field value: the phase remains ordered. Instability leading to phase separation sets in when
\[
\frac{\alpha_1^2}{\beta} - 2 \alpha_0 \left[ 2 \kappa + 4 \eta \alpha_0 + \frac{\alpha_1 \lambda}{\beta} \right] > 0
\] (30)
First, we find the roots of the associated quadratic equation: \( \alpha_1^2 - 2 \alpha_0 \lambda \alpha_1 - 4 \alpha_0 \beta \kappa - 8 \eta \alpha_0^2 \beta = 0 \). The discriminant of this quadratic equation is given by
\[
\Delta = 4 \alpha_0^2 \lambda^2 + 16 \alpha_0 \kappa \beta + 32 \eta \alpha_0 \beta
\] (31)
The sign of the discriminant dictates the existence of real roots. As \( \beta > 0 \), if \( \kappa < -[\alpha_0 \lambda^2 / 4 \beta + 2 \eta] \), the quadratic equation does not have real solutions and the inequality is always met. This means that the homogeneous state is thus always unstable.

Conversely, if \( \kappa > -[\alpha_0 \lambda^2 / 4 \beta + 2 \eta] \), the roots of the quadratic equation are given by
\[
\alpha_1 = \alpha_0 \lambda \pm \sqrt{\alpha_0^2 \lambda^2 + 4 \alpha_0 \kappa \beta + 8 \eta \alpha_0 \beta}
\] (32)
We conclude that if $\kappa > -[\alpha_0 \lambda^2 / 4 \beta + 2\eta]$, the instability sets in only when

$$
\alpha_1 < \alpha_0 \lambda \left[ 1 - \sqrt{1 + \frac{4(\kappa + 2\eta)\beta}{\alpha_0 \lambda^2}} \right] \quad \text{or} \quad \alpha_1 > \alpha_0 \lambda \left[ 1 + \sqrt{1 + \frac{4(\kappa + 2\eta)\beta}{\alpha_0 \lambda^2}} \right]
$$

(33)

Under these instability conditions, the homogeneous ordered phase is unstable. We will see in what follows that in this case, the instability can lead to the emergence of one of the three remaining phase co-existences: cO-dD \textit{(banding)}, cD-dO \textit{(reverse banding)} or cO-dO \textit{(comoving phases)}. Interestingly, in the ordered phase, the homogeneous state can remain stable even if $\kappa$ is negative, as long as the above stability conditions are satisfied.

\section{II. PARTICULAR MODELS}

\subsection{A. A first example}

In general, we are free to choose any functional form for $\alpha(\rho)$ and $\kappa(\rho)$. Here, we consider a specific model and use the previous results to study when this model can lead to non-trivial phase co-existence. We consider the following

\begin{figure}[h]
\begin{center}
\includegraphics[width=\textwidth]{flowchart.png}
\end{center}
\caption{Conditions for instability stemming from the linear stability analysis; the colorscheme used corresponds to the color of the instability regions in the phase diagrams.}
\end{figure}
model

\[ \alpha(\rho) = -A + \rho - \rho^2 \]  
\[ \kappa(\rho) = K - \frac{5}{6}\rho + \rho^2 \]  

and \( \lambda = \mu = \beta = 1 \).

We have seen above that the condition for existence of collective motion is \( \alpha(\rho) > 0 \). In this model, \( \alpha \) can become positive when the discriminant of the quadratic equation \( \rho^2 - \rho + A = 0 \) is positive, i.e. when \( A < 1/4 \). We conclude that collective motion is possible in this model when

\[ \frac{1}{2} \left(1 - \sqrt{1 - 4A}\right) < \rho < \frac{1}{2} \left(1 + \sqrt{1 - 4A}\right) \]  

In particular, within the region where \( \alpha(\rho) < 0 \), we can linearize the expression of \( \alpha \) and write

\[ \alpha(\rho) = (-A + \rho_0 - \rho^2_0) + (1 - 2\rho_0) \delta \rho + O(\delta \rho^2) \]  

where we have introduced \( \rho = \rho_0 + \delta \rho \). Thus, identifying this expression to the results we derived in Section IB, we write \( \alpha(\rho) = \alpha_0 + \alpha_1 \delta \rho \) with

\[ \alpha_0 = (-A + \rho_0 - \rho^2_0) \]  
\[ \alpha_1 = (1 - 2\rho_0) \]  

Substituting this into Equation (30), we obtain the following condition

\[ \frac{(1 - 2\rho_0)^2}{\beta} - 2(-A + \rho_0 - \rho^2_0) \left[ 2 \left( K - \frac{5\rho_0}{6} + \rho^2_0 \right) + 4\eta \left(-A + \rho_0 - \rho^2_0\right) + \frac{1 - 2\rho_0}{\beta} \right] > 0 \]  

