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## Extremely Persistent Dense Active Fluids

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We study the dynamics of dense three-dimensional systems of active particles for large persistence times  $\tau_p$  at constant average self-propulsion force  $f$ . These systems are fluid counterparts of previously investigated extremely persistent systems, which in the large persistence time limit relax only on the time scale of  $\tau_p$ . We find that many dynamic properties of the systems we study, such as the mean-squared velocity, the self-intermediate scattering function, and the shear-stress correlation function, become  $\tau_p$ -independent in the large persistence time limit. In addition, the large  $\tau_p$  limits of many dynamic properties, such as the mean-square velocity and the relaxation times of the scattering function, and the shear-stress correlation function, depend on  $f$  as power laws with non-trivial exponents. We conjecture that these systems constitute a new class of extremely persistent active systems.

## 1 Introduction

Self-propelled or active particles, which use energy from their environment to perform persistent motion, behave in surprising and interesting ways<sup>1–3</sup>. Recently, novel intermittent dynamics was identified in extremely persistent dense homogeneous two-dimensional systems<sup>4–7</sup>. These systems were shown to relax on the time-scale of the persistence time  $\tau_p$  by going through sequences of mechanical equilibria in which self-propulsion forces balance interparticle forces. In contrast, at low and moderate densities and long persistence times active matter undergo motility-induced phase separation (MIPS), and separate into regions with dramatically different densities.

Here we examine the fluid counterparts of systems considered in Refs.<sup>4–7</sup>; extremely persistent homogeneous three-dimensional active fluids in which the interparticle interactions never manage to balance the self-propulsion forces, which influence particles' motion in a nontrivial way. Although our fluid systems are less dense than those considered in Refs.<sup>4–7</sup>, they are dense enough such that we do not observe MIPS.

The parameter space of active systems is much larger than that of passive ones. At a minimum, one has to specify the average strength of active forces and their persistence time in addition to the set of parameters characterizing the corresponding passive system. If one considers athermal active systems, this results in a three-dimensional control parameter space. Thus, when comparing results of diverse studies, one needs to specify the path in the parameter space that one is following.

Early studies of dense homogeneous active systems focused on the glassy dynamics and the active glass transition<sup>8–17</sup>. Many of these studies considered a limited range of persistence times

and often examined the behavior at constant active temperature  $T_a$  that characterizes the long-time motion of an isolated active particle. For many models of self-propulsion, with increasing  $\tau_p$  at constant  $T_a$ , the strength of active forces decreases and dense active systems typically become glassy, see Fig. 2c of Ref.<sup>6</sup> for a recent example. Thus, to investigate the effects of extremely persistent active forces it is common to fix the force strength while increasing the persistence time<sup>4,5,7</sup>.

Recently, it has been shown that interesting behavior emerges for large persistence times in dense active particle systems<sup>4–7,18</sup>. First, it was found that some extremely persistent dense active systems relax only on the time scale of  $\tau_p$ <sup>4–7</sup>. In these systems, the mean squared displacement (MSD) and the two-point overlap function exhibit well-defined large  $\tau_p$  limits when plotted versus time rescaled by the persistence time<sup>5</sup>. Distributions of velocity components were found to exhibit fat exponential tails<sup>6</sup>. An intermediate-time plateau in the MSD was absent. Instead, a region of super-diffusive MSD scaling with time as  $t^\beta$  with  $\beta \approx 1.6$  for times less than  $\tau_p$  was identified<sup>6</sup>.

Second, a recent study of two-dimensional active systems that remain fluid in the large persistence time limit showed that mesoscale advective flows, forming streams and vortices, emerge for large persistence times<sup>18</sup>. These flow patterns resemble turbulent flows, and thus this phenomenon was identified as a type of “active turbulence”.

Here we study dense homogeneous three-dimensional active systems that do not undergo dynamic arrest in the large persistence time limit. In contrast to findings of Refs.<sup>4–7</sup>, for large persistence times  $\tau_p$  many dynamic properties of our systems become  $\tau_p$ -independent and their relaxation functions exhibit well-defined limits when plotted versus time *not* scaled by the persistence time. Also, the large persistence time limits of many prop-

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erties depend on the strength of active forces  $f$  as power laws. We emphasize that both the active systems studied in Refs.<sup>4–7</sup> and the fluid-like active systems studied here exhibit universal behavior, but on very different time scales.

It could be surmised that, in the infinite  $\tau_p$  limit, our systems lie in the un-jammed phase of the active yielding phase diagram studied (in two dimensions) by Liao and Xu<sup>19</sup>. However, these authors studied systems of particles interacting via a harmonic repulsion, in which interparticle forces are bounded, and thus strong enough self-propulsion forces are always able to induce motion. In contrast, in our systems interparticle forces are unbounded, and thus their jamming phase diagram may be different from that obtained by Liao and Xu.

