

Scaling Laws for Passive Polymer Dynamics in Active Turbulence

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Biological systems commonly combine intrinsically out-of-equilibrium active components with passive polymeric inclusions to produce unique material properties. To explore these composite systems, idealized models—such as polymers in active fluids—are essential to develop a predictive theoretical framework. We simulate a single, freely jointed passive chain in two-dimensional active turbulence. Active flows advect the polymer, producing a substantially enhanced diffusivity. Our results reveal that the dimensionless diffusivity obeys scaling laws governed by the Péclet, Weissenberg, and Ericksen numbers, which paves the way for designing active/polymeric hybrid materials with predictable properties that differ significantly from those of nondeformable inclusions.

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The presence of both active and passive components is a hallmark of biomaterials [1]. Examples include intracellular molecular motors interacting with passive cytoskeletal filaments [2–4], polymerases transcribing DNA [5], and the dynamic interplay between euchromatin and heterochromatin [6–11]. In such composite biomaterials, active components are intrinsically out of equilibrium due to continuous energy conversion into coherent movement [12,13], while passive components are driven from equilibrium through their sustained contact with the active phase. Previous work has focused primarily on rigid passive bodies embedded in active flows, including tracer particles in dilute suspensions of swimming bacteria [14–16], colloids in dense bacterial communities [17–21], and disks in active nematic films [22,23].

While previous studies have centered on rigid inclusions in active surroundings [24–30], polymers and flexible filaments embedded in active backgrounds are of specific biological importance. They are crucial for the formation of cytoplasmic flows [31–35], the organization of the extracellular matrix [36], and the transport of DNA [37]. Moreover, macromolecules are commonly found in bacterial environments, including gastrointestinal tracts, mucus, blood, and saliva [36,38]. Such systems are especially intriguing due to the interplay between activity and conformational entropy [39]. Prior work has mainly explored the conformational changes of semiflexible polymers in ensembles of self-propelled particles [40–42] or athermal baths [43–49]. These systems illustrate how intrinsically

out-of-equilibrium surroundings can actively deform passive polymeric inclusions with internal degrees of freedom, yet they do not account for the hydrodynamics of active fluids or the ensuing mesoscale turbulence in continuous active fluids. These spontaneous flows play a vital role in intracellular mixing and transport [50–53], as well as macromolecular nutrient distribution, chemical signal dispersal and molecular waste removal in active microbial communities [54–57]. While the enhanced transport in active flows is widely recognized [58–64], a predictive theoretical framework for the active transport of polymeric inclusions is lacking.

We study the dynamics of a single passive, freely jointed polymer embedded in active nematic (Fig. 1). Active nematics in bulk exhibit active turbulence, a low Reynolds number phenomenon with characteristic length and timescales (Fig. 1) [65,66]. This composite system exemplifies the competition between entropy-driven relaxation processes and activity-driven dynamics. We employ a hybrid simulation approach, combining active nematic multiparticle collision dynamics to model the active nematic background (AN-MPCD; Supplementary Material (SM) [67] and [82,83]) and molecular dynamics to simulate the polymer, which is coupled directly to the velocity but not the orientation field (SM [67]). Because the monomers participate in the MPCD collision step, they both respond to and perturb the local flow via momentum exchange, resulting in two-way coupling. This approach captures fluctuations through the thermal energy $k_B T$, hydrodynamic interactions, nematic elasticity through K , and active stresses via the activity ζ . We find that passive polymers within active nematic turbulence experience a thousandfold increase in diffusivity, arising from short-time advection of the polymers by active flows with a finite correlation time. This effect is characterized by a pair of dimensionless numbers: the Péclet number Pe , quantifying polymer

