

Phase Separation in a Chiral Active Fluid of Inertial Self-Spinning Disks

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We show that systematic particle rotations in a fluid composed of disk-shaped spinners can spontaneously lead to phase separation. The phenomenon arises out of a homogeneous and hydrostatic stationary state, due to a pressure feedback mechanism that increases local density fluctuations. We show how this mechanism induces phase separation, coined as rotation induced phase separation (RIPS), when the active rotation is not properly counterbalanced by translational friction. A low density phase can coexist with a dense chiral liquid due to the imbalance between pressure and stress transmitted through chiral flows when a significant momentum transfer between rotational and translational motion can be sustained. As a consequence, RIPS is expected to appear as a general nonequilibrium phenomenon in chiral fluids.

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Intense and active research on complex fluids dynamics during the last decades has expanded the limits of non-equilibrium statistical mechanics and hydrodynamics, leading to generalized theories that describe a wide variety of nonequilibrium systems [1,2]. The need for these generalized theories stems from strong experimental and computational evidence that cannot be explained within the context of classical fluid and statistical mechanics. For instance, hydrodynamic instabilities [3,4], phase transitions associated with changes in system's symmetry [5,6] and memory formation [7,8] display a rich phenomenology in a wide variety of materials, including granular fluids, active fluids, magnetic colloids, etc. [9]. Most interestingly, phase separation appears to be ubiquitous in granular [10,11] and active fluids [12–14], controlled by mechanisms that extend far beyond their equilibrium counterparts.

In granular matter, energy is dissipated at the single grain scale through frictional forces at collisions. To prevent dynamical arrest, an external action needs to be applied, so that by constant input of kinetic energy the system goes into stationary, inherently out of equilibrium states. In quasi two-dimensional granular layer experiments, energy can be

injected, e.g., shaking the sample along the third direction, and grains gain energy from collisions, both with the confining walls and neighboring particles [15,16]. In such an experimental setup, phenomena like crystallization (hexatic [5] and cubic [17]) and phase separation [18] have been realized.

As opposed to granular fluids, where particle translations and rotations are passive, active fluids are characterized by a mechanism that produces systematic translation and/or rotation at the particle level, leading generically to new types of instabilities and phase transitions, such as flocking [19] and motility-induced phase separation [12,20]. Although phase separation has also been observed recently in chiral fluids with active rotations, no generic mechanism has yet been identified [21,22].

Recently, a system of disks with tilted radial blades has been experimentally proposed. Disks acquire persistent rotation when subjected to air upflow [23–28]. Complementarily, experimental systems with similar chiral interactions, ranging from self-spinning grains and robots [29–33] to magnetic colloids [34–40] and optically activate microrotors [41] have been proposed. A generic property of transport phenomena in these fluids is that their fluxes inherently have an antisymmetric component. As a consequence, a set of new transport coefficients, labeled as *odd*, emerges [42]. The collective behavior of chiral fluids in these systems has been explored theoretically and computationally [4,21,22,43–49], but, despite these efforts, the possibility that a chiral fluid undergoes generically a phase separation with specific features due to its oddness remains an open question. We show here that phase separation may be generically induced in chiral systems from the balance between rotational activity and friction.

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To this end, we introduce a minimal model describing a system of N identical rotating disks with mass m , diameter σ , and moment of inertia I . The disks are subject to forces and torques coming from three sources: friction forces and torques with the embedding medium (with translational γ and rotational γ_θ friction coefficients), interparticle forces, and an active torque. The system is determined by a Langevin dynamics where the thermal bath, of temperature T_{th} , is a white noise for both translation ξ and rotation ξ_θ with zero mean and time correlations $\langle \xi_i(t) \xi_j(t') \rangle = 2\gamma T_{\text{th}} \mathbf{1}_{\delta_{ij}} \delta(t-t')$ and $\langle \xi_{\theta,i}(t) \xi_{\theta,j}(t') \rangle = 2\gamma_\theta T_{\text{th}} \delta_{ij} \delta(t-t')$. Namely,

$$m \frac{d\mathbf{v}_i}{dt} = -\gamma \mathbf{v}_i + \xi_i + \mathbf{F}_i^n + \mathbf{F}_i^t, \quad (1)$$

$$I \frac{d\omega_i}{dt} = -\gamma_\theta \omega_i + \xi_{\theta,i} \hat{\mathbf{u}}_z + \boldsymbol{\tau}_i + \boldsymbol{\tau}_0. \quad (2)$$