## 2 Simulations

We study a three-dimensional 50:50 binary mixture of spherically symmetric active particles interacting via the Weeks-Chandler-Andersen<sup>20</sup> potential,

$$V_{\alpha\beta} = 4\epsilon \left[ \left( \frac{\sigma_{\alpha\beta}}{r} \right)^{12} - \left( \frac{\sigma_{\alpha\beta}}{r} \right)^6 \right] \quad (1)$$

for  $r < \zeta_{\alpha\beta} = 2^{1/6} \sigma_{\alpha\beta}$  and 0 otherwise. Here,  $\alpha, \beta$  denote the particles species A or B and  $\epsilon$  is the unit of energy. The distance unit is set by  $\sigma_{BB} = 1.0$ ,  $\sigma_{AA} = 1.4$ , and  $\sigma_{AB} = 1.2$ . We study the number density  $N/V = 0.451$ , which corresponds to the volume fraction  $\phi = \pi N[\zeta_{AA}^3 + \zeta_{BB}^3]/(12V) = 0.625$ .

We use the athermal active Ornstein-Uhlenbeck particle model without thermal noise<sup>21–23</sup>. The equation of motion for the position  $\mathbf{r}_n$  of particle  $n$  is

$$\xi_0 \dot{\mathbf{r}}_n = \mathbf{F}_n + \mathbf{f}_n, \quad (2)$$

where  $\mathbf{F}_n = -\sum_{m \neq n} \nabla_n V(r_{nm})$  and  $\mathbf{f}_n$  is the active force.  $\xi_0 = 1$  is the friction coefficient of an isolated particle and  $\xi_0 \sigma_{BB}^2/\epsilon$  sets the unit of time.

In turn, the self-propulsion forces evolve according to the Ornstein-Uhlenbeck process. The equation of motion for the self-propulsion force acting on particle  $n$  reads

$$\tau_p \dot{\mathbf{f}}_n = -\mathbf{f}_n + \zeta_n, \quad (3)$$

where  $\tau_p$  is the persistence time of the self-propulsion and  $\zeta_n$  is a Gaussian white noise with zero mean and variance  $\langle \zeta_n(t) \zeta_m(t') \rangle_{\text{noise}} = 2\xi_0 T_a \mathbf{I} \delta_{nm} \delta(t - t')$ , where  $\langle \dots \rangle_{\text{noise}}$  denotes averaging over the noise distribution,  $T_a$  is a single particle effective temperature,  $\mathbf{I}$  is the unit tensor and we set the Boltzmann constant  $k_B = 1$ . The root-mean square strength of active forces is  $f = \sqrt{3T_a/\tau_p}$ . The time step for the simulations ranged from 0.01 to 0.0001 with a smaller time step needed for larger  $f$  and longer  $\tau_p$ . We simulated at least 10  $\tau_p$  for each of 4 productions runs, which required up to  $10^9$  time steps.

To check for finite size effects we studied systems of 1K, 8K, and 64K particles at large persistence times. We did not find any system size dependence of the system's dynamics or properties. The results shown in this work are for 1K particles unless otherwise noted in the figure caption.

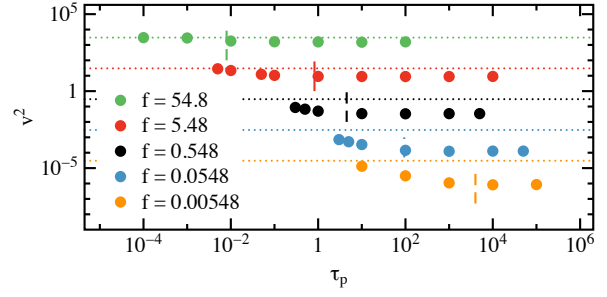


Fig. 1 The persistence time dependence of the mean squared velocity  $v^2 = \langle \dot{\mathbf{r}}^2 \rangle$  for fixed active force strength  $f$ . The velocity decreases and then saturates. The vertical dashed lines indicate approximate  $\tau_p$  at which the saturation occurs. The dotted lines are the free particle value of  $v^2 = f^2$  for each corresponding color.

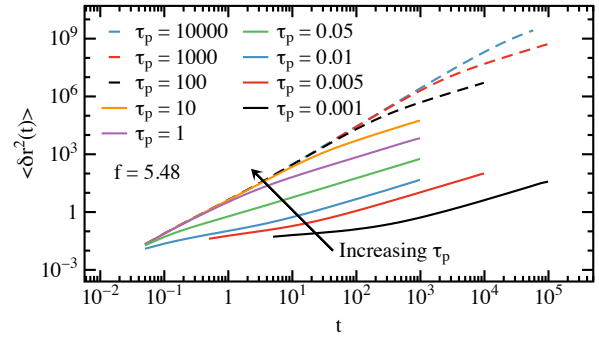


Fig. 2 The mean squared displacement (MSD)  $\langle \delta r^2(t) \rangle$  for  $f = 5.48$  and several  $\tau_p$ . For shorter  $\tau_p$  MSDs display a glassy plateau followed by diffusive motion. With increasing  $\tau_p$  an extended superdiffusive region appears that is analyzed in Fig. 4a.

