Critical mingling and universal correlations in model binary active liquids

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Ensembles of driven or motile bodies moving along opposite directions are generically reported to self-organize into strongly anisotropic lanes. Here, building on a minimal model of self-propelled bodies targeting opposite directions, we first evidence a critical phase transition between a mingled state and a phase-separated lane state specific to active particles. We then demonstrate that the mingled state displays algebraic structural correlations also found in driven binary mixtures. Finally, constructing a hydrodynamic theory, we single out the physical mechanisms responsible for these universal long-range correlations typical of ensembles of oppositely moving bodies.
S

hould you want to mix two groups of pedestrians, or two
er ensembles of colloidal beads, one of the worst possible
trategies would be pushing them towards each other.
Both experiments and numerical simulations have demonstrated
the segregation of oppositely driven Brownian particles into
parallel lanes. Even the tiniest drive results in the formation of
finite slender lanes which exponentially grow with the driving
strength. The same qualitative phenomenology is consistently
observed in pedestrian counterflows. From our daily
observation of urban traffic to laboratory experiments, the
emergence of counter-propagating lanes is one of the most
robust phenomena in population dynamics, and has been at
the very origin of the early description of pedestrians as granular
materials. However, a description as isotropic grains is
usually not sufficient to account for the dynamics of interacting
motive bodies. From motility-induced phase separation, to
giant density fluctuations in flocks, to pedestrian scattering,
the most significant collective phenomena in active matter stem
from the interplay between their position and orientation
degrees of freedom.

In this communication, we address the phase behaviour of a
binary mixture of active particles targeting opposite directions.
Building on a prototypical model of self-propelled bodies
with repulsive interactions, we numerically evidence two non-
equilibrium steady states: a lane state where the two populations
maximize their flux and phase separate, and a mixed state where
all motile particles mingle homogeneously. We show that these
two distinct states are separated by a genuine critical phase
transition. In addition, we demonstrate algebraic density correla-
tions in the homogeneous phase, akin to that recently reported
for oppositely driven Brownian particles. Finally, we construct
a hydrodynamic description to elucidate these long-range
structural correlations, and conclude that they are universal to
both active and driven ensembles of oppositely moving bodies.

Results

A minimal model of active binary mixtures. We consider
an ensemble of self-propelled particles characterized
by their instantaneous positions \( \mathbf{r}_i(t) \) and orientations
\( \mathbf{p}_i(t) = (\cos \theta_i, \sin \theta_i) \), where \( i = 1, \ldots, N \) (in all
that follows \( \mathbf{x} \) stands for \( x/|x| \)). Each particle moves along its orientation vector
at constant speed (|\( \mathbf{r}_i | = 1 \)). We separate the particle ensemble into
two groups of equal size following either the direction \( \Theta = 0 \)
(right movers) or \( \pi \) (left movers) according to a harmonic angular
potential \( V(\theta_i) = \frac{\pi}{2} (\theta_i - \Theta)^2 \). Their equations of motion take the
simple form:

\[
\begin{align*}
\dot{\mathbf{r}}_i &= \mathbf{p}_i, \\
\dot{\theta}_i &= -\partial_{\theta_i} V(\theta_i) + \sum_j T_{ij}.
\end{align*}
\]

In principle, oriented particles can interact by both forces
and torques. We here focus on the impact of orientational couplings
and consider that neighbouring particles interact solely through
pairwise additive torques \( T_{ij} \). This type of model has been
successfully used to describe a number of seemingly different
active systems, starting from bird flocks, fish schools and bacteria
colonies to synthetic active matter made of self-propelled colloids
or polymeric biofilaments. We here elaborate on a minimal
construction where the particles interact only by repulsive
torques. In practical terms, we choose the standard
form \( T_{ij} = -\partial_{\theta_i} E_{ij} \), where the effective angular energy simply
reads \( E_{ij} = -B r_{ij} \cdot \mathbf{p}_i \mathbf{p}_j \). As sketched in Fig. 1a, this interaction
promotes the orientation of \( \mathbf{p}_i \) along the direction of the centre-to-centre vector \( \mathbf{r}_{ij} = (\mathbf{r}_i - \mathbf{r}_j) \): as they interact particles turn their
back to each other (for example, refs 24, 28–30). The spatial decay
of the interactions is given by: \( B(r_{ij}) = B(1 - r_{ij}/(a_i + a_j)) \), where \( B \)
is a finite constant if \( r_{ij}(a_i + a_j) \) and 0 otherwise. In all that
follows, we focus on the regime where repulsion overcomes
alignment along the preferred direction \( B > 1 \). The interaction
ranges \( a_i \) are chosen to be polydisperse to avoid the specifics of
crystallization, and we make the classic choice \( a = 1 \) or 1.4 for
one in every two particles. Before solving equations (1) and (2), two

