

Tuning the viscosity and jamming point in dense active non-Brownian suspensions

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Using particle-based simulations, we study the rheology of dense, frictional, non-Brownian suspensions containing active particles. We find that the viscosity and the shear jamming point of the suspension can be tuned by varying the activity strength of the active particles. Microscopically, variations in the number of frictional contacts and the formation of a more isotropic force-chain network due to the presence of active particles underlie the viscosity changes. We introduce an effective temperature, defined through the control parameters of the active particles, that uniquely governs the viscosity of the suspensions. We also find that while activity drives the system away from jamming by hindering force-chain formation, unjamming an already jammed state by breaking existing force chains requires higher activity strength.

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Introduction. Dense suspensions such as slurries, blood, and molten chocolate are ubiquitous in daily life as well as in industry [1–3]. These materials consist of two components: a suspending viscous fluid with viscosity η_f and solid particles that have the same density ρ as the fluid. The viscosity of such systems is controlled by various properties of the solid particles and exhibits a complex nonlinear growth with increasing solid volume fraction ϕ due to proliferation of mechanical contacts. This process eventually leads to a divergence of viscosity at a property-dependent critical solid volume fraction ϕ_J . Such an increase in viscosity, however, poses a challenge in various industrial processes, as the flow of the suspensions becomes increasingly sluggish and energy consuming.

Recently, researchers have focused on the problem of tuning the viscosity of dense suspensions using various methods. One common approach involves altering the properties of the solid particles by varying their surface smoothness, asphericity, or by adding lubricating chemicals to the system [4–7]. However, these methods are not practical for implementation during the flow of a preprepared suspension. Another approach involves tuning the viscosity by applying external mechanical perturbations [8–10] during flow in a fixed direction, or by appropriately changing the flow direction itself [11–13]. Microscopically, such external perturbations reduce the number of frictional contacts. Recent studies suggest that cyclic deformation applied orthogonally to the primary flow can drive the system toward an absorbing state transition characterized by a vanishing number of particle contacts [14], whereas more general orthogonal perturbations may not

necessarily lead to a transition but can still break existing contacts and prevent the formation of fragile force chains [15].

In the context of viscosity reduction, an interesting class of systems is dense suspensions composed of active particles [16,17]. By active particles, we refer to particles that exhibit self-propelled motion by dissipating their internal energy. Such systems are prevalent in biology, with a common example being various microswimmers suspended in a viscous fluid, such as bacterial suspensions [18]. Recently, artificially prepared active particles such as Janus particles or Quincke rollers [19–26] have emerged as model systems to study the rheology of active dense suspensions. The rheology of microswimmer suspensions has been extensively studied at low volume fractions, showing a significant reduction in viscosity of the original fluid due to the presence of active particles [18,27–31]. From the perspective of numerical studies, most have been limited to Brownian soft particles or far from ϕ_J [32–35]. However, how activity-induced dynamics affect the viscosity of non-Brownian suspensions near ϕ_J remains poorly understood. Moreover, recent studies suggest that the presence of active particles and the application of external cyclic deformation have similar effects on the rheology of the system, where active dynamics can act as a form of local driving [36,37]. In the same spirit, one should expect a similar reduction in viscosity observed in Ref. [10] due to the presence of active particles, in a similar context of microrheology [38–40].

In this Letter, using DEM simulations, we study the rheology of active dense suspensions of non-Brownian frictional particles. The active particles are modeled as non-Brownian run-and-tumble particles [41]. We find that the relative viscosity η_r exhibits a nonmonotonic dependence on activity strength: it first decreases and then increases with increasing activity strength. This occurs because weak activity hinders the formation of an anisotropic force-chain network, thereby reducing frictional contacts, whereas strong activity enhances frictional contacts through activity-driven diffusion.

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This change in frictional contacts with activity strength allows tuning of the jamming volume fraction. We phenomenologically introduce an effective temperature that governs the viscosity of active dense suspensions.