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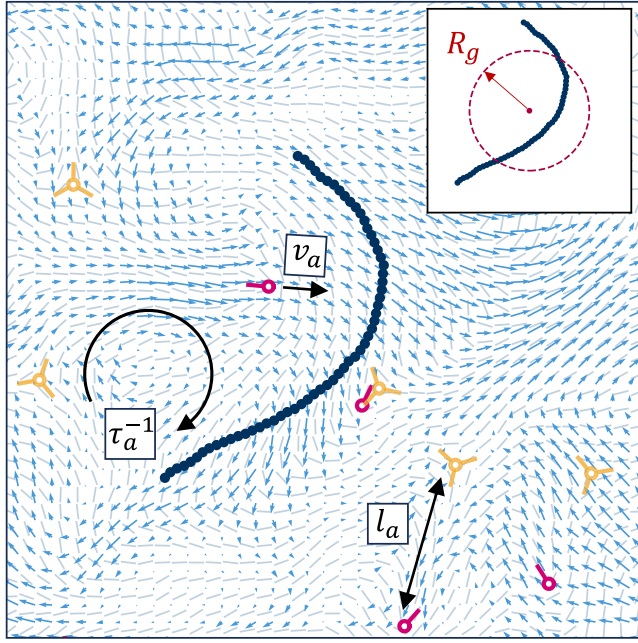


FIG. 1. Snapshot of a polymer in active nematic turbulence. The polymer consists of $N = 60$ navy monomers. The director is shown as silver lines and the flow field as blue vectors. Defects are depicted in ruby (+1/2) and amber (-1/2). The average distance between defects is the active length scale l_a . The active timescale τ_a is represented as the inverse of vorticity. The motile +1/2 defect speed is characterized as $v_a = l_a/\tau_a$. Activity is $\zeta = 0.08$ and system size is $l_{\text{sys}} = 80$. Inset: polymer center of mass and gyration radius R_g .

advection rate relative to the diffusion rate, and the Weissenberg number Wi , comparing the deformation rate with relaxation time. These are shown to be linked via the Eriksen number Er , representing the ratio of characteristic length scales. The interplay between these dimensionless numbers determines the enhanced diffusivity of passive polymers in active turbulence.

This Letter focuses on flow-aligning nematics with extensile activity ($\zeta > 0$) and values are reported in simulation units (SM [67]). The competition between nematic elasticity K , viscosity η , and activity ζ produce active length and timescales (Fig. S1; SM [67]) [84–86]

$$l_a \sim \sqrt{\frac{K}{\zeta}} \quad ; \quad \tau_a \sim \frac{\eta}{\zeta} \quad (1)$$

over which flow remains coherent in active turbulence. As activity ζ decreases, the active length scale l_a grows until it exceeds the system size l_{sys} [Fig. S1(a); SM [67]], at which point active chaotic flows vanish and residual flows are primarily dictated by confinement within periodic boundary conditions [82,87,88]. Our analysis focuses on the active turbulence regime, but the polymer's effective

diffusion monotonically approaches the thermal value D_0 as $\zeta \rightarrow 0$ when $l_a \gtrsim l_{\text{sys}}$ (Fig. S4; SM [67]).

The polymer dynamics are characterized by the mean squared displacement (MSD) of its center of mass $\langle \delta \underline{r}^2 \rangle = \langle (\underline{r}(t + \delta t) - \underline{r}(t))^2 \rangle$, where $\underline{r}$ is the center of mass position and $\langle \cdot \rangle$ is the ensemble and temporal average over all starting times t (SM [67]). Without activity, the polymer interacts with the thermal fluctuations and randomly explores space [Fig. 2(a)], leading to diffusive motion $\langle \delta \underline{r}^2 \rangle = 2dD_0\delta t$, where D_0 is the thermal diffusion coefficient in d dimensions [Fig. 2(c)]. This defines $\tau_r = R_g^2/D_0$, the polymer relaxation time over which it diffuses a distance equal to its gyration radius R_g (Fig. 1, inset).

In active turbulence, the polymer explores a larger region than its passive counterpart [Fig. 2(b)]. This is due to short-time advection by the active flows, where $\langle \delta \underline{r}^2 \rangle = v^2 \delta t^2$ with the center of mass propulsion speed of polymer v [Fig. 2(c)]. The propulsion speed increases with activity for short polymers but saturates at high activities [Fig. 2(d)]. For long polymers ($N \geq 50$), it decreases and becomes relatively independent of activity. However, the turbulent nature of these active flows decorrelates the advective motion. The time over which the polymer is persistently advected by active flows is τ_p [Fig. 2(c); SM [67]]. For all polymer lengths, τ_p decreases monotonically with activity, while increasing polymer length leads to longer propulsion times [Fig. 2(e)].

