We investigate the aggregation and phase separation of thin, living *T. Tubifex* worms that behave as active polymers. Randomly dispersed active worms spontaneously aggregate to form compact, highly entangled blobs, a process similar to polymer phase separation, and for which we observe power-law growth kinetics. We find that the phase separation of active polymer-like worms does not occur through Ostwald ripening, but through active motion and coalescence of the phase domains. Interestingly, the growth mechanism differs from conventional growth by droplet coalescence: the diffusion constant characterizing the random motion of a worm blob is independent of its size, a phenomenon that can be explained from the fact that the active random motion arises only from the worms at the surface of the blob. This leads to a fundamentally different phase-separation mechanism, that may be unique to active polymers.

Driven colloidal particles [1,3], self-propelled bots [4, 5], cells [6, 7], animals [8] and humans [9] belong to the broad field of active matter: a class of non-equilibrium entities composed of many interacting units that individually consume energy and which can at larger scales collectively generate motion or mechanical stresses [2,10,11]. Their collective behavior is fascinating: whether it is flocks of birds, schools of fish, or ants collaborating to survive, the activity and interactions of the individual components give rise to highly non-trivial macroscopic phenomena [12]. Here, we investigate the spontaneous aggregation and eventual phase separation of active-polymer-like living worms. At first sight, the phase separation of active polymers may seem similar to that of solutions [13,14] and crystallizing solids [15], for which detailed theories are available: in both cases, the aggregating particles move randomly and tend to stick together when they are in close proximity. As such, one might naively expect the phase separation of active particles to involve a mechanism similar to Ostwald ripening [16], where the aggregation is driven by the combined effect of diffusion of the aggregating particles and surface tension of the aggregates. However, recent work has shown that the phase separation of active particles can involve unique mechanisms that rely completely on the activity [2, 3, 6, 12, 17, 18], and our results indicate that this also holds for the phase separation of active polymers.

We investigate the phase separation of living *T. Tubifex* worms in water. These worms are active swimmers, and approximately 300 µm thick and 10–40 mm long (Supplementary Information Discussion I.1). The thermal random motion of the worms (estimated from the Stokes-Einstein equation) is negligible compared to their active motion, so they constitute a good model system for active polymers [19–24]. When randomly distributed over a volume of water, the worms aggregate spontaneously (Fig. 1) into highly entangled ‘blobs’. The biological function of this aggregation is to minimize exposure to dissolved oxygen, high levels of which are poisonous to *Tubifex* [25]. We find that we can analyze the aggregation with relatively simple methods, which makes these living worms an excellent system to investigate the phase separation of active polymers.

In the experiments, we disperse a specific number of worms in a thermostated water volume and observe their aggregation in real time by recording videos (see Supplementary Information). We investigate two geometries: in the simplest, the worms are dispersed in a 25 × 25 × 2.5 cm volume of water. In this geometry worm motion is effectively two-dimensional (2D), since the worms are denser than water and therefore always located at or close to the bottom of the water volume (Supplementary Information Fig. S2). In other experiments the water volume is a 51 × 1 × 1 cm channel in which the

![FIG. 1. Aggregation and phase separation of *T. Tubifex* worms.](image-url)
worms are confined in an effectively one-dimensional space (1D). Figures 1a–c show snapshots from a typical 2D experiment; Fig. 1d from a typical 1D experiment (Supplementary Information videos 1 and 2). In both geometries, as time progresses, the worms form ever larger aggregates until after about 1 hour all worms are concentrated into a few large aggregates, which have the shape of a slightly flattened sphere (Supplementary Information Discussion II.1). This shape is a compromise between minimum exposed surface and minimum gravitational energy. In the following, we refer to these worm aggregates as ‘blobs’.

The aggregation of the active worms seems similar to that generally observed for polymer phase separation. Hence, one might expect that the aggregation occurs through Ostwald ripening, in which larger aggregates grow at the cost of smaller ones [26]. This mechanism is driven by the reduction of the total surface tension with increasing average blob size, and we do in fact measure a finite surface tension for the worm blobs (see Supplementary Information). However, closer inspection (Figures 1e–h) shows that blobs of all sizes are growing, and that the growth does not occur by Ostwald ripening, but rather by the merging of smaller aggregates into larger ones. Such growth by coalescence of diffusing droplets has been investigated previously in the context of the growth of vapor-deposited thin films and of droplets on a surface [27–29].

To quantify our observations, we characterize the average blob size $\langle R \rangle$ as a function of time by taking the 2D Fourier transform of the images and determining the wavevector magnitude $q_{\text{max}}$, at which the spectrum reaches its maximum intensity (see Figure 2a for an example). With increasing time, $q_{\text{max}}$ shifts to lower values, corresponding to an increasing average blob size $\langle R \rangle$. Figure 2b shows the average blob size $q_{\text{max}}$ determined in this way as a function of time, at several temperatures. We find that the growth exhibits power law behavior with $\langle R_{2D} \rangle \sim t^{1/3}$ in the 2D experiments, and $\langle R_{1D} \rangle \sim t^{1/4}$ in the 1D experiments, independently of the initial worm concentration (Supplementary Information Fig. S3). In addition, we observe in the 2D experiments that the power-law growth behavior is independent of temperature.

