Supporting Text

1 Dynamic Structure Factor

For the fits to the dynamic light-scattering data we use an expression for the dynamic structure factor derived from the single-polymer transverse mean square displacement (MSD) of a point on the contour (1). By a normal mode analysis as outlined in *Methods*, one finds for the MSD (1–3):

$$\delta r_\perp^2(t) = \langle [\mathbf{r}_\perp(s,t) - \mathbf{r}_\perp(s,0)]^2 \rangle \quad \textbf{[5]}$$

$$\approx \frac{4L^3}{\pi^4 \ell_p} \int_0^{L/a} \mathrm{d}n \frac{1 - \exp[-(t/\tilde{\tau}_n)\tilde{h}(k_n)]}{n^4}, \quad \textbf{[6]}$$

where $k_n = \pi n/L$, $\tau_n = 4\pi\eta/\kappa k_n^4$, and

$$\tilde{\tau}_n = \begin{cases} \tau_n & (n > \ell), \\ \tau_n \exp[\mathcal{E}(l/n - 1)] & (n < \ell), \end{cases} \qquad (\ell = L/\Lambda) \quad \textbf{[7]}$$

is the relaxation time of the glassy wormlike chain and $\tilde{h}(k) = 4\pi\eta h(k)$ is the Fourier transform of the ("Rotne-Prager") mobility function,

$$\tilde{h}(k) = \gamma_{RP} - \log(ka) + \mathcal{O}(k^2 a^2), \quad \textbf{[8]}$$

which represents a dimensionless refining factor to the simple "free-draining" approximation with a constant friction coefficient $\zeta_\perp$ employed in *Methods*. This factor depends on the hydrodynamic constant γ_{RP} and on the backbone diameter a of the filament.

For $n > \ell$ the following approximation to the integral Eq. 6 is used:

$$\delta r^2_{\perp,\Lambda}(t) = \frac{4\Lambda^3}{3\ell_p\pi^4}\Bigg[1-\frac{3}{4}\left(\frac{t}{\tau_\Lambda}\right)^{3/4}\left\{-\log\left[e^{-B}a/\ell_\perp(t)\right]\right\}^{3/4} \Gamma\left(-3/4,(t/\tau_\Lambda)\left\{-\log\left[e^{-B}a/\ell_\perp(t)\right]\right\}\right)\Bigg]. \quad [9]$$

Here, $\tau_\Lambda \equiv \tau_\ell = (4\pi\eta/\kappa)(\Lambda/\pi)^4$ is the relaxation time of a mode of (half) wavelength Λ. The length $\ell_\perp(t) = (\kappa t/4\pi\eta)^{1/4}$ is the transverse elastohydrodynamic correlation length. The constant B is defined as $B = \gamma_{RP} - \log(\tilde{n}_0)$ where $\tilde{n}_0$ is a numerically determined mode number. The (approximated) integral for $n < \ell$,

$$r^{2,G}_\perp(t) = \frac{4L^3}{\ell_p\pi^4}\int_0^\ell \mathrm{d}n \frac{1-\exp\left(-(t/\tau_1)n^4\exp[-\mathcal{E}(l/n-1)]\left\{B-\log\left[a/\ell_\perp(t)\right]\right\}\right)}{n^4} \quad [10]$$

is evaluated numerically (an analytic approximation is given in ref. 4). The dynamic structure factor in the limit $t \gg \tau_q = 4\pi\eta/\kappa q^4$ immediately follows (1),

$$S(q,t)/S(q,0) = \exp(-q^2[\delta r^{2,G}_\perp(t) + \delta r^2_{\perp,\Lambda}(t)]/4). \quad [11]$$

To produce the fits of Eq. 11 to the DLS data in Fig. 5 of the main text, we used the constants $\ell_p = 9$ µm and $a = 9$ nm (5). The free parameters were the stretching parameter $\mathcal{E}$, the interaction length Λ, and the hydrodynamic constant B. The values obtained for $c = 17$ µM and different values of the scattering vector q are given in SI Figs. 7-9.

The applicability of single-polymer theory (1) becomes questionable for low values of q comparable to the inverse mesh size ξ^{-1}. As $\xi^{-1} = 2.6$ µm^{-1} for $c = 17$ µM, a quantitative evaluation of the fits should focus on the largest measurable $q \gg 2.6$ µm^{-1}. The data in SI Figs. 7-9 suggest that (at least) the values obtained for $q \lesssim 10\ \mu\text{m}^{-1}$ have to be discarded as meaningless, while those obtained for $10\ \mu\text{m}^{-1} \lesssim q \lesssim 15\ \mu\text{m}^{-1}$ still exhibit some small but noticeable systematic errors.