Figure 1 | Phase behaviour. (a) Trajectories of two particles interacting
solely by a repulsive torque as defined in equation (2) with \( B = 5 \). The
post-collision orientations \( \mathbf{p}_i(t) \) are along the centre-to-centre axis \( r_{ij} \).
(b–d) Snapshots of a square window at the centre of the simulation box
\((L_x = 168, N = 1,973, \pi \sigma^2 = 0.65) \), respectively, in the lane \((B = 2) \) and
the homogeneous \((B = 5) \) states. The arrows indicate the instantaneous
position and orientation of the particles. Dark blue: right movers. Light blue: left
movers. (c) Phase diagram. \( \pi \sigma^2 \) is the particle area fraction. Filled symbols:
homogeneous state. Open symbols: lanes. (e) Probability distribution
function (p.d.f.) of the density difference \( \bar{\rho} = \rho_i - \rho_i \). Light orange line:
\( B = 2 \), \( \pi \sigma^2 = 0.65 \). Dark blue line: \( B = 5 \), \( \pi \sigma^2 = 0.65 \). Dashed line: best Gaussian fit.
(f) p.d.f. of the orientational fluctuations around the preferred orientation
 Investors and colours as in e. Inset: orientational diffusion \( D_{ij} \) in the
homogeneous state at a fixed repulsion magnitude \( B = 5 \) and
different particle area fractions \( \pi \sigma^2 \). \( D_{ij} \) is defined as the
decorrelation time of the particle orientation. In the mingled state,
the velocity autocorrelation decays exponentially at short time,
\( D_{ij} \) is therefore defined without ambiguity, see also Supplementary Note 1 for a full
description of the numerical computation of \( D_{ij} \). Dashed line: best linear fit.
Critical mingling. Starting from random initial conditions, we numerically solve equations (1) and (2) using forward Euler integration with a time step of $10^{-2}$, and a sweep-and-prune algorithm for neighbour summation. We use a rectangular simulation box of aspect ratio $L_x=2L_y$ with periodic boundary conditions in both directions. We also retrain our analysis to $H=1$, leaving two control parameters that are the repulsion strength $B$ and the overall density $\rho$. The following results correspond to simulations with $N$ comprised between 493 and 197,300 particles.

We observe two clearly distinct stationary states illustrated in Fig. 1b,d. At low density and/or weak repulsion the system quickly phase separates. Computing the local density difference $\Delta N = N_r - N_l$ and analytically demonstrate their universality. The overall pair correlation function of the active liquid, $g_r(r)$, is plotted in Fig. 4a. (a,b) Log plots at five densities for a box of length $L_x=336$ ($N$ ranges from 5,462 to 7,892). (c) Fluctuations of the order parameter plotted versus $B - B_c$ for the same densities as in (b). The fluctuations are defined as $\Delta(W_{ij}) = \langle W_{ij}^2 \rangle - \langle W_{ij} \rangle^2$ with $B < B_c$. (d) Correlation time $\tau_B$ plotted against $B - B_c$. The correlation time is defined as $\langle W(t) \rangle = \tau_B$ for the same densities as in (b). The fluctuations are defined as $\Delta(W_{ij}) = \langle W_{ij}^2 \rangle - \langle W_{ij} \rangle^2$ with $B < B_c$. All error bars correspond to two standard deviations. The error on the estimate of the exponents correspond to one s.d. after considering linear fits for each density.