Simulation details. We simulate a bidisperse system of N [$\sim \mathcal{O}(10^3)$] frictional granular particles with radii a and $1.4a$ in a 50:50 number ratio, contained within a cubic simulation box of length L , with periodic boundary conditions applied [42]. The frictional particles with particle friction coefficient μ_p are suspended in a density-matched fluid. The system is deformed externally using a space-dependent streaming velocity of the suspending fluid $\mathbf{U}^\infty(\mathbf{r})$. A solid particle experiences three different types of interactions with its surroundings. Due to the relative motion of a particle with respect to the streaming fluid, a solid particle experiences a drag force modeled as $\mathbf{f}_i^d = 6\pi\eta_f a_i [\mathbf{U}^\infty(\mathbf{r}) - \mathbf{u}_i]$ and a torque modeled as $\boldsymbol{\tau}_i^d = 8\pi\eta_f a_i^3 [\boldsymbol{\Omega}^\infty(\mathbf{r}) - \boldsymbol{\omega}_i]$. Here, a_i , $\mathbf{u}_i$, and $\boldsymbol{\omega}_i$ are the radius, linear velocity, and angular velocity of the i th particle, and $\boldsymbol{\Omega}^\infty = \frac{1}{2}(\nabla \times \mathbf{U}^\infty)$. In addition, the presence of the viscous fluid resists the relative motion of a pair of particles, which is modeled here as a short-range hydrodynamic lubrication force [43–45]. The leading-order contribution to the lubrication force between a pair of particles is given by [46] $\mathbf{f}_{ij}^h \sim \frac{1}{h_{ij}}(\mathbf{u}_{ij} \cdot \mathbf{n}_{ij})\mathbf{n}_{ij}$. Here, $\mathbf{u}_{ij} = \mathbf{u}_j - \mathbf{u}_i$ is the relative velocity, $h_{ij} = (a_i + a_j) - |\mathbf{r}_{ij}|$ is the shortest distance between the surfaces of two particles with $\mathbf{r}_{ij}$ the vector pointing from the center of particle i to that of particle j , and $\mathbf{n}_{ij}$ is the corresponding unit vector, $\mathbf{n}_{ij} = \mathbf{r}_{ij}/|\mathbf{r}_{ij}|$. The leading-order term of the torque on the i th particle due to the interaction with the j th particle is given by $\boldsymbol{\tau}_{ij}^h \sim \ln(\frac{a_i}{h_{ij}})(\mathbf{u}_{ij} \times \mathbf{n}_{ij})$. For computational advantage, we compute this force only for $h_{ij} < 0.5a'$ and keep $h_{ij} = 0.001a'$ for $h_{ij} \leq 0.001a'$ to allow the particles to come in direct contact. Here, a' is the radius of the smaller particle. When a pair of particles have finite overlap ($h_{ij} < 0$), they experience frictional contact force modeled as Hookean interaction with normal component $\mathbf{f}_{ij}^{c,n} = -k_n h_{i,j} \mathbf{n}_{ij}$ and tangential component $\mathbf{f}_{ij}^{c,t} = -k_t \boldsymbol{\xi}_{ij}$ and torque $\boldsymbol{\tau}_{ij}^c = a_i(\mathbf{n}_{ij} \times \mathbf{f}_{ij}^{c,t})$. Here, k_n and k_t are the normal and tangential stiffnesses. $\boldsymbol{\xi}_{i,j}$ denotes the accumulated tangential displacement between particles, measured from the moment they first come into contact until the contact breaks. This displacement accounts for the history dependence of the frictional force [47]. According to Coulomb's criterion, the maximum allowable tangential force for frictional particles is given by $k_t \boldsymbol{\xi}_{i,j} / k_n h_{i,j} < \mu_p$. We choose $\mathbf{U}^\infty(\mathbf{y}) = \dot{\gamma} \mathbf{y} \hat{\mathbf{x}}$, where $\mathbf{y}$ is the direction of the velocity gradient and $\hat{\mathbf{x}}$ is the unit vector along the shear direction. We deform the system by keeping $\dot{\gamma}$ and ϕ fixed, and measure different components of stress.