Together, the advective polymer speed v and the persistence time τ_p produce the long-time enhanced diffusion coefficient [Fig. 2(c); SM [67]]

$$D_e = D_0 + \frac{v^2 \tau_p}{d}. \quad (2)$$

This persistent dynamics model accurately describes the MSD of polymers in active turbulence [Fig. 2(c), dashed line]. All polymers show a significant increase in their effective diffusivity compared D_0 [Fig. 2(f)]. The observed diffusivity is $\mathcal{O}(10^3)$ times greater than the thermal diffusivity for $N = 200$. The enhanced diffusivity relative to the thermal diffusivity increases with length N . This enhancement arises from the active kicks the polymers receive. As activity increases, the fluid speed rises, advecting the polymers and causing the speed of short polymers to increase [Fig. 2(d)]; however, the coherent time of the flows decreases, corresponding to a reduction in the active timescale τ_a and correspondingly decreasing τ_p [Fig. 2(e)]. These counteracting effects of increased speed and reduced coherent time result in relatively constant effective diffusion coefficients, especially for $N < 100$ [Fig. 2(f)]. We confirm that polymer transport is not significantly influenced by finite-size effects by comparing simulations in larger domains (Fig. S7; SM [67]).

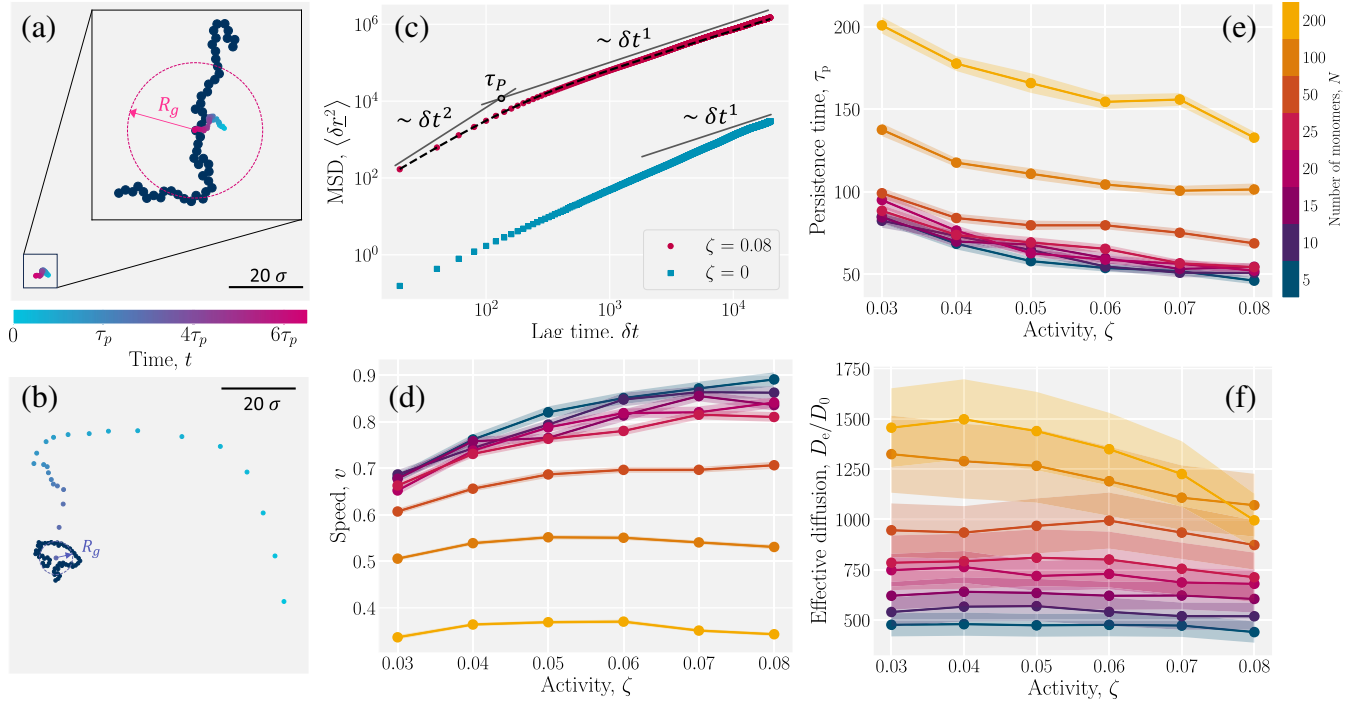


FIG. 2. Mean-squared displacement (MSD) of the polymer's center of mass. (a),(b) Representative center-of-mass trajectories of a polymer consisting of $N = 50$ monomers of size σ , color-coded by time t . (a) Trajectory without activity ($\zeta = 0$) and (b) with activity ($\zeta = 0.08$). Inset: close-up view of the polymer trajectory, conformation, and instantaneous gyration radius R_g . (c) The MSD $\langle \delta r^2 \rangle$ of the examples from (a) and (b) as a function of lag time δt , averaged over 40 realizations and longer times. The passive case ($\zeta = 0$) exhibits long-time diffusive behavior $\langle \delta r^2 \rangle \sim \delta t$. The active case ($\zeta = 0.08$) shows a short-time propulsive regime $\langle \delta r^2 \rangle \sim \delta