## 3 Single-particle dynamics

We start by examining the persistence time dependence of the mean square velocity,  $v^2 \equiv N^{-1} \langle \sum_n \dot{\mathbf{r}}_n^2 \rangle = N^{-1} \langle \sum_n \mathbf{F}_n^2 + \sum_n 2\mathbf{F}_n \cdot \mathbf{f}_n \rangle + f^2$ , which is shown in Fig. 1. With increasing persistence time  $v^2$  decreases and then saturates. For the range of  $f$  that we studied, the cancellation of the interparticle and active forces is never complete and the system does not become arrested on the time scale of the persistence time<sup>5,7</sup>.

Next, we examine the persistence time dependence of the MSD

$$\langle \delta r^2(t) \rangle = N^{-1} \left\langle \sum_n [\mathbf{r}_n(t) - \mathbf{r}_n(0)]^2 \right\rangle, \quad (4)$$

shown in Fig. 2. At short times the motion is ballistic and it is determined by  $v^2$ ,  $\langle \delta r^2(t) \rangle \approx v^2 t^2$ <sup>11</sup>. For large persistence times the short-time dynamics become  $\tau_p$ -independent, which confirms saturation of  $v^2$ . For small  $\tau_p$  and fixed  $f$  we observe a well developed intermediate time plateau crossing over to diffusive behavior at long times. This behavior, characteristic of a system close to dynamic arrest, can be understood by noticing that the state point  $f = 5.48$ ,  $\tau_p = 0.001$  is the same as the state point  $T_{\text{eff}} = 0.01$ ,  $\tau_p = 0.001$  of Ref.<sup>17</sup>. Figs. 2c and 3 of the latter paper show that this state point is on the fluid side of the glass phase diagram, very

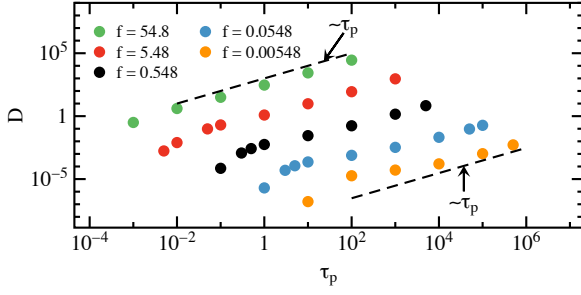


Fig. 3 The long time diffusion coefficient  $D$  for fixed active force strength  $f$  as a function of persistence time.  $D$  grows with increasing  $\tau_p$  and is proportional to  $\tau_p$  in the large  $\tau_p$  limit.

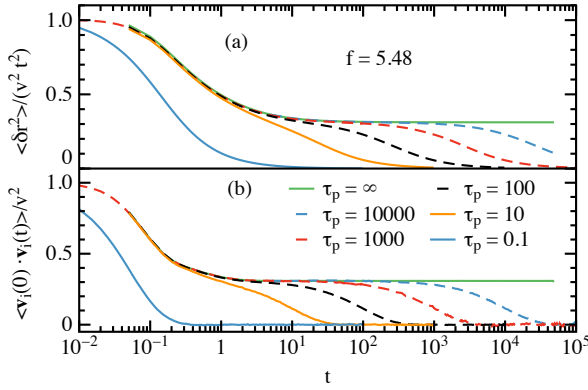


Fig. 4 (a) The mean squared displacement divided by  $v^2 t^2$  for  $f = 5.48$ . For large  $\tau_p$ , at short times the motion is ballistic (not shown), and then there is a super-diffusive regime. For very large  $\tau_p$  the super-diffusive motion speeds up and the resulting behavior is approximately ballistic. For  $t > \tau_p$  long-time diffusive motion is observed. In the  $\tau_p \rightarrow \infty$  limit the system stays in the second ballistic regime, with velocity  $v_1$ . (b) The velocity autocorrelation function. For very large  $\tau_p$  an intermediate time plateau is observed. The level of the plateau is the same in both panels.

close to the dynamic arrest line. With increasing  $\tau_p$  the intermediate time plateau characteristic of glassy dynamics disappears and the system becomes more fluid like. Similar fluidization with increasing  $\tau_p$  and fixed  $f$  was observed in two-dimensional systems by Paoluzzi *et al.*<sup>24</sup>.

At long times, the MSD exhibits diffusive behavior. The self-diffusion coefficient,  $D = \lim_{t \rightarrow \infty} \langle \delta r^2(t) \rangle / (6t)$ , is shown in Fig. 3. For a given  $f$  it monotonically increases with increasing  $\tau_p$ . At large  $\tau_p$  we find that  $D \sim \tau_p$ , indicated by the dashed lines.

