Supporting Information

Derivation of the Smoluchowski equation for FB model : According to the Liouville theorem ($d\varphi/dt = 0$) the time evolution of probability density $\varphi(x, r, t)$ in terms of x and r satisfies

$$\frac{\partial\varphi}{\partial t} = -\frac{\partial}{\partial x}\left(\frac{dx}{dt}\varphi\right) - \frac{\partial}{\partial r}\left(\frac{dr}{dt}\varphi\right). \quad \text{(S1)}$$

Insertion of two Langevin equations (Eq.(1) in the main text) for the fluctuating bottleneck model $\partial_t x = -\zeta^{-1}[\partial_x U_{\text{eff}}(x;r) + F_x(t)]$ and $\partial_t r = -\lambda r + F_r(t)$ into Eq.S1 leads to

$$\begin{aligned}\frac{\partial\varphi}{\partial t} &= \frac{\partial}{\partial x}\left(\frac{1}{\zeta}\frac{dU_{\text{eff}}(x)}{dx}\varphi\right) + \frac{\partial}{\partial r}(\lambda r\varphi) - \frac{\partial}{\partial x}\left(\frac{1}{\zeta}F_x(t)\varphi\right) - \frac{\partial}{\partial r}(F_r(t)\varphi) \\ &\equiv -\mathcal{L}\varphi - \frac{\partial}{\partial\vec{a}}\cdot\left(\vec{F}(t)\varphi\right)\end{aligned} \quad \text{(S2)}$$

where $\vec{a} \equiv (x, r)$ and $\vec{F}(t) \equiv (\frac{1}{\zeta}F_x(t), F_r(t))$. Using the vector notation as in the second line of Eq.S2, one can formally solve for the probability density $\varphi(\vec{a}, t)$ as

$$\varphi(\vec{a}, t) = e^{-t\mathcal{L}}\varphi(\vec{a}, 0) - \int_0^t ds e^{-(t-s)\mathcal{L}}\frac{\partial}{\partial\vec{a}}\cdot\left(\vec{F}(s)\varphi(\vec{a}, s)\right) \quad \text{(S3)}$$

Averaging $\varphi(\vec{a}, t)$ over noise after iterating $\varphi(\vec{a}, t)$ into the noise related term in the integrand and exploiting the fluctuation-dissipation theorem, we obtain the Smoluchowski equation for $\varphi(x, r, t)$ in the presence of a reaction sink, $\mathcal{S}(x, r) = k_r r^2\delta(x - x_{\text{ts}})$,

$$\frac{\partial\overline{\varphi}(x, r, t)}{\partial t} = [\mathcal{L}_x(x) + \mathcal{L}_r(r) - \mathcal{S}(x, r)]\,\overline{\varphi}(x, r, t), \quad \text{(S4)}$$

where $\mathcal{L}_x \equiv D\partial_x\left(\partial_x + (k_BT)^{-1}\partial_x U_{\text{eff}}(x)\right)$ and $\mathcal{L}_r \equiv \lambda\theta\partial_r(\partial_r + r/\theta)$. Integrating both sides of the equation over x by defining $\overline{C}(r, t) \equiv \int_{-\infty}^{\infty} dx\overline{\varphi}(x, r, t)$ leads to $\partial_t\overline{C} = \mathcal{L}_r\overline{C}(r, t) - k_r r^2\overline{\varphi}(x_{\text{ts}}, r, t)$. By setting $\overline{\varphi}(x_{\text{ts}}, r, t) = \phi_x(x_{\text{ts}})\overline{C}(r, t)$ where $\phi(x_{\text{ts}}) = e^{-U_{\text{eff}}(x_{\text{ts}})/k_BT}/\int dx e^{-U_{\text{eff}}(x)/k_BT} \approx \sqrt{U''_{\text{eff}}(x_{\text{b}})/2\pi k_BT}\,e^{-(U_{\text{eff}}(x_{\text{ts}})-U_{\text{eff}}(x_{\text{b}}))/k_BT}$, we get

$$\partial_t\overline{C}(r, t) = \left[\lambda\theta\partial_r(\partial_r + r/\theta) - kr^2\right]\overline{C}(r, t), \quad \text{(S5)}$$