t^2$, followed by effective diffusive behavior at times longer than the persistence time τ_p . (d) The center-of-mass speed v of polymers with N monomers. (e) Persistence time of advection by the flow. (f) Effective diffusion coefficients D_e normalized by thermal diffusivity D_0 . Shaded areas indicate the standard error of the mean, computed from fits to individual simulation repeats.

While polymers with $N < 100$ monomers maintain a relatively constant diffusivity, the $N \geq 100$ polymers exhibit decreasing diffusivity as activity rises. By Eq. (2), this can be seen to be due to the fact that the polymer speed becomes constant with activity [Fig. 2(d)], while the persistence time still decreases [Fig. 2(e)]. For the largest polymer ($N = 200$), a minor degree of nonmonotonicity can be observed in the effective diffusivity for the smallest activity [Fig. 2(f)], which is inherited from the speed squared [Fig. 2(d); Eq. (2)]. The observed nonmonotonicity is small compared to the uncertainty, but may be due to the largest polymers suppressing the onset of active turbulence.

We have seen that the speed v , persistence time τ_p and, consequently, effective diffusivity D_e vary with the activity of the surrounding fluid ζ and the polymer length N . For macromolecules embedded in a flowing fluid, the dimensionless Péclet number (Pe) quantifies the importance of advection relative to thermal diffusion [89,90] as

$$\text{Pe} = \frac{v/R_g}{D_0/R_g^2} = \frac{vR_g}{D_0}, \quad (3)$$

where R_g is the gyration radius of the polymer in equilibrium, which in $d = 2$ dimensions [91,92] is verified

numerically to scale as

$$R_g \sim N^{3/4} \quad (4)$$

(Fig. S2; see SM [67]). Computational studies in 2D have demonstrated the theoretical expectation for the diffusion coefficient is

$$D_0 = \frac{k_B T}{4\pi\eta} \left[\ln \frac{l_{\text{sys}}}{R_g} + \text{const} \right] \equiv \mathcal{D}_N^{-1}, \quad (5)$$

where l_{sys} is the system size (SM [67]) [93,94]. The change of variables to inverse diffusion $\mathcal{D}_N$ highlights the logarithmic dependence on N via R_g (SM [67]).