We find a surprising time-dependence of the MSD between the initial ballistic and the long-time diffusive regimes. In Fig. 4a we show the MSD divided by  $v^2 t^2$  to show this time-dependence more clearly. The MSD exhibits a superdiffusive behavior that does not follow a single power law. Instead, a second, intermediate time ballistic regime appears, with velocity  $v_1$ . In the  $\tau_p \rightarrow \infty$  limit the system stays in the second ballistic regime.

The time-dependence of the MSD shown in Fig. 4a is reflected in the time-dependence of the velocity auto-correlation function,  $VACF(t) = N^{-1} \sum_n \langle \dot{\mathbf{r}}_n(t) \cdot \dot{\mathbf{r}}_n(0) \rangle$ , shown in Fig. 4b. We observe a two-step decay, which becomes increasingly more pronounced with increasing  $\tau_p$ . This agrees with the observation in Ref.<sup>18</sup>,

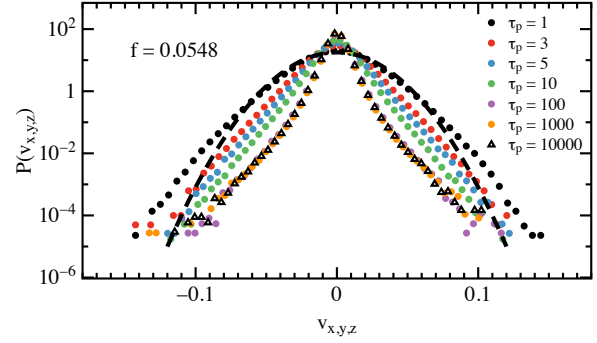


Fig. 5 The distribution of velocities for  $f = 0.0548$  for a range of persistence times. The distributions are non-Gaussian; the tails become more prominent with increasing  $\tau_p$  until the persistence time at which  $v^2$  saturates and then the distributions overlap. The dashed line represents a Maxwell distribution.

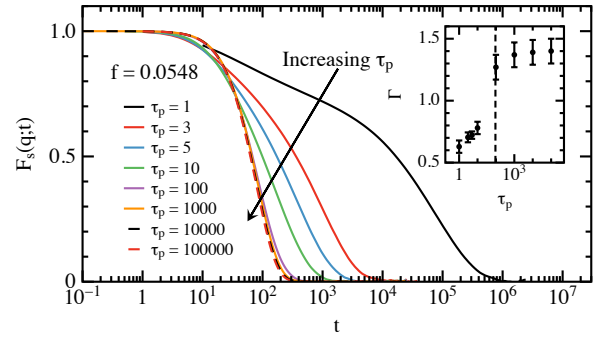


Fig. 6 The self-intermediate scattering function  $F_s(k; t)$  for  $f = 0.0548$  and a range of persistence times. The relaxation time decreases with increasing persistence time until around  $\tau_p(f)$  where  $F_s(k; t)$  becomes independent of persistence time. The inset shows the  $\tau_p$  dependence of parameter  $\Gamma$  of fit to  $F_s(k; t) = ae^{-(t/\tau_s)^\Gamma}$ .

except that, since we keep strength of active forces  $f$  constant, we find the plateau level to be  $\tau_p$  independent. We note that the plateau levels are the same in both panels of Fig. 4. The plateau level decreases with increasing self-propulsion force  $f$ .

Results shown in Fig. 4 suggest that, in the large  $\tau_p$  limit, the diffusion can be thought of as a random walk consisting of steps of length  $v_1 \tau_p$  taken every  $\tau_p$ . This reasoning rationalizes the observed scaling  $D \sim \tau_p$ . We note that while the physical picture of a random walk of steps of length  $v_1 \tau_p$  taken every  $\tau_p$  is identical to the behavior expected at low densities, our systems exhibit highly non-trivial dependence of  $v_1$  on the average self-propulsion force  $f$ , which we will discuss at the end of this section. This dependence suggests that the observed effective behavior is strongly influenced by interparticle interactions.

In Fig. 5 we show velocity distributions. As found by Keta *et al.*<sup>6</sup>, the distributions are strongly non-Gaussian. Their broad tails become more prominent with increasing  $\tau_p$  and then saturate.

The evolution of the MSD with the persistence time is reflected

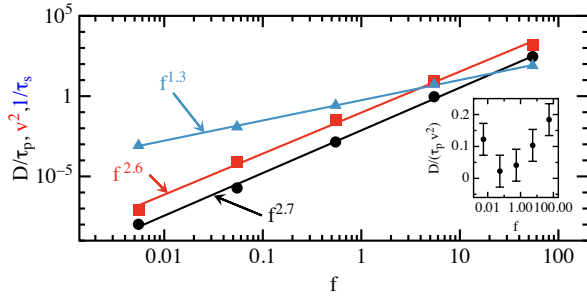


Fig. 7 The large  $\tau_p$  limit of  $v^2$ ,  $D/\tau_p$  and  $1/\tau_s$  as a function of the self-propulsion strength  $f$ . The lines are power law fits where the exponent for  $v^2$ ,  $D/\tau_p$  and  $1/\tau_s$  are  $2.6 \pm 0.1$ ,  $2.7 \pm 0.1$  and  $1.3 \pm 0.1$ , respectively. The inset shows  $D/(\tau_p v^2)$  versus  $f$ , which varies between 0.02 to 0.2 over the range of  $f$  examined, with error bars of about 0.05.