The interactions between disks are controlled by normal and tangential forces. The former, $\mathbf{F}_i^n = -\nabla_i \sum_{j \neq i} U(|\mathbf{r}_j - \mathbf{r}_i|)$, is due to a purely repulsive Weeks-Chandler-Anderson (WCA) potential [50], $U(r) = \varepsilon[2(\sigma/r)^{12} - (\sigma/r)^6]$, cut at its minimum $r_c = \sigma$. The tangential force, $\mathbf{F}_i^t = -\sum_{j \neq i} \eta \mathbf{v}_{ij}^t$, of range $r_{ij} < \sigma$ (r_{ij} being the interparticle distance), is due to friction between particles and characterized by a friction coefficient η . The relative tangential velocity at the contact point is defined as $\mathbf{v}_{ij}^t = \mathbf{v}_{ij} - (\mathbf{v}_{ij} \cdot \hat{\mathbf{n}}_{ij}) \hat{\mathbf{n}}_{ij} - \omega_{ij} \times \hat{\mathbf{n}}_{ij}$, where $\mathbf{v}_{ij} = \mathbf{v}_j - \mathbf{v}_i$, $\omega_{ij} = (\omega_i + \omega_j)\sigma/2$, and $\hat{\mathbf{n}}_{ij} = (\mathbf{r}_i - \mathbf{r}_j)/r_{ij}$. The tangential force induces an angular momentum impulse upon particle encounters, due to the torque $\boldsymbol{\tau}_i = -\frac{1}{2} \sum_{j \neq i} (\mathbf{r}_{ij} \times \mathbf{F}_{ij}^t)$. Finally, disks are subject to a homogeneous and constant torque $\boldsymbol{\tau}_0$, which generates a persistent rotation at a fixed rate $\omega_0 = \tau_0/\gamma_\theta$, also referred to as spin.

Figure 1 shows the combined action of tangential friction and active rotation in disks' collisions, which breaks the time parity when converting rotational into translational kinetic energy. In particular, the rotational kinetic energy may drop at collisions, then relaxing to ω_0 , as imposed by the external forcing, after the collision. Consequently, the translational kinetic energy shows a sharp increase in collisions, and subsequently decreases because of medium friction. Overall, rotational kinetic energy is injected into the system at particle level and transformed into translational kinetic energy upon collisions, breaking detailed balance and driving the system toward an out of equilibrium and chiral state.

We have performed molecular dynamics (MD) simulations [51] of $N = 256^2$ disks following the equations of movement Eqs. (1) and (2) in a system of size L , and explore the effect of intrinsic spin ω_0 and surface fraction $\phi = N\pi\sigma^2/4L^2$ on the collective behavior of the system.

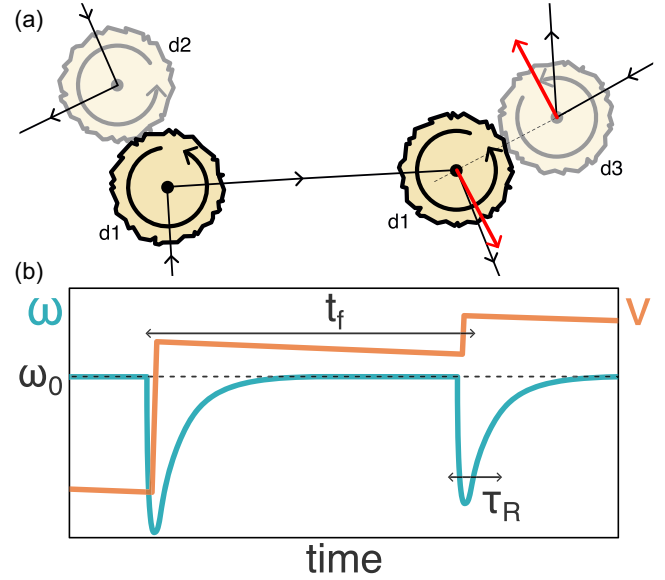


FIG. 1. (a) A sketch of two consecutive collisions of a disk $d1$ (solid colors), with corresponding incoming and outgoing trajectories (black lines and arrows). The first collision is with disk $d2$, the second one is with disk $d3$. The disks are represented with a rough surface, denoting the action of a tangential frictionlike force (symbolized as red arrows in the case of the second collision). (b) Time evolution of disk $d1$ speed (v , orange line) and angular velocity (ω , cyan line), over a time interval that includes the two consecutive collisions. The rotational inertia timescale τ_R and the free flight time t_f are the two relevant timescales. In this case, rotational dynamics is much faster, hence $\tau_R < t_f$, and the disk recovers the intrinsic spinning velocity ω_0 before colliding again. The translational inertial timescale τ_t , not displayed here, is associated with friction and is much larger than the other two, since in our system energy loss due to friction in between two collisions is small.