We randomly select a fraction C_a of particles from the presheared state and designate them as active particles. In addition to contact and hydrodynamic forces, each active particle experiences an additional active or self-propulsion force $\mathbf{f}_{a,i} = F_a \mathbf{n}_i$. Here, $\mathbf{n}_i$ is a unit vector oriented randomly with direction changing independently after a persistence time τ_p with $\sum_{i=1}^{N C_a} \mathbf{n}_i = 0$ at any time of the simulation. This induces self-propulsion in the viscous fluid in the absence of interactions. The stress reported here is the virial stress given by $\sigma_{\alpha\beta} = \frac{1}{L^3} \sum_{i \neq j} \mathbf{f}_{a,i,j} \cdot \mathbf{r}_{\beta,i,j}$. For the passive system, the shear and

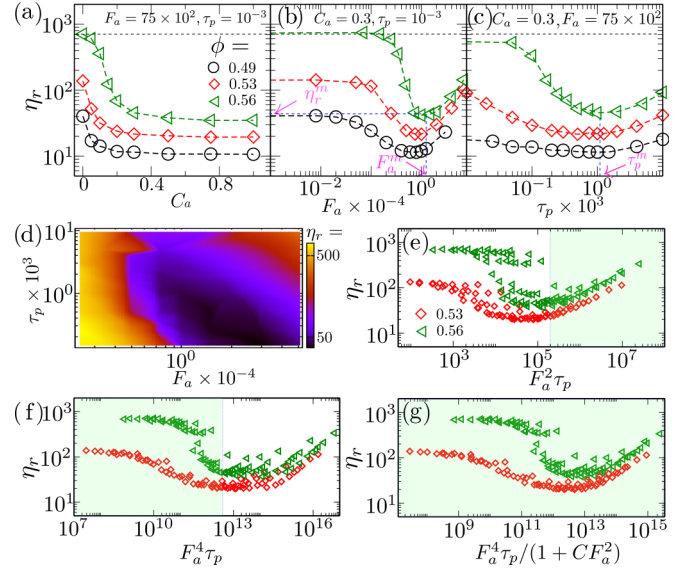


FIG. 1. Variation of relative viscosity η_r as a function of (a) active particle fraction C_a , (b) active force F_a , and (c) persistence time τ_p , for three different volume fractions ϕ . The black dashed line represents the viscosity of the passive system with $\phi = 0.56$. Blue dashed lines show the coordinate of the minimum of the η_r - F_a or τ_p curve. (d) Variation of η_r with F_a and τ_p at fixed $C_a = 0.3$ is shown for volume fractions $\phi = 0.56$. (e)–(g) η_r expressed as a function of an effective temperature T_{eff} constructed by different combinations F_a and τ_p that collapse data at different regimes (green region) for $C_a = 0.3$. The legends in (f) and (g) are the same as (e).

normal stress σ_{xy} comprise three contributions: contact force, lubrication force, and drag force. In active systems, there is an additional active stress given by $\sigma_{a,\alpha\beta} = \frac{1}{L^3} \sum_{i=0}^{N C_a} F_{a,\alpha,i} r_{\beta,i}$. However, in our simulations, for the chosen range of active particle parameters, the active stress components remain significantly smaller than the components arising from particle interactions [48,49].

We choose the units of length, time, and viscosity as $a^* = a = 1$, $t^* = 1/\dot{\gamma} = 10^3$, and $\eta^* = \eta_f = 0.1$, respectively. This choice sets the force unit to $f^* = \dot{\gamma} \eta_f a^2 = 10^{-4}$, the stress unit to $\sigma^* = \dot{\gamma} \eta_f = 10^{-4}$, and the density unit to $\rho^* = \eta_f / (a^2 \dot{\gamma})$. We choose $\mu_p = 1$ and $\rho = 10^{-2}$ in units of ρ^* , and $k_n = 10^8$ and $k_t = 7 \times 10^7$ in units of $\sigma^* a^*$. To ensure that the system without activity remains in the rate-independent regime, we choose the strain rate $\dot{\gamma}$ such that $\dot{\gamma} \ll \eta_f / (\rho a^2)$, $\dot{\gamma} \ll \sqrt{k_n / (\rho a^3)}$, and $\dot{\gamma} \leq \tau_p^{-1}$ [50,51]. F_a and τ_p are expressed in units of f^* and t^* throughout the paper. All data presented here are for systems with 10^3 particles and $\mu_p = 1$, unless stated otherwise.