in the  $\tau_p$  dependence of the self-intermediate scattering function

$$F_s(k; t) = \frac{1}{N} \left\langle \sum_n e^{i\mathbf{k} \cdot (\mathbf{r}_n(t) - \mathbf{r}_n(0))} \right\rangle. \quad (5)$$

We chose  $k = 5.3$ , which is approximately equal the first peak of the total static structure factor. In Fig. 6 we show  $F_s(k; t)$  for  $f = 0.0548$ . With increasing  $\tau_p$  the intermediate time glassy plateau disappears and the decay changes from stretched exponential, to exponential, then to compressed exponential. Shown in the inset to Fig. 6 is the parameter  $\Gamma$  obtained from fits to  $F_s(k; t) = ae^{-(t/\tau_s)^\Gamma}$  where we restrict  $a \leq 1$ .  $\Gamma$  increases with increasing  $\tau_p$  and reaches a plateau above  $\tau_p \approx 94$ .

We find that the large persistence time limits of several properties discussed above depend on the strength of the active forces as power laws. In Fig. 7 we show the large  $\tau_p$  limits of  $v^2$  (squares),  $D/\tau_p$  (circles) and  $\tau_s$  (triangles). We find that the former two quantities follow a power law with  $f$  with statistically the same exponent,  $2.6 \pm 0.1$  for  $v^2$  and  $2.5 \pm 0.1$  for  $D/\tau_p$ . We note that  $v^2$  is bound from above by  $f^2$ , which corresponds to the limit in which interactions become irrelevant compared to active forces. Thus, we do not expect the power law for  $v^2$  to extend up to arbitrary large  $f$ . The power law of the relaxation time,  $\tau_s \sim f^{-1.3}$  can be related to that of  $v^2$ ; in the large  $\tau_p$  limit  $F_s$  decays on the time scale on which a particle moves over its diameter, which scales as  $v$ .

We also analyzed the dependence of  $v_1$  on the strength of the active forces. To this end, we extracted  $v_1^2$  for each  $f$  by finding the plateau in  $\langle \delta r^2(t) \rangle / t^2$  for large  $\tau_p$ . We find that  $v_1^2$  grows with  $f$  as a power law with an exponent of  $2.5 \pm 0.2$ , which is statistically the same as the exponent for  $v^2$ .

We emphasize that, although in the large  $\tau_p$  limit the motion of individual particles is close to ballistic, it is still influenced by the interparticle interactions in a highly non-trivial way. If there were no interparticle interactions, scaling exponents discussed above would have values 2 and  $-1$ .

## 4 Collective dynamics

The quantities discussed in Sec. 3 describe the single particle motion in our many-particle systems. To examine collective properties we investigated the persistence time dependence of the stress

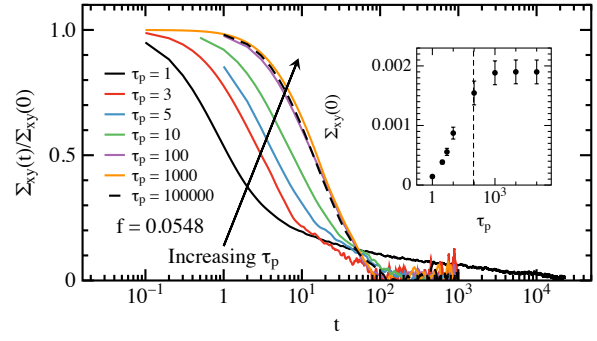


Fig. 8 Normalized shear-stress correlation function for  $f = 0.0548$  and a range of persistence times. The shear stress relaxation time initially increases with  $\tau_p$ , and then saturates. Additionally,  $\Sigma_{xy}(0)$  grows with  $\tau_p$  and then it becomes approximately constant (inset).

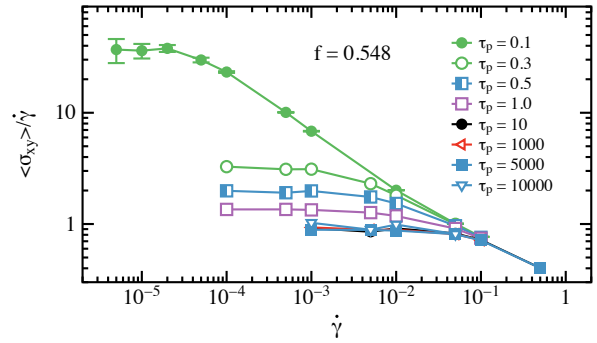


Fig. 9 Average shear stress  $\langle \sigma_{xy} \rangle$  divided by the shear rate  $\dot{\gamma}$  at constant  $f = 0.548$  for a range of persistence times. The zero-shear-rate viscosity  $\eta$  can be obtained from the small shear rate plateau.