where $k \equiv k_r\sqrt{U''_{\text{eff}}(x_{\text{b}})/2\pi k_BT}\,e^{-\Delta U^{\ddagger}/k_BT}$ with $\Delta U^{\ddagger} \equiv U(x_{\text{ts}}) - U(x_{\text{b}})$. In all likelihood, $k_r\left(= D\times\sqrt{U''_{\text{eff}}(x_{\text{ts}})/2\pi k_BT}\right)$ represents the product of diffusion coefficient D associated with barrier crossing dynamics and the contribution of dynamics at the barrier top. Thus, under tension f, one can set $k \to k(f) = k_0 e^{f\Delta x^{\ddagger}/k_BT}$ where $k_0 \equiv (\xi D\sqrt{U''_{\text{eff}}(x_{\text{b}})U''_{\text{eff}}(x_{\text{ts}})}/2\pi k_BT)e^{-\Delta U^{\ddagger}/k_BT}$ and ξ describes the correction due to geometrical information of the cross section of bottleneck [1, 2]. Therefore, under tension f, Eq.S5 becomes Eq.(2) in the main text.

Solution of the Smoluchowski equation with time-dependent sink : For the problem with a constant loading rate, the sink function of our Smoluchowski equation becomes time-dependent, resulting in the following equation for the flux $\overline{C}(r, t)$,

$$\frac{\partial\overline{C}(r, t)}{\partial t} = \lambda\theta\frac{\partial}{\partial r}\left(\frac{\partial}{\partial r} + \frac{r}{\theta}\right)\overline{C}(r, t) - k_0 r^2 e^{t(\gamma\Delta x^{\ddagger}/k_BT)}\overline{C}(r, t) \quad \text{(S6)}$$

with $\overline{C}(r, t = 0) = \sqrt{\frac{2}{\pi\theta}}e^{-r^2/2\theta}$. Although a time-dependent sink term, in general, makes Smoluchowski equations analytically intractable, the ansatz $\overline{C}(r, t) \sim e^{\nu(t)-\mu(t)r^2}$ used in the Ref. [1] allows us to solve the above problem exactly. Substitution of $\overline{C}(r, t) \sim e^{\nu(t)-\mu(t)r^2}$ leads to two ODEs for $\nu(t)$ and $\mu(t)$ (with $'$ denoting derivative with respect to t),

$$\nu'(t) = -2\lambda\theta\mu(t) + \lambda \tag{S7}$$

and

$$\left(\mu(t) - \frac{1}{4\theta}\right)' = -4\lambda\theta\left(\mu(t) - \frac{1}{4\theta}\right)^2 + \frac{\lambda}{4\theta}\left(1 + \frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}}\right) \tag{S8}$$

with $\mu(0) = 1/2\theta$. The equation for $\mu(t)$ in Eq.S13 is the Riccati equation, $y' = q_0(t) + q_1(t)y + q_2(t)y^2$ with $y(t) \equiv \mu(t) - 1/4\theta$. In general, the Riccati equation can be reduced to a second order ODE. The variable is changed in two steps : (i) $v(t) = q_2(t)y(t)$ leads to $v' = v^2 + P(t)v + Q(t)$ where $Q = q_0q_2 = -\lambda^2\left(1 + \frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}}\right)$ and $P = q_1 + q_2'/q_2 = 0$. (ii) Another substitution $v(t) = -u'(t)/u(t)$ leads to $u''(t) - P(t)u'(t) + Q(t)u(t) = 0$, i.e.,

$$u''(t) - \lambda^2\left(1 + \frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}}\right)u(t) = 0. \tag{S9}$$

Introducing the variable $\rho = \frac{2\lambda}{\tilde{\gamma}}\sqrt{\frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}/2}} = \beta\kappa(t)$ with $\beta \equiv \frac{2\lambda}{\tilde{\gamma}}$ and $\kappa(t) \equiv \sqrt{\frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}/2}}$ one can modify the second-order ODE in Eq.S9 into a more familiar modified Bessel equation,

$$\rho^2U_{\rho\rho} + \rho U_\rho - \left[\beta^2 + \rho^2\right]U = 0 \tag{S10}$$

where $u(t) = U(\rho)$. The solution of the above ODE is the linear combination of $I_{\pm\beta}(\rho)$ for non-integer β, and the linear combination of $I_\beta(\rho)$ and $K_\beta(\rho)$ when β is integer. Thus, the solution of Eq.S10 is

$$U(\rho) = \begin{cases} c_1I_\beta(\rho) + c_2I_{-\beta}(\rho) & \beta \neq n,\ \beta > 0 \\ c_1I_\beta(\rho) + c_2K_\beta(\rho) & \beta = n \text{ where } n = 0, 1, 2\cdots \end{cases} \tag{S11}$$