Tunable viscosity using active particles. First, we focus on how $\eta_r = \sigma_{xy} / \dot{\gamma} \eta_f$ depends on different control parameters of active particles by running simulations for varying values of one parameter while keeping the other two parameters fixed (Fig. 1). Here, σ_{xy} is shear stress. We find that η_r decreases monotonically with C_a for all studied ϕ values, but exhibits a nonmonotonic dependence on F_a and τ_p [shown in Figs. 1(a)–1(c)]. In the case of dependence on F_a , η_r remains more or less constant at small F_a , then decreases with F_a to a ϕ -dependent minimum viscosity η_r^m (shown by blue dashed line) at a

ϕ -dependent active force F_a^m . This suggests that when the active force is too small, it is not able to break the force chains or prevent their formation and does not contribute significantly to particle diffusion. In this regime, the system behaves like a passive system. Only when F_a exceeds a critical value does the activity affect the rheology of the system. When τ_p is varied, a similar response is observed with minimum viscosity at τ_p^m , except with no clear plateau at small τ_p . When $F_a > F_a^m$ and $\tau_p > \tau_p^m$, the trend reverses, with η_r increasing with both F_a and τ_p . The viscosity reduction is more than an order of magnitude for the highest ϕ , comparable to reported values in Refs. [12,14]. Such a nonmonotonic dependence of viscosity on the activity strength is in agreement with the rheological response of dry granular systems reported in Ref. [52]. However, in the case of soft particles, the relaxation time always decreases with increasing F_a [41,53], while a nonmonotonic dependence on τ_p has been observed in theoretical studies [54,55] as well as in simulations [53,56,57].

In Fig. 1(d), the color map of η_r is shown as a function of F_a and τ_p for $\phi = 0.56$, with a fixed $C_a = 0.3$. Here, F_a is chosen larger than the critical value below which no variation of η_r is observed. It is clear that when F_a is small, η_r is weakly dependent on τ_p within the studied range. However, as F_a increases, the effect of τ_p becomes prominent, with F_a^m showing strong dependence on τ_p . This observation suggests that although both F_a and τ_p control η_r , they do not have equal weight.

To understand with what weight F_a and τ_p contribute to η_r , we introduce an effective temperature T_{eff} in our non-Brownian active dense suspensions. This effective temperature originates from the motion of particles due to self-propulsion. In the case of active systems of run-and-tumble Brownian particles such T_{eff} is usually defined as $T_{\text{eff}} \sim F_a^2 \tau_p$ [19,41,58,59]. If this definition holds in our case, we should observe the same value of η_r for simulations with different combinations of F_a and τ_p that yield the same T_{eff} . In Fig. 1(e), η_r vs $F_a^2 \tau_p$ is shown for a fixed value of the latter using multiple combinations of τ_p and F_a . Although this definition of T_{eff} appears to work at large T_{eff} , it fails to collapse data at smaller values. This is expected since at small F_a , η_r is weakly dependent on τ_p within the studied range. To account for this, we modify the definition of the effective temperature as $T_{\text{eff}} \sim F_a^4 \tau_p$. The dependence of η_r on $F_a^4 \tau_p$ is shown in Fig. 1(f), which successfully collapses the data in the small-value regime. However, the microscopic origin of this form of effective temperature is not studied here. Finally, using the harmonic combination of these two effective temperatures, we define a $T_{\text{eff}} \sim F_a^4 \tau_p / (1 + CF_a^2)$ which can collapse data for all activity strength as shown in Fig. 1(g). Here, C is a constant with dimensions of inverse force squared, having a value of 10^{-8} .