fluctuations and the rheological response. First, we examined the shear-stress correlation function  $\Sigma_{\alpha\beta}(t) = V^{-1} \langle \sigma_{\alpha\beta}(t) \sigma_{\alpha\beta}(0) \rangle$ , where

$$\sigma_{\alpha\beta}(t) = -\frac{1}{2} \sum_n \sum_{m \neq n} \frac{r_{nm}^\alpha r_{nm}^\beta}{r_{nm}} \frac{dV(r_{nm})}{dr_{nm}}, \quad (6)$$

and  $r_{nm}^\alpha$  is the  $\alpha$  component of the distance vector between particle  $n$  and particle  $m$ .

In Fig. 8 we report the normalized shear stress correlation function,  $\Sigma_{xy}(t)/\Sigma_{xy}(0)$ , for  $f = 0.0548$  and a large range of  $\tau_p$ . For small  $\tau_p$  there is a rapid decay to an emerging plateau followed by a slow decay to zero. With increasing persistence time, the decay of  $\Sigma_{xy}(t)$  becomes less stretched; it is exponential for a long enough  $\tau_p$ . In the inset we show the dependence of the initial value,  $\Sigma_{xy}(0)$ , on  $\tau_p$ . We see that the initial value first grows with  $\tau_p$  and then plateaus.

To probe the rheological response of our active systems we simulated shear flow by adding to the equations of motion (2) a bulk non-conservative force  $\mathbf{F}_n^\gamma = \xi_0 \dot{\gamma}_n \mathbf{e}_x$  with Lees-Edwards boundary conditions<sup>25</sup>. In Fig. 9 we show the average shear stress,  $\langle \sigma_{xy} \rangle / \dot{\gamma}$ , for  $f = 0.548$  and a large range of  $\tau_p$ . In Fig. 10 we show the  $\tau_p$  dependence of the zero-shear-rate viscosity obtained from the small  $\dot{\gamma}$  plateaus of  $\langle \sigma_{xy} \rangle / \dot{\gamma}$ . We find that  $\eta$  initially decreases and reaches a  $\tau_p$ -independent plateau.

We find that large  $\tau_p$  limits of collective properties also scale

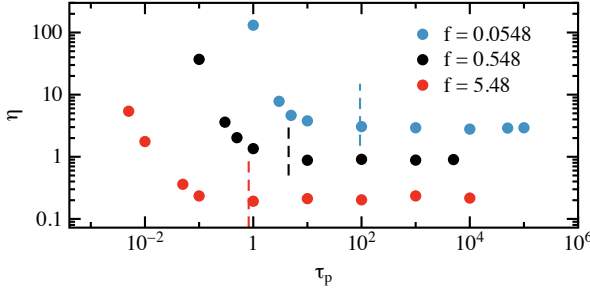


Fig. 10 The viscosity  $\eta$  at fixed average active forces strength versus  $\tau_p$ . The viscosity initially decreases with increasing  $\tau_p$  and then becomes constant. The vertical dashed lines indicate approximate  $\tau_p$  at which the saturation occurs.

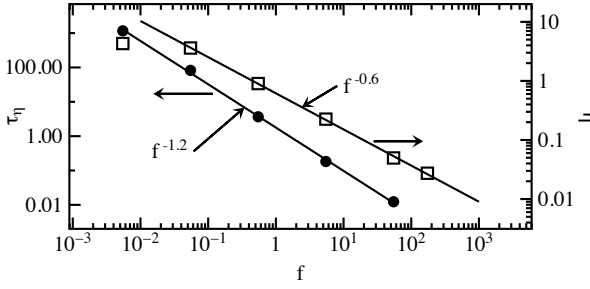


Fig. 11 The large  $\tau_p$  limit of the relaxation time of  $\Sigma_{xy}(t)/\Sigma_{xy}(0)$  (defined as when this function equals  $e^{-1}$ ) and of viscosity  $\eta$  as a function of the self-propulsion strength  $f$ . The lines are power law fits where the exponents for  $\tau_\eta$  and  $\eta$  are  $-1.2 \pm 0.1$  and  $-0.60 \pm 0.1$ , respectively.

with  $f$  as power laws. In Fig. 11 we show the dependence of the large  $\tau_p$  limits of the relaxation time of the normalized stress tensor autocorrelation function and of the viscosity on the strength of the self-propulsion.

## 5 Static structure

When analyzing the dynamics in passive systems, one usually tries to make connections between the average distribution of the particles and their dynamics. To check how the average arrangement of the particles in our active systems changes with increasing  $\tau_p$  we evaluated the total steady-state structure factor

$$S(k) = \frac{1}{N} \left\langle \sum_{n,m} e^{i\mathbf{k} \cdot (\mathbf{r}_n - \mathbf{r}_m)} \right\rangle, \quad (7)$$

with summation over all particles in the system, and the pair correlation function between the large particles,

$$g(r) = \frac{V}{(N/2)^2} \left\langle \sum_n \sum_{m \neq n} \delta(\mathbf{r} - (\mathbf{r}_n - \mathbf{r}_m)) \right\rangle, \quad (8)$$

where the summation extends over the large particles only.