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Long-range correlations in mingled liquids. We now evidence long-range structural correlations in this active-liquid phase, and analytically demonstrate their universalality. The overall pair correlation function of the active liquid, $g_r(r)$, is plotted in Fig. 4a. At a first glance, deep in the homogeneous phase, the few visible oscillations would suggest a simple anisotropic liquid structure. However, denoting $z$ and $\beta$ the preferred direction of the populations (left or right), we find that the asymptotic behaviours of all pair correlation functions $g_{ij}(x,y) = 0$ decay algebraically as $|1-g_{ij}(x,0)| \sim x^{-\nu_z}$ with $\nu_z \approx 1.5$, Fig. 4b. This power-law

\[
\langle W \rangle = (1 - \cos(\theta_i - \theta_j)), \quad \langle W \rangle = (1 - \cos(\theta_i - \theta_j)).
\]
behaviour is very close to that reported in numerical simulations\(^4\) and fluctuating density functional theories of oppositely driven colloids at finite temperature\(^20\).

**Hydrodynamic description.** To explain the robustness of these long-range correlations, we provide a hydrodynamic description of the mingled state, and compute its structural response to random fluctuations. We first observe that the orientational diffusivity of the particles increases linearly with the average density \(\bar{\rho}\) in Fig. 1f inset. This behaviour indicates that binary collisions set the fluctuations of this active liquid, and hence suggests using a Boltzmann kinetic-theory framework, for example, refs 32,33 from an active-matter perspective. In the large \(B\) limit, the microscopic interactions are accounted for by a simplified scattering rule anticipated from equation (2) and confirmed by the inspection of typical trajectories (Fig. 1a). Upon binary collisions the self-propelled particles align their orientation with the centre-to-centre axis regardless of their initial orientation and external drive. Assuming molecular chaos and binary collisions the self-propelled particles align their orientation around the average homogeneous state (see Methods section for technical details). Within a linear response approximation, they take the compact form:

\[
\partial_t \delta \rho_z(r, t) + \nabla \cdot (J_z + \bar{J}) = 0, \tag{5}
\]

where \(J_z\) describes the convection and the collision-induced diffusion of the \(z\) species, and \(\bar{J}\) is the coupling term, crucial to the anomalous fluctuations of the active liquid:

\[
J_z = v_0 \hat{h}_z \delta \rho_z - \mathbf{D} \cdot \nabla \delta \rho_z, \tag{6}
\]

\[
\bar{J} = -\bar{v} \hat{h}_z \delta \bar{\rho} - \bar{D} \cdot \nabla \bar{\delta}. \tag{7}
\]

The two anisotropic diffusion tensors \(\mathbf{D}\) and \(\bar{D}\) are diagonal and their expression is provided in Supplementary Note 3 together with all the hydrodynamic coefficients. \(\bar{J}\) is a particle current stemming from the fluctuations of the other species and has two origins. The first term arises from the competition between alignment along the driving direction \(\hat{h}_z\) and orientational diffusion caused by the collisions: the higher the local density \(\bar{\rho}\), the smaller the longitudinal current. The second term originates from the pressure term \(\propto \nabla \bar{\rho}\): a local density gradient results in a net flow of both species (see Methods section for details). This diffusive coupling is therefore generic and enters the description of any binary compressible fluid. Two additional comments are in order. First, this prediction is not specific to the small-density regime and is expected to be robust to the microscopic details of the interactions. As a matter of fact, the above hydrodynamic description is not only valid in the limit of strong repulsion and small densities discussed above but also in the opposite limit, where the particle density is very large while the repulsion remains finite as detailed in Supplementary Note 5. Second, the robustness of this hydrodynamic description could have been anticipated using conservation laws and symmetry considerations, as done for example, in ref. 16 for active flocks.
Here the situation is simpler, momentum is not conserved and no soft mode is associated to any spontaneous symmetry breaking. As a result the only two hydrodynamic variables are the coupled (self-advected) densities of the two populations. The associated mass currents are constructed from the only two vectors that can be formed in this homogeneous but anisotropic setting: \( \mathbf{h}_a \) and \( \nabla \rho_a \). These simple observations are enough to set the functional form of equations (5)–(7).

By construction the above hydrodynamic description alone cannot account for any structural correlation. To go beyond this mean-field picture we classically account for fluctuations by adding a conserved noise source to equation (5) and compute the resulting density-fluctuation spectrum. At the linear response level, without loss of generality, we can restrain ourselves to the case of an isotropic additive white noise of variance \( 2 \lambda^2 \). As a result the only two hydrodynamic variables are the coupled soft mode is associated to any spontaneous symmetry breaking.