Data presented in Fig. 1 suggest that the most important parameter controlling the viscosity is the active force. We further explore, for fixed C_a and τ_p , how F_a^m and η_r^m depend on ϕ . Figure 2(a) shows the dependence of η_r on F_a for $C_a = 1$ and $\tau_p = 10^{-3}$ over a range of ϕ values spanning above and below $\phi_J^{\mu_p=1}$. Both F_a^m and η_r^m increase with ϕ and diverge close to isotropic jamming volume fraction ϕ_{RCP} [60] as shown in Fig. 2(b), where green and red solid lines represent the functional forms $F_a^m(\phi) = \mathcal{F}_1(\phi) \sim (\phi_J - \phi)^{-\alpha}$ and

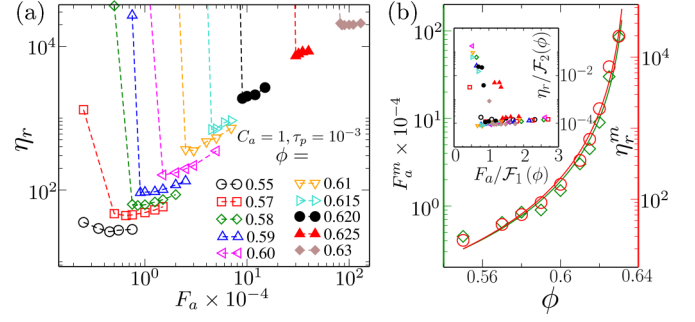


FIG. 2. (a) Relative viscosity η_r as a function of active force F_a for different volume fractions ϕ above and below $\phi_J^{\mu_p=1}$. (b) The divergence of active force F_a^m that yields minimum viscosity η_r^m , and minimum viscosity η_r^m as $\phi \rightarrow \phi_{\text{RCP}}$. Green and red solid lines stand for the scaling functions $\mathcal{F}_1(\phi)$ and $\mathcal{F}_2(\phi)$. Inset: Collapse of data shown in (a) with the help of $\mathcal{F}_1(\phi)$ and $\mathcal{F}_2(\phi)$.

$\eta_r^m(\phi) = \mathcal{F}_2(\phi) \sim (\phi_J - \phi)^{-\beta}$, respectively, with $\alpha = 2.04$, $\beta = 2.64$, and $\phi_J = 0.636$. The inset of Fig. 2(b) shows the collapse of data presented in Fig. 2(a) by scaling F_a and η_r using $\mathcal{F}_1(\phi)$ and $\mathcal{F}_2(\phi)$. Similar dependence on τ_p is also expected, but not studied here.

Microscopic understanding of tunable viscosity. To understand the microscopic origin of tunable viscosity due to self-propulsion, we study the dependence of the mean number of frictional contacts $n_{\text{fc}} = n_{\text{fc}}^{\text{total}}/N$ and force chain network on different control parameters of active particles. Here, $n_{\text{fc}}^{\text{total}}$ is the total number of frictional contacts in the system. In Fig. 3(a), the variation of frictional contacts with active force (bottom, in red, for $\tau_p = 10^{-3}$) and with persistence time (top, in green, for $F_a = 75 \times 10^2$) is shown for $\phi = 0.56$ and