Unless otherwise stated, $T_{\text{th}} = 1/100$, $\gamma = 1/10$, $\gamma_\theta = 100$, $\eta = 100$, in units of mass m , length σ , and energy ε .

Spinners possess translational and rotational inertia, characterized by timescales $\tau_t = m/\gamma$ and $\tau_R = I/\gamma_\theta$, respectively. We focus on the regime $\tau_R \ll \tau_t$, where chiral effects are strong. For an arbitrary rotating disk, the postcollisional spin ω relaxes rapidly to ω_0 , while the translational velocity v slowly decays due to friction with the surrounding medium [Fig. 1(b)].

For large enough ϕ and ω_0 , the system segregates into dense regions that leave nearly empty areas. Figure 2(a) shows representative configurations of the system in the late-time regime. The phenomenology of this rotation induced phase separation (RIPS) is very rich with regard to the changes in the structure of the coexisting phases, the shape of their interface, and their dynamics.

Increasing ω_0 at constant ϕ , we identify a value above which the system develops a gaseous void surrounded by a dense phase. This critical persistence angular velocity $\omega_{0,c}$ is identified by a first increase of the root mean square time

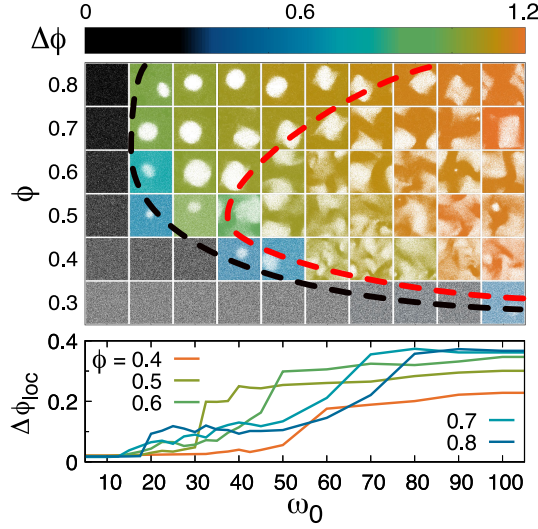


FIG. 2. A collection of representative configurations, sampled in the late time regime, in the ω_0 - ϕ plane. Particles are colored according to the dispersion in local density, $\Delta\phi = \phi_{\text{loc}}^{\text{max}} - \phi_{\text{loc}}^{\text{min}}$. The bottom panel displays the diameter of the gaseous voids for stationary phase separation (d_v), as a function of ω_0 , for different values of ϕ . The bottom panel displays the root mean square time deviation $\Delta\phi_{\text{loc}}$ of the local surface fraction. At fixed ϕ , $\Delta\phi_{\text{loc}} \simeq 0$ for small ω_0 , corresponding to a steady homogeneous state with standard fluctuations; then, at intermediate ω_0 , we observe a first small increase of $\Delta\phi_{\text{loc}}$, indicated with the black dashed line in the phase diagram, corresponding to a steady phase separated pattern with the void slowly moving in the system; then a second increase at larger ω_0 , indicated with the red dashed line in the phase diagram, corresponds to a transition to the nonsteady state.

deviation of the local velocity $\Delta\phi_{\text{loc}} = \langle (\phi_{\text{loc}}(\mathbf{r}, t) - \langle \phi(\mathbf{r}) \rangle_t)^2 \rangle_t^{1/2}$ (shown in the bottom panel of Fig. 2) and indicated by a black dashed line in the phase diagram. Above $\omega_{0,c}$, the void size increases monotonically until ω_0 reaches a second critical value $\omega_{0,ns}$, where the void is destabilized in favor of an unsteady regime identified by a second increase of $\Delta\phi_{\text{loc}}$ and indicated by a red dashed line in the phase diagram. Movies of the system dynamics in the different regimes are provided in Appendix A. We observed that, regardless of the values of ω_0 and ϕ , RIPS is suppressed in the absence of large enough translational inertia, see Appendix B. A similar path is followed when ϕ increases at constant ω_0 , with a critical value of ϕ for the appearance of the void. Voids reach a maximum typical finite size, independent of the system size, as discussed in detail in Appendix A.