We also investigate the motion of individual worms by recording image sequences of single, isolated worms at different temperatures (Supplementary Information Discussion I). Figure 2c shows an example of the centre-of-mass trajectory of an individual worm at various temperatures. The motion is a random walk with an effective diffusion constant...
that increases with temperature, as is confirmed by extracting the mean square displacement from which we retrieve the diffusion coefficient (Supplementary Information Discussion I.2), see Fig. 2. We find that by rescaling the time axis \( \tau \propto D_{\text{warm}}t \) all the growth curves collapse onto a single master curve \( \propto \tau^{1/3} \) (Fig. 2), confirming that the aggregation kinetics is determined by the random motion of the worms.

To shed more light on the observed power-law growth kinetics, we perform computer simulations of 2D growth by coalescence of diffusing droplets, using an approach similar to that of Ref. [28] (Supplementary Information Discussion III). We assume that the blobs are spherical and move randomly in a 2D space, and that two blobs with radii \( R_1 \) and \( R_2 \) at positions \( r_1 \) and \( r_2 \) coalesce when \( |r_1 - r_2| < R_1 + R_2 \) to form a new blob with radius \( R = \sqrt{R_1^2 + R_2^2} \) at position \( r = (R_1^2 r_1 + R_2^2 r_2) / (R_1^2 + R_2^2) \). The simulation starts with a random distribution of mono-disperse spheres with radius 1 which represent the individual worms, in a 2D space of size 200×200. At every time step, each droplet is moved in a random direction by a distance equal to the diffusion constant \( D_{\text{blob}} \).

**FIG. 3. Simulation of phase-separation dynamics.** a Snapshots of simulated blob growth by coalescence of randomly moving spherical blobs, with blob diffusion constant \( D_{\text{blob}} \) inversely proportional to the blob radius (top), and independent of blob radius (bottom). b Average blob radius \( \langle R \rangle \) as a function of time obtained from the simulations. In both cases the growth follows a power law (indicated by lines), with exponents of 0.15 and 0.3 respectively.

**FIG. 4. Blob diffusion.** Effective diffusion constant as a function of average blob size \( \langle R_{\text{blob}} \rangle \) at \( T=20^\circ C \) as determined from the slopes of experimental MSDs (see Supplementary Information). The error bars are mostly due to sample-to-sample variability. The purple line shows the expected scaling for particles undergoing Brownian random motion \( (D_{\text{blob}} \propto \langle R_{\text{blob}} \rangle^{-1}) \). The experimental data (blue symbols) indicate a diffusion constant independent of blob size (dotted line). Lower insets: photographs of blobs of different sizes from which we measured the diffusion coefficient.
worm blobs differs fundamentally from conventional, Brownian random motion. In the latter case, \(D_{\text{blob}} = k_B T / 6 \pi \eta R_{\text{blob}}\), with \(k_B\) Boltzmann’s constant and \(\eta\) the viscosity of the liquid surrounding the particle [26, 30]. The difference can be explained by considering the origin of the random motion: Brownian random motion of a particle is caused by the thermal motion of the surrounding molecules, whereas the random motion of a worm blob is due to the active motion of the constituent worms, and this can lead to a different dependence of \(D_{\text{blob}}\) on \(R_{\text{blob}}\).

To rationalize the observed size-independent blob-diffusion constant for active worms, let us assume that an individual worm exerts a swimming force of fixed magnitude \(|F_0|\) in a direction that varies randomly with a correlation time \(\tau\). Since the worms inside the entangled blob are effectively immobilized, only the worms at the outer surface of a blob contribute swimming force. The number of worms at the surface of a blob of radius \(R_{\text{blob}}\) is \(N_{\text{surf}} \sim R_{\text{blob}}^2\), so the total random force exerted by these \(N_{\text{surf}}\) worms has an average magnitude \(\langle |F_{\text{drive}}| \rangle \sim |F_0| \sqrt{N_{\text{surf}}} \sim |F_0| R_{\text{blob}}^2 \) (and a correlation time \(\tau\)). The drag force on the blob as a function of speed \(v\) is given by Stokes’ law, \(F_{\text{drag}} \sim v R_{\text{blob}}\). To obtain the steady-state speed \(v\) we equate the driving and drag forces, and obtain a velocity \(\nu\) that is independent of the blob radius \(R_{\text{blob}}\). Assuming that the random walk of the blobs occurs by random steps in which it has approximately the steady-state speed, this result implies that a blob-diffusion constant that is independent of blob size.