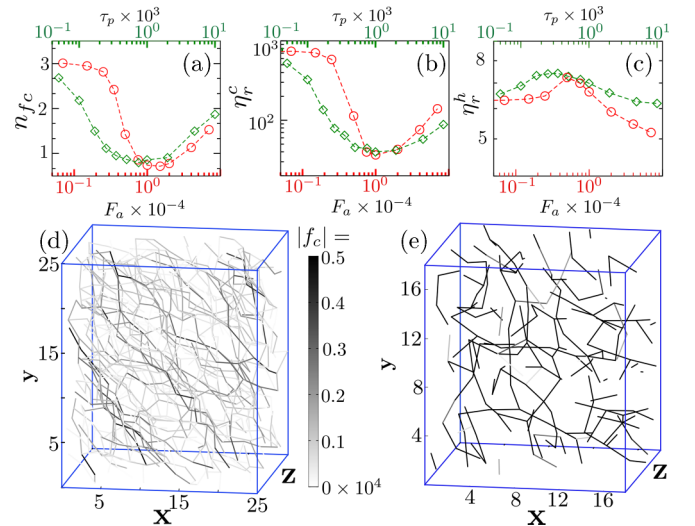


FIG. 3. (a) Variation of the mean number of frictional contacts n_{fc} with active force F_a (red) and persistence time τ_p (green). (b) Variation of the viscosity contributions from frictional contacts η_{fc}^c with F_a (red) and τ_p (green). (c) Same as (b) for the hydrodynamic interactions. (d) Force chain network for a passive system with $\phi = 0.49$. (e) Same as (d) for an active system with $\phi = 0.56$, $C_a = 0.3$, $F_a = 75 \times 10^2$, and $\tau_p = 10^{-3}$, which has the same η_r as the passive system.

$C_a = 0.3$ in both cases. The trend is similar to that observed in Figs. 1(b) and 1(c), suggesting that the number of frictional contacts varies depending on the magnitude of F_a and τ_p , which in turn leads to changes in η_r . In the case of small F_a , the number of frictional contacts remains unchanged, indicating that the activity is not sufficient to break existing contacts or prevent their formation. However, as F_a increases, contacts begin to break or fail to form. At large values, the number of contacts starts increasing again due to enhanced particle diffusion driven by strong active forcing. The dependence on τ_p is similar except for the plateau observed at small F_a . In Fig. 3(b), the viscosity contributions from contact forces η_r^c and hydrodynamic forces η_r^h are shown. As frictional contacts break or fail to form, η_r^c decreases, leading to an increase in purely hydrodynamic interactions and thus a higher η_r^h .

We further explore the force chain network [Figs. 3(d) and 3(e)] to illustrate the microscopic differences between passive and active systems at different ϕ but with the same η_r . Force chains are represented by lines connecting the centers of contacting particles, with the color of each line indicating the magnitude of the contact force. For clarity, the same color scale is used in all panels, with force magnitudes capped at a maximum of 5×10^3 . When the force exceeds 5×10^3 , it is truncated to this upper limit. We observe that the passive system with $\phi = 0.49$ exhibits long, branched, and anisotropic force chains. In contrast, the active system with $\phi = 0.56$, $C_a = 0.3$, $F_a = 75 \times 10^2$, and $\tau_p = 5 \times 10^{-4}$, having the same η_r as the passive system, displays shorter, more isotropic chains, although some contacts exhibit significantly higher force magnitudes compared to others. This suggests that activity reduces the anisotropy in the structure that originates from straining.

Tuning the jamming volume fraction. The systematic reduction of contact numbers with increasing activity strength suggests that it is possible to tune the jamming volume fraction ϕ_J with the help of active particles. Note that, since shear jamming depends on various factors such as preparation protocol, stress anisotropy, friction, etc., by ϕ_J we mean effective jamming volume fraction under steady shear in the presence of activity. Shear jamming occurs due to anisotropy in the microscopic structure, which leads to $n_{fc} \geq 4$. However, active particles introduce random local straining in addition to global straining, which either breaks the force chains or prevents their formation and reduces structural anisotropy. In Fig. 4(a), the data for the η_r - ϕ curve for a passive system with $\mu_p = 1$ are shown by black circles. For this system, shear jamming occurs at $\phi_J^{\mu_p=1} \sim 0.58$. Here, we keep $\tau_p = 10^{-3}$ for all active systems. The data points corresponding to $F_a = 25 \times 10^2$ and 5×10^3 for $C_a = 0.3$ introduce mild activity into the system. These systems attain a low viscosity at small ϕ , but as ϕ increases, the effect of the active forcing becomes insignificant, and the η_r - γ curves start to approach that of the passive system, eventually converging with it. With a further increase of F_a to 75×10^2 , keeping C_a the same, the viscosity remains significantly lower than that in the passive system at higher ϕ , but begins to increase dramatically as ϕ approaches the jamming volume fraction. However, jamming occurs at $\phi > \phi_J^{\mu_p=1}$. Increasing F_a further to 3×10^4 shifts the jamming volume fraction close to 0.60. The data points for $F_a = 3 \times 10^4$ and 9×10^5 with $C_a = 1$ suggest it is possible to