RIPS voids are circle shaped under the regime of stationary pattern, and there is a counterclockwise edge current over the interface (i.e., opposite to single-particle intrinsic spin). This is shown in the velocity field of Fig. 3(a) and the radial profile of the azimuthal (perpendicular to the radial component) flow velocity u_ϕ and translational temperature T_{kin} of Fig. 3(b). These

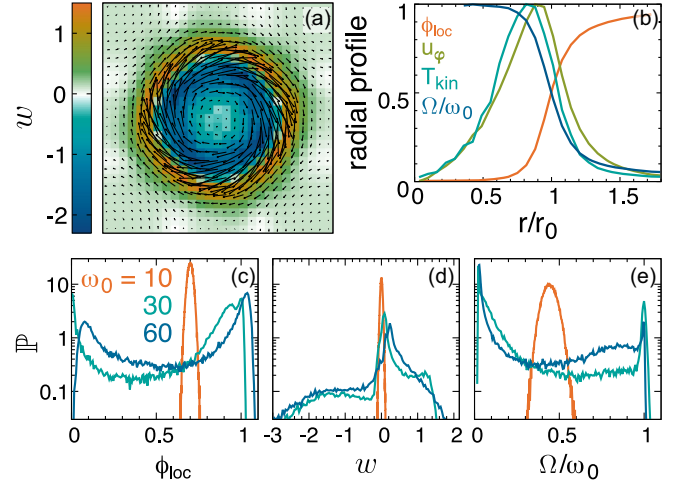


FIG. 3. (a) Velocity field at $\omega_0 = 30$ and $\phi = 0.700$. Colors represents vorticity $w = \partial_y u_x - \partial_x u_y$, where $\mathbf{u}$ is the flow velocity field. (b) Radial profile corresponding to the configuration in (a), computed with respect to the center of the void, of the local density $\bar{\phi}$, the azimuthal component of the flow velocity u_ϕ , the kinetic temperature $2T_{\text{kin}}/m \equiv (\mathbf{v} - \mathbf{u})^2$, and the local average angular velocity field Ω , normalized by ω_0 . The values of u_ϕ and T_{kin} are scaled with respect to their maximum values, $u_\phi^M \approx 50$ and $T_{\text{kin}}^M \approx 300$. In the lower panels are shown the probability density functions of (c) local density, (d) local vorticity, (e) local angular velocity, for $\phi = 0.700$ and $\omega_0 = 10, 30, 80$ in the homogeneous phase, separated phase with a stationary void, nonstationary phase, respectively.

currents generate a vorticity pattern, with a negative (counterclockwise) vorticity inside the void that switches to positive (clockwise) at the interface and then decays in the dense phase. Figure 3(b) displays the characteristic pattern of the local spin $\Omega = \langle \omega_i \rangle_{\text{loc}}$, with $\Omega \simeq \omega_0$ in the interior of the void, where particles are free to move and unlikely to collide. Ω drops in the dense phase, where rotational motion is suppressed by tangential forces over more frequent collisions. These features of the phase separated state can be quantified through the probability density function $\mathbb{P}$ of the local surface fraction ϕ_{loc} , the local vorticity w , and the angular velocity field Ω , as shown in Figs. 3(c)–3(e), respectively, and compared with the case of a homogeneous system at $\omega_0 = 10$, where we see no net vorticity and a normal-distributed spin. In the phase separated regime, we observe a bimodal distribution for ϕ_{loc} and Ω . The maxima for the bimodal distributions correspond to a dense phase with packing fraction peaked close to 1 and spin peaked around zero and a loose phase with packing fraction peaked close to zero and spin peaked around ω_0 . The edge currents produce the concurrent appearance of net positive and negative local vorticity.

To understand the origin of RIPS, we compute the Irving-Kirkwood pressure [52] $P_{\text{IK}} = (1/V) \sum_{i=1}^N [mv_i^2 + \mathbf{r}_i \cdot (\mathbf{F}_i^t + \mathbf{F}_i^n)]$ from a small system that remains

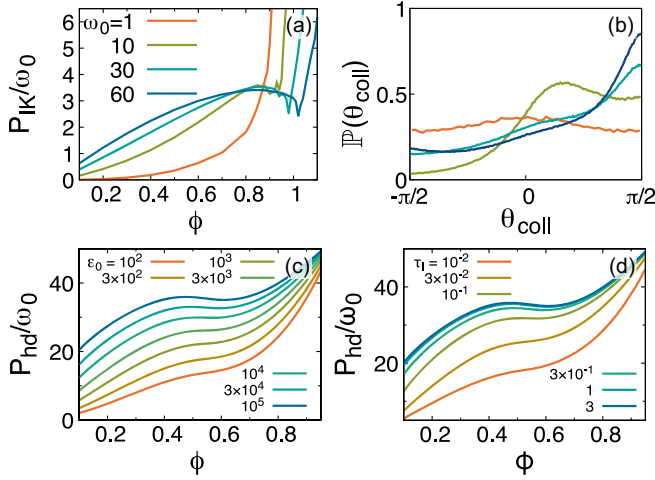


FIG. 4. (a) Numerical pressure P_{IK} for different values of ω_0 . (b) Probability density function $\mathbb{P}$ of the collisional angle, measured from simulations, for $\phi = 0.700$ and the same values of ω_0 as in (a). Hard disk pressure, P_{hd} , computed within the kinetic model for $\alpha = 0.1$, $\tau_R = 10^{-2}$, $\nu = 3$, for (c) $\tau_I = 10$ and increasing values of ϵ_0 and (d) $\epsilon_0 = 10^5$ and increasing τ_I . Values of numerical and effective pressure are scaled by ω_0 for graphical convenience.