For simplicity, we use the notation $\mathcal{Q}_\beta(\rho)$ to represent either $I_{-\beta}(\rho)$ or $K_\beta(\rho)$,

$$\mathcal{Q}_\beta(\rho) = \begin{cases} I_{-\beta}(\rho) & \beta \neq n,\ \beta > 0 \\ K_\beta(\rho) & \beta = n \text{ where } n = 0, 1, 2\cdots \end{cases} \tag{S12}$$

Thus one obtains $\mu(t)$ using $y(t) = -\frac{u'(t)}{q_2(t)u(t)}$.

$$\mu(t) = \frac{1}{4\theta} + \frac{\kappa(t)}{4\theta}\left(\frac{I_\beta'(\rho) + c\mathcal{Q}_\beta'(\rho)}{I_\beta(\rho) + c\mathcal{Q}_\beta(\rho)}\right). \tag{S13}$$

Note that $I_\beta'(\rho) \equiv dI_\beta(\rho)/d\rho$. The initial condition $\mu(0) = 1/2\theta$ determines the constant c in Eq.S13

$$c = \frac{I_\beta'(\rho_0) - [\kappa(0)]^{-1}I_\beta(\rho_0)}{[\kappa(0)]^{-1}\mathcal{Q}_\beta(\rho_0) - \mathcal{Q}_\beta'(\rho_0)}. \tag{S14}$$

where $\rho_0 \equiv \beta\kappa(0)$. Thus, one obtains

$$\frac{\mu(t)}{\mu(0)} = \frac{1}{2}\left[1 + \kappa(t)\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)}\right] \tag{S15}$$

where $\mathcal{I}(\rho) \equiv \left(I_\beta'(\rho_0)\mathcal{Q}_\beta(\rho) - \mathcal{Q}_\beta'(\rho_0)I_\beta(\rho)\right) - [\kappa(0)]^{-1}\left\{I_\beta(\rho_0)\mathcal{Q}_\beta(\rho) - \mathcal{Q}_\beta(\rho_0)I_\beta(\rho)\right\}$. Recall that $\rho \equiv \beta\kappa(t)$ with $\beta \equiv 2\lambda/\tilde{\gamma}$, $\kappa(t) \equiv \sqrt{\frac{4k_0\theta}{\lambda}e^{t\tilde{\gamma}/2}}$, and $\rho_0 \equiv \rho(0)$. Note that $\kappa(0)(\mathcal{I}'(\rho_0)/\mathcal{I}(\rho_0)) = 1$ is satisfied. Integration of Eq.S7 with t using Eq.S15 and change of variable $d\rho = \frac{\beta\tilde{\gamma}}{2}\kappa(t)dt = \lambda\kappa(t)dt$ results in the expression for $\nu(t)$:

$$\nu(t) = \frac{\lambda t}{2} - \frac{1}{2}\log\left(\frac{\mathcal{I}(\rho)}{\mathcal{I}(\rho_0)}\right). \tag{S16}$$

With $\mu(t)$ (Eq.S15) and $\nu(t)$ (Eq.S16) in hand, we can solve

$$\overline{C}(r,t) = \sqrt{\frac{2}{\pi\theta}}\left[\frac{\mathcal{I}(\rho)}{\mathcal{I}(\rho_0)}\right]^{-1/2}\exp\left[\frac{\lambda t}{2} - \frac{r^2}{4\theta}\left\{1 + \kappa(t)\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)}\right\}\right], \tag{S17}$$

from which the survival probability is obtained as

$$\Sigma_\lambda^\gamma(t) = \int_0^\infty dr\overline{C}(r,t) = \frac{1}{\sqrt{2\theta}}\frac{e^{\nu(t)}}{\sqrt{\mu(t)}} = \sqrt{2}e^{\frac{\lambda t}{2}}\left[\frac{\mathcal{I}(\rho)}{\mathcal{I}(\rho_0)}\right]^{-1/2}\left[1 + \kappa(t)\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)}\right]^{-1/2}. \tag{S18}$$