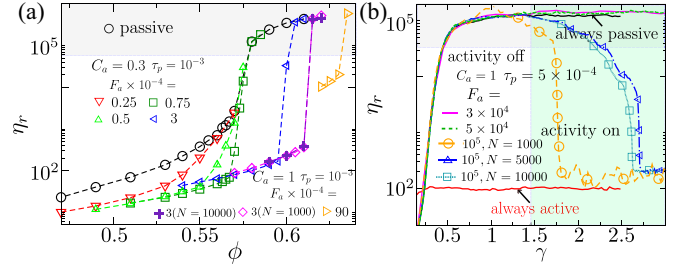


FIG. 4. (a) Relative viscosity η_r plotted as a function of volume fraction ϕ for various activity strengths, for $\mu_p = 1$. Data suggest the effective jamming volume fraction under steady shear in the presence of activity is pushed toward ϕ_{RCP} . (b) η_r - γ curves with activity turned off for $\gamma \leq 1.4$ and turned on for $\gamma > 1.4$ (green region). The black and red lines represent the η_r - γ curves for systems without activity and with activity always on. All other curves are with activity off for $\gamma \leq 1.4$ and on for $\gamma > 1.4$ for two different activity strengths.

push $\phi_J^{\mu_p=1}$ to ϕ_{RCP} . Also, data for two different system sizes, $N = 10^3$ (magenta diamonds) and $N = 10^4$ (violet pluses), at the same activity parameters suggest that there is no finite-size effect in the jamming mechanism for the presented data. Interestingly, jamming transition under steady shear in the presence of activity is different from that in passive systems. With activity, the system remains in a low-viscosity state for $\phi > \phi_J^{\mu_p=1}$ and then jams abruptly within a very narrow window of ϕ [see Fig. 4(a)]. A similar behavior is also observed with varying F_a [see Fig. 2(a)]. Note that with increasing F_a , η_r at smaller ϕ is higher than what is obtained for smaller F_a . This is consistent with the existence of a ϕ -dependent F_a^m , as discussed in Fig. 1. Such an increase in ϕ_J due to active dynamics is also observed in soft point particles [33,61,62].

Jamming-unjamming transition using activity. We further study the jamming-unjamming transition facilitated by activity by first driving the system while keeping activity off for $\gamma \leq 1.4$, where the system gets jammed, and then switching on activity for $\gamma > 1.4$ [green region in Fig. 4(b)] to study unjamming while the system is being sheared. We choose $\phi = 0.58$, which, in the absence of activity, jams upon straining (in the hard-sphere limit) with a very large viscosity, $\eta_r \sim 4 \times 10^4$ (gray region), as shown by the black line in Fig. 4(b). In the presence of activity with $F_a = 3 \times 10^4$, $C_a = 1$, and $\tau_p = 10^{-3}$, the system attains a low-viscosity steady state, shown by the red line. The pink line represents data for the same activity, but the activity is kept off for $\gamma \leq 1.4$, where the system attains a high-viscosity state. Interestingly, when activity is turned on after $\gamma > 1.4$, the system remains jammed. Even when F_a is increased to 5×10^4 (green dotted line), the system remains jammed. The orange line (circles) corresponds to $C_a = 1$, $F_a = 10^5$, and $\tau_p = 5 \times 10^{-4}$ with $N = 10^3$, which unjams once activity is switched on. There is a mild finite-size effect, with larger systems ($N = 5 \times 10^3$, blue triangles, and $N = 10^4$, cyan rectangles) fluidizing at larger strain values. This strongly suggests that if activity is applied from the beginning of deformation, a smaller F_a can lead to fluidization. In contrast, once the system is jammed, a much larger F_a is required to unjam it, highlighting the different roles of active particles in these two distinct scenarios. For unjammed systems, the self-propulsion of active parti-