homogeneous for all values of ω_0 and ϕ . Figure 4(a) shows that P_{IK} develops, at high enough ω_0 , a nonmonotonic behavior with a regime of negative compressibility. This can lead to an instability that eventually drives phase separation, as also observed in heated granular fluids [18].

In order to gain insight into the instability mechanism, we propose a kinetic model for 2D hard disks, to compare the hard disk pressure, $P_{hd} = 4/(\pi\sigma^2)\phi T(1 + 2\phi\chi_{hd})$ [53], with MD results. As has been proposed in shaken granular gases [18], we use an equilibrium pair distribution function at contact, $\chi_{hd} = (1 - 7\phi/16)/(1 - \phi)^2$, to derive T as an effective temperature resulting from the balance of energy exchange between colliding particles and between particles and the substrate.

Between collisions, particles lose translation energy due to the friction with the substrate at a rate controlled by the particles' inertia, which can be expressed as

$$\langle \Delta E_k^{\text{fric}} \rangle = \left\langle \frac{1}{2} \int_0^{t_f} dt \gamma v_{\text{out}}^2 e^{-2t/\tau_I} \right\rangle = \frac{\langle t_f \rangle}{2\langle t_f \rangle + \tau_I} T, \quad (3)$$

where v_{out} is the outgoing velocity after a collision, and in the averages we have assumed a Maxwell-Boltzmann distribution for the velocities and an exponential free flight time distribution, with $\langle t_f \rangle = \sigma\sqrt{m\pi}/(8\phi\chi_{hd}T^{1/2})$ [54].

The gain of translational kinetic energy of a pair of colliding particles, as prescribed by Eqs. (1) and (2), reads

$$\Delta E_k^c = m\alpha[\Delta v^t((\alpha - 1)\Delta v^t + (1 - 2\alpha)\sigma\omega) + \alpha(\sigma\omega)^2], \quad (4)$$

where $\Delta v^t = \Delta v \sin \theta$ is the incoming relative tangential velocity and θ the collision angle [see Fig. 7 and Appendix C for the definition of θ and the derivation of Eq. (4)]. We have defined $\alpha = \eta\Delta t/m$, with Δt the typical collision time and ω the average incoming angular velocities of the pair. As illustrated in Fig. 1, ω evolves dynamically prior to the collision, as $\omega = 2\omega_0[1 - e^{-t/\tau_R}]$, assuming that ω drops to zero at all collisions. Hence, the average energy exchange at collisions reads

$$\langle \Delta E_k^c \rangle = \frac{8\alpha^2\epsilon_0\langle t_f \rangle^2}{(\tau_R + \langle t_f \rangle)(\tau_R + 2\langle t_f \rangle)} + 2\alpha(\alpha - 1)T\langle \sin^2\theta \rangle_\theta + \alpha(1 - 2\alpha)\sqrt{2\pi\epsilon_0 T} \left(\frac{\langle t_f \rangle}{\tau_R + \langle t_f \rangle} \right) \langle \sin\theta \rangle_\theta, \quad (5)$$

where $\epsilon_0 = m\sigma^2\omega_0^2/4$ stands for the particle rotational energy input, and $\langle \rangle_\theta$ for an average over collision angles.

As interactions are noncentral, inhomogeneous states develop chiral flows with a relevant impact on the collision angle distribution $f(\theta)$. Particles flowing along a void edge are more likely to undergo grazing collisions with those in the bulk, thus with a preferred angle around $\theta = \pi/2$. This effect is captured by

$$f(\theta) = \frac{1}{\mathcal{N}} \left[\frac{\phi^\nu}{\pi} + 2(1 - \phi)^\nu \delta(\theta - \pi/2) \right], \quad (6)$$

with $\theta \in [-\pi/2, \pi/2]$ and $\mathcal{N}$ a normalization constant. This is a generic nonlinear interpolation between a uniform distribution at high density and a δ -peaked one at low density. Enhanced collisions at $\theta = \pi/2$ in the gaseous phase is backed up by MD as shown in Fig. 4(b), which displays a peak at $\pi/2$ as ω_0 increases and voids start forming.

A steady state is achieved when Eqs. (3) and (5) balance each other, determining the translational kinetic temperature, T (the expression can be found in Appendix C), which results a monotonically decreasing function of ϕ .