The polymer's advection speed v is expected to scale with the characteristic active flow speed v_a , which itself goes as

$$v_a \sim \frac{l_a}{\tau_a} \sim \zeta^{1/2} \quad (6)$$

by Eq. (1) [85]. Moreover, the polymer's speed v exhibits a length-dependent behavior [Fig. 2(d)]. Specifically, shorter polymers ($N \leq 50$) show increasing speed with activity, while longer polymers exhibit a saturation, with their speed

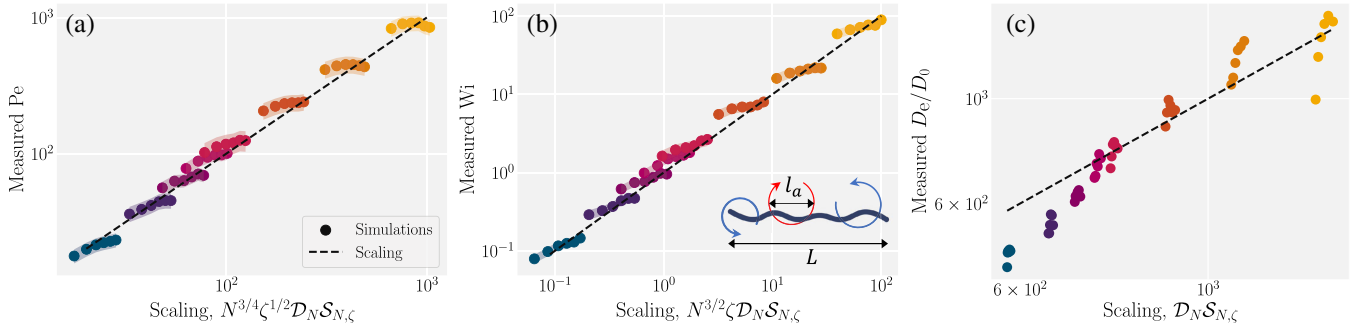


FIG. 3. Dynamics of polymers expressed in terms of dimensionless numbers and relevant scaling. (a) Péclet number Pe . (b) Weissenberg number Wi . (c) Effective diffusion D_e normalized by thermal diffusivity D_0 . The y axes are measured from simulations and the x axes are the scaling predictions for polymers of length N in active turbulence with activity ζ . $\mathcal{D}_N$ is the inverse of the thermal diffusion coefficient and $\mathcal{S}_{N,\zeta}$ is the crossover factor on the polymer velocity. Inset: the polymer contour length L can be subdivided into $n = L/l_a$ segments that are as long as the activity length scale l_a .

generally decreasing with length. To explain the lower speeds of longer polymers, consider a polymer that can span many decorrelated regions of active turbulence (Fig. 3 inset). Such a long polymer interacts with multiple “patches,” each of which is a coherent flow, but decorrelated from neighboring patches. The cumulative effect of these randomly misaligned flows reduces the average velocity of the polymer’s center of mass, explaining the decline in enhanced diffusivity for long polymers ($N = 200$) and high activities [Fig. 2(f)]. To model this effect, consider a sufficiently long polymer as a chain of $n = L/l_a \sim N\zeta^{1/2}$ segments, each subject to uncorrelated flow. For large n , the variance scales as $\langle v^2 \rangle \sim n$. The center of mass speed is then the average of these across the n segments, leading to $v \sim 1/\sqrt{n}$. Hence, the polymer speed requires a correction factor $v = \mathcal{S}_{N,\zeta} v_a$. The function $\mathcal{S}_{N,\zeta}$ should be unity for small lengths $n \ll 1$ and cross over to scale as $n^{-1/2}$ for $n \gg 1$. We pragmatically choose

$$\mathcal{S}_{N,\zeta}(\alpha, \beta; n) = \alpha \left(\frac{1}{1 + \frac{n^{1/2}}{\beta}} \right), \quad (7)$$

which has the correct limiting behaviors, where the fitting parameters α absorbs any nonscaling factors and β controls the crossover point. The subscripts emphasize that the crossover function depends on both N and ζ through n . In the short polymer limit, $\lim_{n \ll 1} \mathcal{S}_{N,\zeta} \rightarrow \text{constant}$, which means that $v = \lim_{n \ll 1} \mathcal{S}_{N,\zeta} v_a \sim \zeta^{1/2}$. On the other hand, in the long polymer limit, $\lim_{n \gg 1} \mathcal{S}_{N,\zeta} \rightarrow \alpha n^{-1/2}/2$ and $v = \lim_{n \gg 1} \mathcal{S}_{N,\zeta} v_a \sim \zeta^{1/4} N^{-1/2}$.