Again simulating the blob growth, but now using a size-independent blob diffusion constant \((D_{\text{blob}} = 1\)), we obtain a growth exponent of \(\sim 0.3\) (red points and line in Fig. 3b), in good agreement with the experimentally observed value. Thus a size-independent diffusion constant for the random motion of the blobs explains the observed power-law growth of the worm blobs (Figs. 2 and 3), at least in the 2D experiments. In the 1D experiments, the situation is complicated by the fact that blob motion slows down when the blob size becomes comparable to the channel width, possibly due to friction at the glass/worm surface. This makes quantitative analysis of the blob diffusion difficult. However, the fact that in this case the diffusion constant does decrease with \(R_{\text{blob}}\), and that at the same time the growth exponent is closer to that predicted by the conventional droplet-coalescence model does provide a qualitative confirmation of the above ideas.

To conclude, the active motion of \(T.\) Tubifex worms leads to a phase-separation mechanism that is fundamentally different from that of normal polymers. It involves growth by blob coalescence rather than Ostwald ripening, and with power-law behavior that is different from that of conventional growth by diffusing-droplet coalescence. The observed phase-separation mechanism and power-law kinetics seem to be due only to the active nature of the living worms and to the immobilization of the worms at the inside of the entangled blobs; it may therefore be a generic phenomenon, also occurring in other types of polymeric active matter. We hope that the results presented here will stimulate further experimental and theoretical work in this direction.

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[1] J. Palacci, S. Sacanna, A. P. Steinberg, D. J. Pine, and P. M. Chaikin, Science 339, 936 (2013).
[2] D. Geyer, D. Martin, J. Tailleur, and D. Bartolo, Phys. Rev. X 9, 031043 (2019).
[3] M. N. van der Linden, L. C. Alexander, D. G. Aarts, and O. Dauchot, Phys. Rev. Lett. 123, 098001 (2019).
[4] G. Briand and O. Dauchot, Phys. Rev. Lett. 117, 098004 (2016).
[5] A. Deblais, T. Barois, P. Delville, R. Vaudaine, J. S. Lintuuvori, J. Boudet, J. Barei, and H. Kelley, Phys. Rev. Lett. 120, 188002 (2018).
[6] J. Schwarz-Lineik, C. Valeriani, A. Cacciuto, M. E. Cates, D. Marenduzzo, A. N. Morozov, and W. C. K. Poon, Proc Natl Acad Sci USA 109, 4052 (2012).
[7] G. Duclos, C. Blanch-Mercader, V. Yashunsky, G. Salbreux, J.-F. Joanny, J. Prost, and P. Silberzan, Nature Physics 14, 728 (2018).
[8] M. Tennenbaum, Z. Liu, D. Hu, and A. Fernandez-Nieves, Nat. Mat. 15, 54 (2015).
[9] N. Bain and D. Bartolo, Science 363, 46 (2019).
[10] T. B. Liverpool, A. C. Maggs, and A. Ajdari, Phys. Rev. Lett. 86, 4171 (2000).
[11] T. Speck, J. Bialké, A. M. Menzel, and H. Löwen, Phys. Rev. Lett. 112, 218304 (2014).
[12] C. Bechinger, R. D. Leonardo, A. Cacciuto, and M. E. Cates, Rev. Mod. Phys. 88, 045006 (2012).
[13] A. Dobry and F. Boyer-Kawenoki, Journal of Polymer Science 2, 90 (1947).
[14] H. Tanaka and Y. Nishikawa, Phys. Rev. Lett. 95, 078103 (2005).
[15] G. Olson and J. Roth, Materials Science Reports 3, 1 (1988).
[16] P. W. Voorhees, Annu. Rev. Mater. Sci. 22, 197 (1992).
[17] M. E. C. A. G. Thompson, J. Tailleur and R. A. Blythe, J. Stat. Mech. 2, 02029 (2011).
[18] M. E. Cates and J. Tailleur, Annu. Rev. Condens. Matter Phys. 6, 219 (2015).
[19] R. G. Winkler, J. Elgeti, and G. Gompper, J. Phys. Soc. Jpn. 86, 101014 (2017).
[20] A. Martín-Gómez, G. Gompper, and R. G. Winkler, Polymers 2018, 0837 (2018).
[21] V. Bianco, E. Locatelli, and P. Malgaretti, Phys. Rev. Lett. 121, 217802 (2018).
[22] A. Martín-Gómez, T. Eisenstein, G. Gompper, and R. G. Winkler, Soft Matter 15, 3957 (2019).
[23] Z. Mokhtari and A. Zippelius, Phys. Rev. Lett. 123, 028001 (2019).
[24] A. Deblais, S. Woutersen, and D. Bonn, “Rheology of active-polymer-like T. Tubifex worms,” (2019), arXiv:1910.12301
[25] J. G. Walker, Biol. Bull. 138, 235 (1971).
[26] R. A. L. Jones, *Soft Condensed Matter* (Oxford University Press, Oxford, 2002).
[27] F. Family and P. Meakin, Phys. Rev. Lett. 61, 428 (1988).

[28] A. Steyer, P. Guenon, D. Beyens, and C. M. Knobler, Pys. Rev. A 44, 8271 (1991).
[29] P. Meakin, Rep. Prog. Phys. 55, 157 (1992).
[30] A. Einstein, Ann. Phys. 17, 549 (1905).