In agreement with MD results, Figs. 4(c) and 4(d) show a transition from an increasing pressure to a nonmonotonic one with respect to ϵ_0 and/or τ_I . This can be interpreted as a competition among the relevant timescales at play. Finite rotational inertia makes the rotational to translational kinetic energy exchange increasingly concealed for increasing density since, as $\tau_R > t_f$, active rotation is not fully recovered in between collisions, and hence the fluid becomes less chiral. This explains why the temperature decreases with density, thus yielding a decreasing pressure. τ_I decreases upon increasing friction with the embedding medium, causing the energy gained in collisions to be quickly dissipated, cooling the system at low density. The last crucial ingredient for the appearance of a van der Waals loop is the form of $f(\theta)$ in Eq. (6), which accounts for the effect of chiral currents, and modulates the dependence on

density of the energy contribution in Eq. (5). Further details on the dependence of the pressure on the parameter ν in Eq. (6) are reported in Appendix D. This means that when a density gradient is created in the system due to spontaneous fluctuations, the emergence of the circulating currents favors an extra energy input through the selection of preferred collision angles, corresponding to grazing collisions, which triggers the mechanical instability responsible for phase separation.

In summary, we have shown that the RIPS develops as a result of the pressure that emerges from the competition between active rotation and interparticle friction. Rotation transfers linear momentum more efficiently at low densities, while contact friction inhibits particles' persistent rotation in the denser region. As a result, a positive feedback mechanism is induced, which increases pressure in low-density regions and reduces pressure in high-density regions. Consequently, particles are expelled out of the void region toward the denser region. This is reflected in the emergence of a van der Waals loop, as shown in Fig. 4. We observe that RIPS is characterized by voids of finite size, as they grow until they become so rarefied that particle encounters are not frequent enough, angular to linear momentum transfer becomes inefficient, and the system reaches a stationary state. A more detailed description of the mechanical balance between the two coexisting phases, including the odd components of the stress that generates the circulating currents [4], is required to predict the void size. Further analysis of interface stability in the presence of edge currents would also clarify the transition to the unsteady regime, observed for large spinning velocity, which is not predicted by our kinetic theory.

The proposed model has the minimal features that explain the emerging mechanical instability leading to the phase separation observed for inertial spinners. A similar phenomenon has been observed in systems that interact through noncentral forces and generate chiral flows [21,22]. Therefore, RIPS is a general phenomenon associated with a large class of systems that share a few simple features, such as short range tangent forces, and translational inertia.

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Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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End Matter

Appendix A: Description of the movies and finite-size effects—The Supplemental Movie 1 [55] shows the late time dynamics of the system at global surface fraction $\phi = 0.700$, for three different values of the persistent angular velocity, namely $\omega_0 = 10$, in the homogeneous phase, $\omega_0 = 30$, in the phase separated stationary phase, and $\omega_0 = 80$, in the phase separated nonstationary phase (see the phase diagram in Fig. 2 of the main text). Particles are colored by their instantaneous speed.

As mentioned in the main text, the typical size of the voids in the RIPS regime is finite, reaching a characteristic maximum value for each system configuration. This implies that, simulating a larger system, compared to the one seen in the main text, it shows that RIPS is characterized by arrested phase separation, forming an array of voids of approximately the same size. The Supplemental Movie 2 [55] shows that this configuration is stationary and the voids do not coarsen, while generating the same edge current pattern as in the smaller system.

The Supplemental Movie 3 [55] shows the dynamics of the system in the chaotic regime, for $\omega_0 = 80$ and $\phi = 0.600$ and large system size, displaying the same unsteady state as the one reported in the phase diagram of Fig. 2.

Appendix B: Suppression of phase separation for small inertia and analysis of the inertial timescales—As discussed in the main text, phase separation is observed only in the regime of large particles' translational inertia τ_I . Figure 5 displays a set of configurations for decreasing τ_I (increasing translational damping coefficient γ) at constant $\omega_0 = 30$ and $\phi = 0.700$, showing how the void continuously shrinks and phase separation disappears. As we showed with our effective model, this behavior stems from a balance of energy injected in the system at collisions between particles and lost through friction, controlled by two relevant timescales: the average free flight t_f enslaved to density, and τ_I . Figure 5 shows that τ_I directly controls the size of the void in the phase separated state.