In Fig. 12 we show that the peak height of  $S(k)$  initially decreases with increasing  $\tau_p$ , which nicely correlates with relaxation getting faster and viscosity decreasing. The peak height then saturates at a large persistence time. The large  $\tau_p$  limit of the structure factors still looks liquid-like, and thus it does not signal the

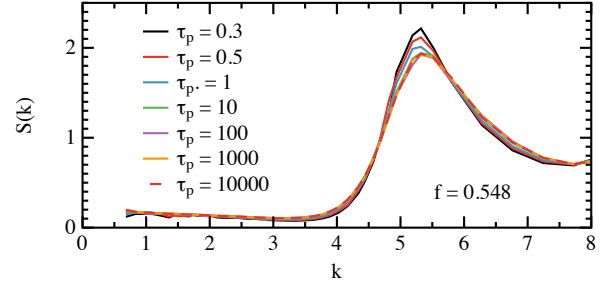


Fig. 12 The persistence time dependence of the steady-state structure factor  $S(k)$  for  $f = 0.548$ . The peak height decreases with increasing persistence time and then saturates at a large persistence time.

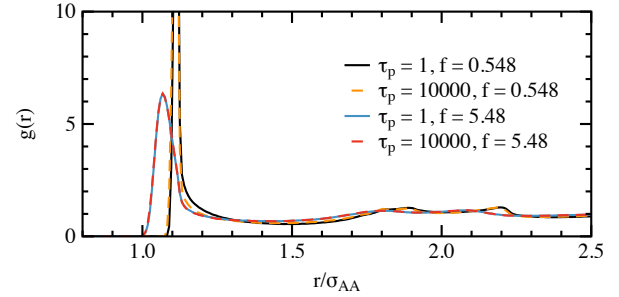


Fig. 13 The pair correlation function between the large particles for  $f = 0.548$  and  $f = 5.48$  and two representative persistence times  $\tau_p = 1$  and  $\tau_p = 10000$ . For small  $f$  the first peak in  $g(r)$  is sharply peaked; it is shown in Fig. 14. There is little dependence of  $g(r)$  on the persistence time.

interesting dynamics we uncovered. To describe the dynamics of extremely persistent dense active fluids one cannot rely upon static structure factors only.

We note that the absence of any small wavevector peak implies that our systems are homogeneous. We confirmed this observation by evaluating local density histograms at several simulated state points. Incidentally, the homogeneity allows us to rationalize the absence of shear thickening that was observed at large self-propulsions in Ref.<sup>26</sup>, which was attributed to clustering and was dubbed motility-induced shear thickening.

In Figs. (13-14) we show the pair correlation function between the large particles. In contrast to the behavior exhibited by the structure factor,  $g(r)$  is not very sensitive to the changes in the persistence time. However, the height of the first peak of  $g(r)$  decreases rapidly with increasing self-propulsion force  $f$  while the position of the first peak shifts to the smaller  $r$ , suggesting that the system becomes effectively less crowded with increasing  $f$ . This rationalizes decreasing large persistence time limits of the relaxation times and of the viscosity and increasing large persistence time limit of the diffusion coefficient

## 6 Conclusions

We conjecture that the systems we presented here form a new class of extremely persistent active matter systems. Earlier investigations<sup>5,7</sup> revealed systems that relax on the time scale of the self-propulsion and exhibit intermittent dynamics. The sys-



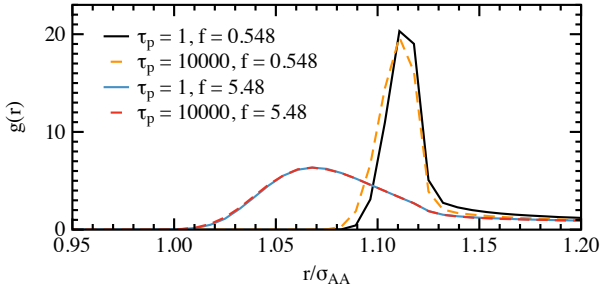


Fig. 14 The first peak of the pair correlation function between the large particles for  $f = 0.548$  and  $f = 5.48$  and two representative persistence times  $\tau_p = 1$  and  $\tau_p = 10000$ . The peak height decreases and the peak width increases with increasing  $f$ . There is little dependence of the first peak of  $g(r)$  on the persistence time.

tems with intermittent dynamics slow down with increasing  $\tau_p$  at fixed  $f$ , but they exhibit universal dynamics when investigated as a function of time rescaled by  $\tau_p$ . In contrast, the systems we presented exhibit universal dynamics for large  $\tau_p$  but as a function of un-scaled time. We expect that for higher volume fractions there is a transition between the regime in which the relaxation becomes independent of the persistence time of the self-propulsion, which is the regime we analyzed, and the regime in which the system flows only on the time scale of the self-propulsion, which is the regime investigated earlier<sup>5,7</sup>. At a fixed strength of the active forces the transition is driven by the density. Since inter-particle interactions in our systems are not bounded, it is not obvious whether, at constant density, the transition can be induced by changing the force strength, like in two-dimensional systems with harmonic interactions, which were investigated by Liao and Xu<sup>19</sup>.