This correction factor is applied to the polymer speed and, along with the scaling expectations for R_g [Eq. (4)] and D_0 [Eq. (5)], is substituted into Eq. (3) to find

$$Pe \sim N^{3/4} \zeta^{1/2} \mathcal{D}_N \mathcal{S}_{N,\zeta}. \quad (8)$$

The predicted scaling agrees well with measured simulation values [Fig. 3(a)]. The motion of all polymers is primarily governed by advection, as indicated by $Pe \gg 1$ [Fig. 3(a)],

consistent with the large effective diffusivity (Fig. 2). The fitting parameter $\alpha = 1.47 \pm 0.03$ captures the combined effects of absorbed constants, while $\beta = 11.3 \pm 1.3$. While the empirical crossover factor $\mathcal{S}_{N,\zeta}$ is seen to improve the prediction (SM [67]), it does not fully collapse the data, especially at longest polymer lengths considered, where steps between contour lengths remain visible. Theoretically predicting this crossover function to better describe the long-polymer regime is expected to further narrow the steps and is a direction for future work.

Having collapsed the polymer’s advective velocity through the Péclet number [Fig. 3(a)], the same can be done for the persistent time via the Weissenberg number [95,96]

$$Wi = \frac{\tau_r}{\tau_p}, \quad (9)$$

where $\tau_r = R_g^2/D_0$ is the thermal relaxation time of the polymer and τ_p is the timescale of polymer advection. The polymer’s persistent advection time is estimated as

$$\tau_p = \frac{l_a}{v} = \frac{l_a}{v_a \mathcal{S}_{N,\zeta}}, \quad (10)$$

which is the time it takes for the polymer to explore a distance l_a at speed v . The ratio of τ_r and τ_p causes the Weissenberg number to scale as

$$Wi = \frac{R_g^2}{D_0} \frac{v}{l_a} \sim N^{3/2} \zeta \mathcal{D}_N \mathcal{S}_{N,\zeta}. \quad (11)$$

The expected value of Wi matches well with the simulations [Fig. 3(b)]. The shorter polymers ($N < 20$) follow the expectation and relax to a relatively undeformed conformation, as indicated by $Wi < 1$. Longer polymers with $Wi \gg 1$ are substantially elongated, extending over many active length scales.

Since the Péclet number describes advection speed and the Weissenberg number characterizes propulsion time, the effective diffusion coefficient [Eq. (2)] can be expressed in terms of these dimensionless numbers as

$$\frac{D_e}{D_0} - 1 = \frac{\text{Pe}^2}{d\text{Wi}} \sim \mathcal{D}_N \mathcal{S}_{N,\zeta}. \quad (12)$$

Surprisingly, the power law scaling of Pe^2 with N and ζ cancels with the scaling of Wi in $d = 2$ dimensions, leaving only the logarithmic scaling $\mathcal{D}_N$ [Eq. (5)] and the crossover scaling $\mathcal{S}_{N,\zeta}$ [Eq. (7)]. Equation (12) is found to predict the effective diffusivity [Fig. 3(c)].

Most studies of active nematics have focused on quasi-two-dimensional films, largely due to the experimental challenges of accessing and resolving nematic fields in three dimensions [97–99]. However, recent works—including numerical simulations and experimental methods—are exploring the structure of active turbulence in 3D [100–103]. The scaling relations are expected to become $\text{Pe} \sim N^{2\nu} \zeta^{1/2} \mathcal{S}_{N,\zeta}$ and $\text{Wi} \sim N^{3\nu} \zeta \mathcal{S}_{N,\zeta}$, leading to $D_e/D_0 - 1 \sim N^\nu \mathcal{S}_{N,\zeta}$ with $\nu \approx 3/5$ in a three-dimensional good solvent.