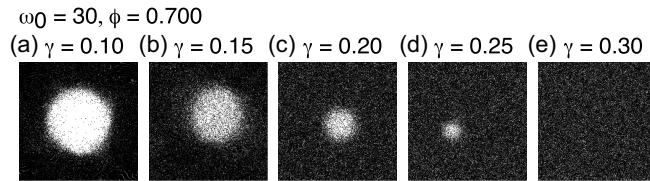


FIG. 5. Representative configurations of the system, sampled within the late time stationary regime, $\omega_0 = 30$, $\phi = 0.700$ and increasing values of γ . Starting from (a) $\gamma = 0.1$, the value used for all the cases shown in the main text, and increasing γ , we see the phase separation progressively reenter (b)–(d), until it is suppressed at around $\gamma = 0.30$ (e).

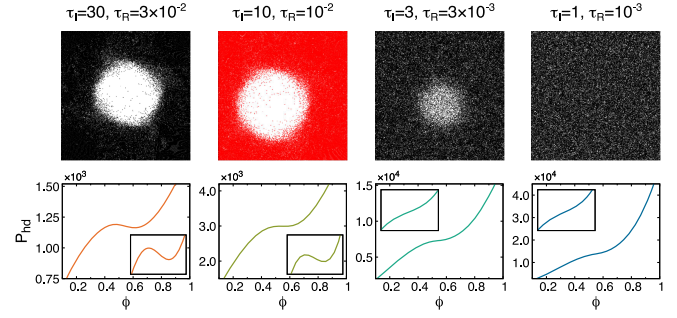


FIG. 6. In the top row are snapshots of the steady state of the system, for $\omega_0 = 30$, $\phi = 0.700$, $\tau_I = 30, 10, 3, 1$, respectively, and $\tau_R = 10^{-3}\tau_I$, such that $\tau_R \ll \tau_I$. The red snapshot is the one corresponding to the values of τ_I and τ_R used in the main text, for the phase diagram shown in Fig. 2. In the bottom row is shown the equation of state for the hard disk system, as obtained from the kinetic theory developed in the manuscript.

We also investigate the nature of the system as τ_I decreases, while $\tau_R \ll \tau_I$. Figure 6 (top row) shows a series of snapshots of the microscopic system, for decreasing values of τ_I , while keeping constant the ratio $\tau_I/\tau_R = 10^3$. Decreasing both translational and rotational inertia in the microscopic model has the effect of gradually shrinking the void, which eventually becomes unstable, in favor of a homogeneous state. The bottom row of Fig. 6 shows the corresponding behavior of our kinetic theory. The van der Waals loop in the pressure predicted theoretically progressively fattens and then disappears for small inertia, indicating the absence of mechanical instability of the homogeneous state, which is compatible with the observations from the molecular simulations.

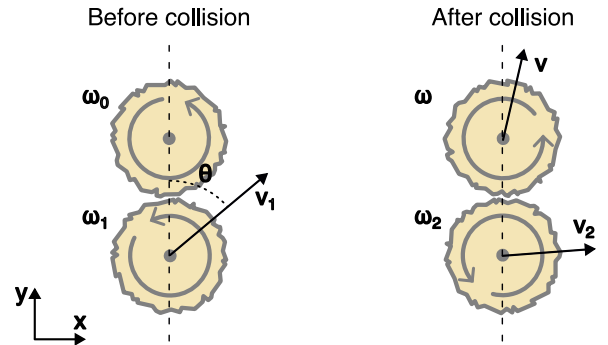


FIG. 7. Diagram of a general collision between two particles, in the reference frame of the top particle. The two particles collide at a generic angle θ , with incoming angular velocities ω_0 and ω_1 . After the collision, both the translational and rotational velocities change, according to the interaction between the two particles, which is an overlap between a repulsive WCA potential and the transversal force, as described the main text.

TABLE I. Coefficients of the polynomial expression for $x = T^2$, as it results in the kinetic model.

x^0	$4\alpha^2\varepsilon_0(m\pi)^{3/2}\sigma^3$
x^1	$16\alpha^2\varepsilon_0m\pi\sigma^2\phi\chi\tau_I + \sqrt{2\pi}(m\pi)^{3/2}\sigma^3\alpha(1-2\alpha)\sqrt{\varepsilon_0}\langle\sin\theta\rangle_\theta$
x^2	$4\alpha\sqrt{2\pi}m\pi\sigma^2\phi\chi(1-2\alpha)(\tau_I + \tau_R)\sqrt{\varepsilon_0}\langle\sin\theta\rangle_\theta + 2\alpha(\alpha-1)(m\pi)^{3/2}\sigma^3\langle\sin^2\theta\rangle_\theta - (m\pi)^{3/2}\sigma^3$
x^3	$16\pi\sqrt{2\pi}\sigma\alpha\phi^2\chi^2\tau_R\tau_I\sqrt{\varepsilon_0}\langle\sin\theta\rangle_\theta + 16m\pi\sigma^2\alpha(\alpha-1)(\tau_I + 3\tau_R)\phi\chi\langle\sin^2\theta\rangle_\theta - 6m\pi\sigma^2\phi\chi\tau_R$
x^4	$32\sigma\sqrt{m\pi}\alpha(\alpha-1)\phi^2\chi^2\tau_R(3\tau_I + 2\tau_R)\langle\sin^2\theta\rangle_\theta - 16\sigma\sqrt{m\pi}\alpha\phi^2\chi^2\tau_R^2$
x^5	$256\alpha(\alpha-1)\phi^3\chi^3\tau_R^2\tau_I\langle\sin^2\theta\rangle_\theta$