For the systems studied here, the single-particle motion exhibits two ballistic regions separated by a superdiffusive regime. Classic signatures of two-step relaxation seen in glassy systems are absent both in the MSD and  $F_s(k; t)$ . Therefore, our systems are not close to a glass transition. Many properties that quantify the large persistence time limit of the relaxation depend on the strength of the active forces as a power law.

In contrast to Keta *et al.*<sup>18</sup>, we did not observe essential features of “active turbulence”. In particular, mean-squared displacement difference of initially close by particles monitored by Keta *et al.* was only slightly smaller than the MSD. This may not be surprising since our systems have somewhat lower density and larger active forces.

Finally, we note that approximate theories developed to describe the relaxation in active fluids<sup>29–34</sup> we tested against computer simulations for rather limited range of self-propulsion force  $f$  and persistence time  $\tau_p$ . The discovery of a new paradigm of extremely persistent active fluids with non-trivial power laws calls for additional theoretical work. The single-particle motion in the systems that we investigated can perhaps be described as a persistent random walk with renormalized density and self-propulsion-dependent velocity. However, it is unlikely that such a simple picture could also account for the collective dynamic properties.

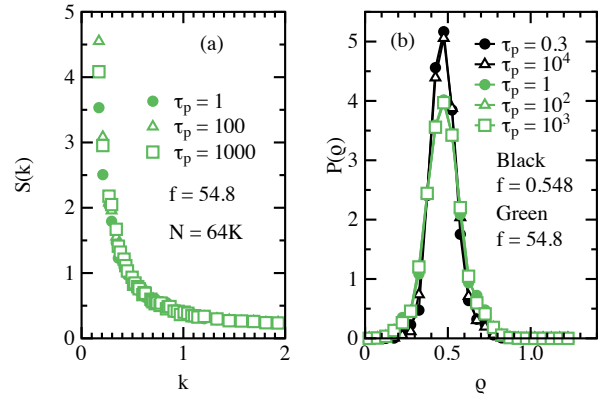


Fig. 15 (a) The structure factor  $S(k)$  for  $f = 54.8$  for our 64K system. There is a distinct upturn at small wavevectors which is not seen for our other  $f$ . (b) The density distribution for the same  $f$ ,  $N$ , and  $\tau_p$  as shown in (a), and for two state points for which structure factors were shown in Fig. 12. We do not observe the two peak structure expected for systems undergoing motility induced phase separation.

## 7 Appendix

Motility induced phase separation (MIPS) is often reported and extensively studied in active systems. In two-dimensional systems it has been shown that MIPS will occur with increasing  $\tau_p$  and fixed  $f$  over a range of densities<sup>24</sup>. We note that, in three-dimensional systems states with MPIS, coexistence is metastable with respect to active crystallization in a large region of the parameter space<sup>35</sup>. Here we claim that our system is not undergoing MIPS at any of the state points studied.

To determine if the system undergoes MIPS one can examine structural signatures and density distributions. One structural signature is an upturn of the structure factor  $S(k)$  at small wavevectors  $k$ . However, an upturn can occur if there are large density fluctuations but no MIPS. The density distribution is another method to examine MIPS, where two peaks indicate that the system has separated into high density and low density regions. We use these two methods.

To calculate the density distribution we divide the system into boxes of length  $\ell = 2.6082$  and determine the density inside each box. For each configuration used in the calculation we determine the density distribution of the small boxes and then average these distributions. To examine smaller wavevectors and to get a more accurate density distribution we present results for our 64K systems.

As shown in Fig. 12 we don’t see any evidence of an upturn in the low wavevector values of  $S(k)$  for  $f = 0.548$ . The only  $f$  we do observe an upturn at small wavevectors is  $f = 54.8$ , Fig. 15(a). We note that the small  $k$  upturn does not increase with increasing  $\tau_p$ , and thus we don’t expect any changes with increasing  $\tau_p$ . The density distribution for these state points do not show a two peak structure and remains statistically unchanged when increasing  $\tau_p$ , Fig. 15(b). We conclude that there are large density fluctuations, but the system is not yet undergoing MIPS. For larger  $f$  we expect that this system will exhibit MIPS.

## Conflicts of interest

There are no conflicts to declare.

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## Data Availability

All the data presented in this manuscript can be found on Dryad<sup>36</sup>.

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