So far, the dynamics has been described in terms of timescales, but this can be interpreted through a length-scale perspective via the Eriksen number $\text{Er} = \eta v_a R_g / K$, which compares the active stress to the elastic stress of the nematic over a polymer-sized region [104,105]. Physically, Er characterizes whether active flows v_a are sufficiently large to disrupt nematic alignment on the scale of the polymer R_g . The Eriksen number is found to connect the timescale-based numbers as

$$\text{Er} = \frac{\text{Wi}}{\text{Pe}} = \frac{R_g}{l_a}. \quad (13)$$

Thus, Er is the ratio of characteristic lengths. The ratio of lengths is precisely what is employed as a dimensionless activity number in studies of confined active nematics [24,106–114] and Eq. (13) shows that the dimensionless activity numbers commonly used in the literature can be interpreted as Ericksen numbers. This relation highlights that these common dimensionless groups encode the same underlying physics of competition between activity and elasticity. Therefore, the effective diffusion can be rewritten as

$$\frac{D_e}{D_0} - 1 = \frac{\text{Wi}}{d\text{Er}^2}. \quad (14)$$

In fact, effective diffusivity can be directly written as a ratio of time and length scales as

$$\frac{D_e}{D_0} - 1 = \frac{1}{d} \left(\frac{\tau_r}{\tau_p} \right) \left(\frac{l_a}{R_g} \right)^2. \quad (15)$$

These expressions unify the scaling arguments and highlight that, no matter how it is expressed, enhanced diffusion arises from a balance between both the time- and length scales of the polymer and active fluid.

This Letter investigated the dynamics of single flexible polymers embedded in active nematic turbulence, finding that activity significantly enhances polymer diffusion. The enhancement arises from the short-time advection of the polymer and decorrelation of active turbulence. The advection rate for shorter polymers increases with activity and is directly proportional to the characteristic velocity of the active turbulence, whereas longer polymers reach a plateau. From the ratio of Péclet number, which reveals the scaling for advection speed, and Weissenberg number, which characterizes the shear, the scaling for effective diffusion of polymers can be predicted.

The scaling relation $D_e/D_0 - 1 \sim \mathcal{D}_N \mathcal{S}_{N,\zeta}$ demonstrates that, by selecting polymer length N and activity ζ , one can tune the enhancement of diffusivity by orders of magnitude. This provides a direct “design rule” for engineering active-polymer composites: targeting a desired effective diffusion D_e in 2D can be achieved by solving $D_e/D_0 - 1 \sim \mathcal{D}_N \mathcal{S}_{N,\zeta}$ for the required N at fixed ζ or vice versa.

Experiments have recently demonstrated DNA experiences enhanced transport in dense suspensions of swimming bacteria [115]. To verify our predictions for the dependence of polymer diffusivity across a broad range of polymer lengths, DNA molecules of varying contour lengths, including linearized plasmids [116], fosmid and bacterial artificial chromosome (BAC)-DNAs [117,118], λ -DNA constructs [118], and longer genomic fragments [119–121] could feasibly be suspended in bacterial suspensions to measure the effective diffusivity. Another approach might be to suspend flexible colloidal-bead chains in active bacterial suspensions [122]. Experimentally verifying these predictions through DNA/bacteria composites would open the path for future active/polymer hybrids, such as embedding fibers in microtubule/kinesin films or intermediate filaments in actomyosin. This Letter sets the groundwork for future research on polymer dynamics in active media and may have direct consequences for macromolecular transport in bacterial swarms [55,56,123].

Studies of the dynamics of polymers suspended in classical inertial turbulence have revealed a wide array of dynamics [124–127] and the present Letter suggests a similar variety of dynamics may occur in active turbulence. The success of scaling arguments suggests that integrating polymer physics with active matter could provide valuable insights into biological systems where active forces influence polymer dynamics. For instance, recent studies of long polymers confined in spherical cavities with active nematics have helped to illuminate chromatin dynamics [39]. Similarly, dense suspensions of polymers in active fluids can serve as a model for cytoskeletal assemblies [128–131], polymer brushes in contact with active surroundings may model for biosensors such as protein channels [132–135] and polymer gels in an active solvent may represent the cell cortex [136–139]. Thus, studying the interplay between the polymeric degrees of freedom and active media can

reveal the pathways through which biological systems orchestrate their functions.

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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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