Appendix C: Energy exchange at collisions—Because of the tangential force acting between the particles at collision, a conversion of rotational kinetic energy into translational kinetic energy takes place when two particles collide, resulting in a net translational kinetic energy gain that depends on the collision details.

According to the diagram of a generic collision shown in Fig. 7, the transversal force acting along the x axis is $F_x^t = -\eta[v_{1,x} - R(\omega_0 + \omega_1)]$. The interaction along the y axis is conservative, hence it does not produce energy variations.

Outgoing velocities are therefore

$$v_{2,x} - v_{1,x} = \int dt \frac{F_x^t}{m} \simeq -\frac{\eta\Delta t}{m} \left[v_{1,x} - \frac{\sigma}{2}(\omega_0 + \omega_1) \right],$$

$$v_x = -\int dt \frac{F_x^t}{m} \simeq \frac{\eta\Delta t}{m} \left[v_{1,x} - \frac{\sigma}{2}(\omega_0 + \omega_1) \right], \quad (\text{C1})$$

where K is related to the particles' moment of inertia, $I = Km\sigma^2/4$, and the integral is intended over the duration of the collision Δt . The integral is evaluated assuming a very short collision duration, such that the force can be considered constant. One can then compute the outgoing particle velocities, from which Eq. (4) in the main text is derived.

Balancing the average rate of energy gain at collisions $\langle\Delta E_k^c\rangle$ and the average rate of energy loss due to friction $\langle\Delta E_k^{\text{fric}}\rangle$, we obtain an equation of the form $W(\phi, T^2) = 0$, where ϕ is the hard disk surface fraction and T the kinetic temperature. At fixed density ϕ , W is a 5th order polynomial function of the variable $x = T^2$, with coefficients with coefficients reported in Table I. This equation is solved numerically using the Newton-Raphson method. We observe that only one positive solution is present, which represents the stationary temperature of the kinetic model.

Appendix D: Model phase diagram—In Fig. 8, we show the spinodal line computed from the hard disk

equation of state as derived from the kinetic model, alongside the phase diagram of the microscopic model, as reported in Fig. 2 of the main text. The orange curve in Fig. 8(a) represents the value of the hard-disk surface fraction ϕ at which we observe the slope of the equation of state $P_{\text{hd}}(\phi)$ becoming negative, as a function of $\omega_0 = 2\sqrt{\varepsilon_0}$. These points represent the beginning of the van der Waals loop identified in the kinetic theory, corresponding to the boundary of the spinodal region, highlighted in orange. For small $\omega_0 \lesssim 20$, the equation of state does not show a loop and the pressure is a strictly increasing function of ϕ . The two phase diagrams are compatible, showing that the kinetic theory is able to capture all the main features of this phase separation phenomenon. However, the values of spinodal densities predicted by the kinetic theory differ slightly from the ones observed in the microscopic model.

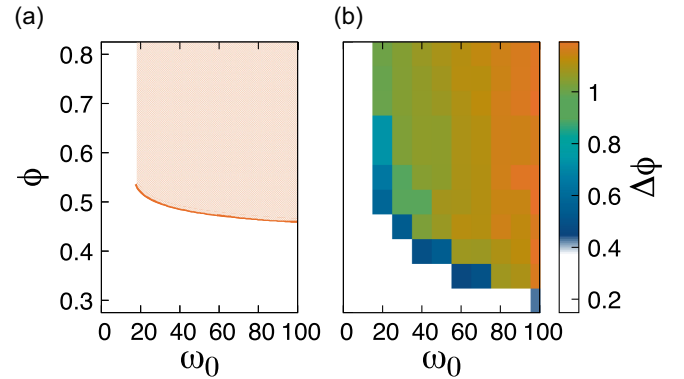


FIG. 8. (a) Spinodal line (orange), corresponding to the value of density where we observe the slope of hard disk pressure $P_{\text{hd}}(\phi)$ becoming negative, as a function of ω_0 , in the kinetic model. The light-orange region corresponds to a phase-separated region. (b) The phase diagram of the continuous theory, as reported in Fig. 2 of the main text.