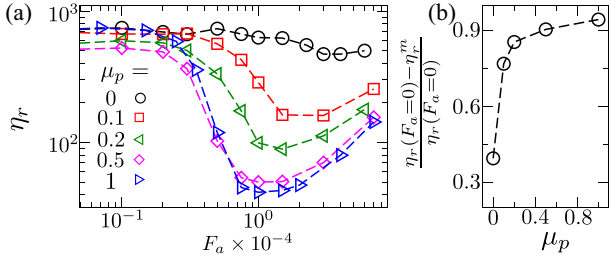


FIG. 5. (a) Variation of relative viscosity η_r as a function of active force F_a for different friction coefficients $\mu_p = 0, 0.1, 0.2, 0.5$, and 1 for volume fractions $\phi = 0.63, 0.61, 0.595, 0.57$, and 0.56 , respectively. (b) Relative viscosity drop defined as $[\eta_r(F_a = 0) - \eta_r^m]/\eta_r(F_a = 0)$ due to self-propulsion as a function of friction coefficient μ_p .

cles prevents force-chain formation, while in already jammed systems, active forcing reduces viscosity by breaking force chains, which requires higher activity strength. Figure 4(b) shows data for one volume fraction, however, the finding remains valid for other $\phi > \phi_j^{\mu_p=1}$ as well.

Role of friction in tunable viscosity. In Fig. 5, we show how η_r varies with F_a for a range of μ_p , with suitable ϕ for different μ_p chosen such that the passive systems have similar viscosities. We find that the trend in the variation of η_r remains the same for all μ_p . However, the relative drop in viscosity, quantified as $\frac{\eta_r(F_a=0) - \eta_r^m}{\eta_r(F_a=0)}$, increases with μ_p . The frictional constraints allow the system to have the same viscosity at lower ϕ . Self-propulsion reduces the effect of these constraints and therefore results in a much lower viscosity, as a lower ϕ allows particles to access a larger amount of available free volume.

Conclusion. To conclude, we find that the presence of active particles in dense suspensions of frictional granular particles significantly tunes the viscosity and the effective jamming volume fraction under steady shear. This suggests that the presence of self-propelled particles in the suspension and the application of mechanical perturbation orthogonal to

the primary flow [14] have similar effects on the rheology; both tune the number of frictional contacts. The viscosity of the suspension is uniquely determined by an effective temperature, which is defined using various control parameters of the active particles. Similar to our study, shear thinning followed by shear thickening with activity strength are also found in systems of soft active Brownian particles. However, the microscopic mechanisms are different. In the latter case, shear thinning occurs due to the breaking of particle clusters [35], or due to elastic barrier or cage breaking [32] induced by active forcing, while thickening occurs due to cluster formation. In the former case, shear thinning occurs due to the prevention of force-chain network formation, and thickening occurs due to activity-induced diffusion, which may also lead to clustering at large τ_p . The effective temperature discussed in this work is constructed purely as a means to achieve the best data collapse. The microscopic origin of such a form is unknown and will be explored in the future. We note that close to $\phi_j^{\mu_p}$, the rheology is dominated by the frictional interactions rather than the hydrodynamic interaction. Therefore, our findings are expected to be applicable for dry granular systems as reported in Ref. [52]. Our findings will be useful for designing experiments to study the rheology of active suspensions. This study is limited to a fixed strain rate and a minimal model of active particles. Effects of self-propulsion on the rate-dependent phenomena such as thickening or thinning [63,64] and active particles modeled in more practical ways [52,65] are left for future studies.

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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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