The MSD given by Eq. 6 exhibits a remarkable symmetry. It is straightforward to check that upon simultaneously rescaling $\Lambda \to \Lambda' = \gamma\Lambda$ and $\mathcal{E} \to \mathcal{E}' = \gamma\mathcal{E}$, the long-time tails of the original MSD $\delta r_\perp^2$ and the MSD of the rescaled variables, $\delta r_\perp'^2$, can be superimposed, that is, $\delta r_\perp'^2(t) = \delta r_\perp^2(\alpha t)$ for $t \gg \tau_\Lambda, \tau_{\Lambda'}$ with $\alpha = \exp[\mathcal{E}(\gamma - 1)]$ (the weak mode-number dependence of the mobility function is neglected). Inset of Fig. 5 in the main text demonstrates that this symmetry is well obeyed for $q = 8.04\ \mu\text{m}^{-1}$.

2 Linear Viscoelastic Modulus of a Glassy wormlike Chain

In the theory of Soft Glassy Rheology (SGR) (6), the noise temperature $1 < x < 2$ is directly monitored by the power-law exponent $x - 1$ of the linear viscoelastic moduli for low frequencies ($G'(\omega), G''(\omega) \sim \omega^{x-1}$). In this section, x will be compared to the stretching parameter $\mathcal{E}$ of a glassy wormlike chain (GWLC).

By applying the the prescription of the GWLC to the high-frequency limiting form of the dynamical shear modulus of a wormlike chain (7), we calculate the macrorheological modulus of a GWLC (4). Its expression for vanishing prestress is

$$G(\omega) = \frac{1}{5}\Lambda/\xi^2\alpha(\omega), \qquad \textbf{[12]}$$

here $\xi = \sqrt{3/c_p L}$ is the mesh size and

$$\alpha(\omega) = \alpha_L \sum_{n=1}^{\infty} \frac{1}{n^4 + i\omega\tilde{\tau}_n/2} \qquad \textbf{[13]}$$

with the prefactor $\alpha_L = L^4/(kT\pi^4\ell_p^2)$ is the susceptibility of the GWLC to a point force at the ends. The relaxation times $\tilde{\tau}_n$ are modified as described in section 1. By an analytic approximation it is possible to determine the functional dependence of $G(\omega)$ on the frequency and the stretching parameter $\mathcal{E}$. The viscoelastic modulus of a GWLC is not a simple power law but a function that depends essentially logarithmically on frequency. An approximation to the storage modulus valid for $\omega\tau_\Lambda \ll 1$, $\mathcal{E} \gg 1$ is

$$G'(\omega) = \frac{\Lambda}{5\xi^2\alpha_\Lambda} \frac{3}{\left(1 - \frac{4}{\mathcal{E}}\log\left[\frac{4}{\mathcal{E}}\left(\frac{\omega\tau_\Lambda}{2}\right)^{1/4}\right] - \frac{4}{\mathcal{E}}\log\left\{\frac{\mathcal{E}}{4} - \log\left[\left(\frac{\omega\tau_\Lambda}{2}\right)^{1/4}\right]\right\}\right)^3} \qquad \textbf{[14]}$$

From Eq. 14 it is straightforward to derive the local power law exponent of the elastic modulus for $\mathcal{E} \gg 1$ at a fixed frequency. Asymptotically for $\mathcal{E} \to \infty$ (at fixed ω) the result is $x - 1 = 3/\mathcal{E}$ (note that the limits $\mathcal{E} \to \infty$ and $\omega \to 0$ do not commute). This result and the exact slope valid for all $\mathcal{E}$ are shown in SI Fig. 10.

A similar analysis can be carried out for the loss angle $\delta = \arctan(G''/G')$. The exact value as a function of $\mathcal{E}$ at a fixed frequency and the asymptotic analytical approximation $\delta = \arctan(5/\mathcal{E}) \approx 5/\mathcal{E}$ (valid for $\mathcal{E} \to \infty$ at $\mathcal{E}\omega/\omega_\Lambda =$ const. $\ll 1$) are given in SI Fig. 10. Our result ($5(x-1)/3 = \delta$) is compatible with power-law rheology, where the exact relation $(\pi/2)(x-1) = \delta$ is expected (8). The deviation of our factor from the exact value is an artefact of the analytical approximations made.

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