The locus corresponds to the roots of a quartic equations and the analytical expressions are not illuminating. However, the instability regions can be easily obtained numerically. Figure 1 summarizes the conditions for instability. In Figures 2 and 3, colored regions correspond to the colors in Figure 1, i.e.:

- **white region**, homogeneous disordered;
- **red region**, \( \alpha_0 < 0 \) and \( \kappa_0 - \eta \alpha_0 < \) leading to a phase co-existence;
- **orange region**, \( \alpha_0 > 0 \) and \( \kappa_0 < -[\alpha_0 \lambda^2/4\beta + 2\eta] \) leading to a phase co-existence;
- **blue region**, \( \alpha_0 > 0 \) and \( \kappa_0 > -[\alpha_0 \lambda^2/4\beta + 2\eta] \) and \( \alpha_1^2/\beta - 2\alpha_0 [2\kappa + 4\eta \alpha_0 + \alpha_1 \lambda/\beta] > 0 \) leading to a phase co-existence;
- **grey region**, \( \alpha_0 > 0 \) but no instability condition met leading to a homogeneous ordered phase.

Note that the color scheme used here refers to the parameter ranges that lead to instability, which is distinct from the color scheme used in Fig. 2 in the main text, which depicts the types of phase co-existence of the system.
FIG. 2. Instability regions for the model defined by Eqs. (34) and (35) with fixed $K$ with values (a) 0.175, (b) 0.15, (c) 0.125 and (d) 0.1.

FIG. 3. Instability regions for the model defined by Eqs. (34) and (35) with fixed $A$ with values (a) 0.25, (b) 0.225, (c) 0.2 and (d) 0.175.
B. Phase diagrams

Using the method outlined in the main text, we can construct qualitatively the corresponding phase diagrams given the instability regions. Briefly, the crucial ingredient here is that the instability regions is analogous to the spinodal decomposition region in thermal phase separation, which is always flanked by the nucleation and growth regions, which are metastable (and hence stable under linear stability analysis). Therefore, the phase separation boundaries (or the binodal lines) always extend further from the instability boundaries.

For instance, the corresponding phase diagrams of Fig. 2(a) and (b) are depicted qualitatively in Fig. 4, and those of Fig. 3(a) and (b) are shown in Fig. 5. In Figs. 4 and 5, all instability regions (obtained via our linear stability analysis) are shown as a shaded (grey) area while homogeneous regions are shown in white.

C. A model displaying the new comoving phase co-existence

We now discuss the model system discussed in the main text, in which

$$\alpha(\rho) = -A + 18\rho - 10/3\rho^2$$  \hspace{1cm} (41)
$$\kappa(\rho) = 140 - 145\rho + 30\rho^2 .$$ \hspace{1cm} (42)

The instability regions are shown in Fig. 6 and the corresponding phase diagram is shown in Fig. 2 in the main text. In the next section, we provide details of the numerical methods that enables us to determine the various phase boundaries for this system.
FIG. 6. Instability regions for the system of Equations (41–42).

III. NUMERICAL METHODS

The stationary profiles displayed in the main text are obtained via direct numerical integration of Eqs. (1) and (2). To do so, we discretize space using a linearly-spaced grid with spacing $\Delta x$ and use a central finite difference approximation accurate to order $O(\Delta x^8)$ for the spatial derivatives. The resulting set of ordinary differential equations is integrated using a variable order implicit scheme with adaptive time-stepping (ODE15s from MATLAB — https://www.mathworks.com/help/matlab/ref/ode15s.html).

The results of numerical simulations shown in the main text were obtained for a domain size $L = 100$ and $\Delta x \in [0.1, 0.2]$, with periodic boundary conditions. In all our simulations, we used the following parameters $\eta = 2$, $\lambda = 1$, $\beta = 0.5$ and $\mu = 1$. We made sure to reach a stable steady-state in our simulations by simulating the system for at least a total time $T = 500$. To confirm the consistency of our results, we compared the steady-state profiles obtained for a variety of initial conditions (e.g., by varying the values of $\rho_c$, $\rho_d$, $p_d$, and $p_c$ as described below). For the data shown in Fig. 1 and Fig. 2 of the main text, we used as initial conditions a double sigmoid with sharp interfaces:

$$
\rho_1(x) = \rho_d + \frac{\rho_c - \rho_d}{1 + \exp \left[-(x - L/4)\right]}
$$