Relation functions take the form\( h_j \sim |q|^{2j}/(2j)! \) readily demonstrated noting that \( \langle \gamma^2 \rangle \sim \langle \gamma \rangle^2 \) from different directions yields different limits, which is anisotropy, and Supplementary Note 4). The structure factor has been computed from hydrodynamic equations common to driven binary mixtures. We first recall that this structure factor between a phase separated and a mingled state in binary active and driven particles. Altogether these observations confirm the universality of the long-range structural correlations found in both classes of non-equilibrium mixtures.

In conclusion, we have demonstrated that the interplay between orientational and translational degrees of freedom, inherent to motile bodies, can result in a critical transition between a phase separated and a mingled state in binary active mixtures. In addition, we have singled out the very mechanisms responsible for long-range structural correlations in any ensemble of particles driven towards opposite directions, should they be passive colloids or self-propelled agents.

**Methods**

**Boltzmann kinetic theory.** Let us summarize the main steps of the kinetic theory employed to establish equations (5)–(7). The so-called collision integral on the r.h.s of equation (4) includes two contributions, which translate the behaviour illustrated in Fig. 1a:

\[
\mathcal{I}^a_{\text{coll}}(\mathbf{r}) = \mathcal{I}^a_{\text{coll}}(\mathbf{r} - 2a\mathbf{p}) - \mathcal{I}^a_{\text{coll}}(\mathbf{r}, \mathbf{p}, \mathbf{r} \cdot \mathbf{p}).
\]

The first term indicates that a collision with any particle located at \( (\mathbf{r} - 2a\mathbf{p}) \) reorients the \( a \) particles along \( \mathbf{p} \) at a rate \( \mathcal{I}^a_{\text{coll}} \). The second term accounts for the random reorientation, at a rate \( \mathcal{I}^{\text{coll}} \), of a particle aligned with \( \mathbf{p} \) on collision with any other particle. Within a two-fluid picture, the velocity and nematic texture of the \( a \) particles are given by \( v_a = v_a^{-1} (\mathbf{p} \cdot \mathbf{v}) \) and \( Q_a = Q_a^{-1} (\mathbf{p} \cdot \mathbf{v}) \). The mass conservation relation, \( \partial_t \rho_a + \nabla \cdot (\rho_a \mathbf{v}_a) = 0 \), is obtained by integrating equation (4) with respect to \( \theta \) and constraints \( 2(\mathcal{I}^{\text{coll}}) = \mathcal{I}^{\text{coll}} = D \). The time evolution of the velocity field is also readily obtained from equation (4):

\[
\partial_t \rho_a \mathbf{v}_a = \nabla \cdot \left( \rho_a \mathbf{F}_a \right) = \mathbf{F}_a,
\]

where the second term on the l.h.s is a convective term stemming from self-propulsion. The force field \( \mathbf{F}_a \) on the r.h.s of equation (10) reads:

\[
\mathbf{F}_a = \rho_a^{-1/2} \Big( \mathbf{h}_a + (aD_p) \nabla \mathbf{p} - (D_p) \mathbf{p} \Big) \rho_a^{-1} \mathbf{v}_a.
\]

This term originates from the alignment of particles along the \( h_a \) direction, the second term is a repulsion-induced pressure, and the third one echoes the collision-induced rotational diffusivity of the particles. An additional closure relation between \( Q_a \) and \( \rho_a \) is required to yield a self-consistent hydrodynamic description. Deep in the homogeneous phase, we make a wrapped Gaussian approximation for the orientational fluctuations in each population. The hypothesis is equivalent to setting \( Q_a = v_a^{-1} \left( \langle \mathbf{v} \cdot \mathbf{v} \rangle - 1 \right) \) (refs 24,37). As momentum is not conserved, the velocity field is not a hydrodynamic variable; in the long wavelength limit the velocity modes relax much faster than the (conserved) density modes. We therefore ignore the temporal variations in equation (10) and use this simplified equation to eliminate \( v_a \) in the mass conservation relation, leading to the mass conservation equation (5).

**Data availability.** The data that support the findings of this study are available from the corresponding author upon request.

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**Author contributions**

D.B. designed the research. N.B. performed the numerical simulations. D.B. and N.B. performed the theory, discussed the results and wrote the paper.

**Additional information**

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