$$
\rho_2(x) = \rho_d + \frac{\rho_c - \rho_d}{1 + \exp \left[(x - 3L/4)\right]}
$$

$$
\rho(x, 0) = \min(\rho_1(x), \rho_2(x))
$$

and if $p_c > p_d$,

$$
p_1(x) = p_d + \frac{p_c - p_d}{1 + \exp \left[-(x - L/4)\right]}
$$

$$
p_2(x) = p_d + \frac{p_c - p_d}{1 + \exp \left[(x - 3L/4)\right]}
$$

$$
p(x, 0) = \min(p_1(x), p_2(x))
$$

while if $p_c < p_d$,

$$
p_1(x) = p_c - \frac{p_c - p_d}{1 + \exp \left[-(x - L/4)\right]}
$$

$$
p_2(x) = p_c - \frac{p_c - p_d}{1 + \exp \left[(x - 3L/4)\right]}
$$

$$
p(x, 0) = \max(p_1(x), p_2(x))
$$

To produce the binodal lines in Fig. 2 of the main text, we varied $A$ (by small increments of $\Delta A = 0.25$) and initialized the system with densities taken at the edges of the corresponding instability region. We summarize in the
Table I the parameters used for the 6 profiles shown in Fig. 1 of the main text and the associated supplementary movies S1 to S6.

| Supp. Movie | Phase co-existence | Fig. Panel | A  | p₀  | ρ₀ | ρ₋ | ρ₊ | p₋ | p₊ |
|-------------|-------------------|------------|----|-----|----|----|----|----|----|
| S1          | dD-cD             | Fig. 1(a)  | 24 | 2.5 | 1.5| 3.5| 0  | 0  |
| S2          | dO-cO             | Fig. 1(b) (blue line) | 7  | 2.5 | 1.0| 4.0| 1  | 2  |
| S3          | dO-cO             | Fig. 1(b) (red line) | 8  | 2.5 | 1.0| 4.0| −1 | 1  |
| S4          | dD-cO             | Fig. 1(c) (blue line) | 7  | 0.45| 0.4| 0.5| 1  | 0  |
| S5          | dD-cO             | Fig. 1(c) (red line) | 14.75 | 2.35| 0.7| 4.0| 0  | 1  |
| S6          | dO-cD             | Fig. 1(d)  | 14 | 4.45| 4.4| 4.5| 0  | −1 |

TABLE I. Simulation parameters \((A, p₀)\) for steady-state profiles from Fig. 1 of the main text (and corresponding movies S1 to S6); we also include in the table the initial conditions parameters \((ρ₀, ρ₋, ρ₊, p₋, p₊)\).

IV. MULTICRITICAL POINT

At the linear level around the multi-critical point \((κ₀ = κ₁ = α₀ = α₁ = 0)\), the EOM are

\[
\partial_t \delta \rho + \nabla \cdot p = \eta \nabla^2 \delta \rho , \quad \partial_t p = \mu \nabla^2 p + f ,
\]

\(f\) is a Gaussian noise term with a non-zero standard deviation.

Performing the following re-scalings:

\[
\begin{align*}
    r &\mapsto e^\ell r , \\
    t &\mapsto e^\chi p \ell t , \\
    \delta \rho &\mapsto e^{\chi ρ \ell} \delta \rho , \\
    p &\mapsto e^{\chi p \ell} p ,
\end{align*}
\]

the EOM become

\[
e^{(\chi ρ - z)\ell} \partial_\ell \delta \rho + e^{(\chi ρ - 2)\ell} \eta \nabla^2 \delta \rho , \quad e^{(\chi ρ - z)\ell} \partial_\ell p = e^{(\chi p - 2)\ell} \mu \nabla^2 p + e^{-(z+d)\ell/2} f .
\]

Therefore, the above linear EOM remain invariant if we make the following choice:

\[
\begin{align*}
    z &= 2 , \\
    \chi p &= \frac{2 - d}{2} , \\
    \chi ρ &= \frac{4 - d}{2} .
\end{align*}
\]

Substituting these values into all possible nonlinear terms in the EOM of \(p\) (which is identical to the Toner-Tu equation), we find that as one decreases \(d\) from infinity, the first nonlinear term whose re-scaling based on the linear result diverges as \(\ell \to \infty\) is the term \(\nabla \delta \rho^2\). This is because \(\nabla \delta \rho^2 \mapsto e^{(1-d-\frac{1}{2})\ell} \nabla \delta \rho^2\), and the exponent \((3-d)\ell\) becomes greater than the re-scaling exponent of \(\partial_\ell p\), which is \((\chi ρ - z)\ell = -(d+2)\ell/2\), at the upper critical dimension \(d_c = 8\).