The λ-dependent unbinding time distribution $P_\lambda(t)$ are obtained from the relation $P_\lambda(t) = -d\Sigma_\lambda^\gamma(t)/dt$,

$$P_\lambda(t) = \frac{\lambda e^{\lambda t/2}}{\sqrt{2}}\left[\kappa^2(t)\frac{\mathcal{I}''(\rho)}{\mathcal{I}(\rho)} + \frac{1}{\beta}\kappa(t)\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)} - 1\right]\left[\frac{\mathcal{I}(\rho)}{\mathcal{I}(\rho_0)}\right]^{-1/2}\left[1+\kappa(t)\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)}\right]^{-3/2}. \tag{S19}$$

Transformation to the unbinding force distribution $P_\lambda(\tilde{f})\left[=\tilde{\gamma}^{-1}P_\lambda(t)\right]$ is made through the relationship between dimensionless scaled-force ($\tilde{f}$) and time t: $\tilde{f} = \tilde{\gamma}t$ with $\tilde{\gamma} = \gamma\Delta x^\ddagger/k_BT$.

Illustration using synthetic data : Although $P_\lambda(\tilde{f})$ in Eq.S19 is complicated, the familiar expression used in the Dynamic Force Spectroscopy (DFS) for $P(f)$ is restored when $\lambda \to \infty$ (see below). In order to obtain insight into the behavior of $P_\lambda(\tilde{f})$ we generated several synthetic distributions for varying λ values and loading rates. We find that $P_\lambda(\tilde{f})$ with varying $\tilde{\gamma}(=\gamma\Delta x^\ddagger/k_BT)$ shows the standard pattern of force distribution in DFS (Fig.S1-A, B) [2, 3]. The effect of varying λ on $P_\lambda(\tilde{f})$ is shown in Fig.S1-C, D. It is of particular interest that if $\tilde{\gamma} \gg k_0\theta$ then the most probable forces f_λ^* from $P_\lambda(\tilde{f})$ are insensitive to the variation in λ even though the shapes of $P_{\lambda\to 0}(\tilde{f})$ and $P_{\lambda\to\infty}(\tilde{f})$ are very different from each other (Fig.S1-C). However, when $\tilde{\gamma} \sim k_0\theta$, $\tilde{f}_\lambda^*$ changes with λ (Fig.S1-E) and the shape of $P_{\lambda\to 0}(\tilde{f})$ differs from $P_{\lambda\to\infty}(\tilde{f})$ qualitatively (Fig.S1-D).

Asymptotic behavior at $\lambda/\tilde{\gamma} \to \infty$: To obtain the asymptotic behavior we will use the following uniform asymptotic expansion of the modified Bessel function for large orders ($\nu \to \infty$) [4].

$$\begin{aligned}
I_\nu(\nu z) &\sim \frac{1}{\sqrt{2\pi\nu}}\frac{e^{\nu\eta}}{(1+z^2)^{1/4}}\left(1+\mathcal{O}(\nu^{-1})\right)\\
K_\nu(\nu z) &\sim \sqrt{\frac{\pi}{2\nu}}\frac{e^{-\nu\eta}}{(1+z^2)^{1/4}}\left(1+\mathcal{O}(\nu^{-1})\right)\\
I'_\nu(\nu z) &\sim \frac{1}{\sqrt{2\pi\nu}}\frac{(1+z^2)^{1/4}}{z}e^{\nu\eta}\left(1+\mathcal{O}(\nu^{-1})\right)\\
K'_\nu(\nu z) &\sim -\sqrt{\frac{\pi}{2\nu}}\frac{(1+z^2)^{1/4}}{z}e^{-\nu\eta}\left(1+\mathcal{O}(\nu^{-1})\right)
\end{aligned} \tag{S20}$$

where $I'_\nu(\nu z) \equiv \frac{d}{d(\nu z)}I_\nu(\nu z)$ and $\eta = \sqrt{1+z^2} + \log\left(\frac{z}{1+\sqrt{1+z^2}}\right)$

The asymptotic behavior at large negative orders can be obtained by using the relation $I_{-\nu}(z) = \frac{2}{\pi}\sin(\nu\pi)K_\nu(z) + I_\nu(z)$

$$\begin{aligned}
I_{-\nu}(\nu z) &\sim \left(\frac{2}{\sqrt{2\pi\nu}}\sin(\nu\pi)\frac{e^{-\nu\eta}}{(1+z^2)^{1/4}} + \frac{1}{\sqrt{2\pi\nu}}\frac{e^{\nu\eta}}{(1+z^2)^{1/4}}\right)\left(1+\mathcal{O}(\nu^{-1})\right)\\
I'_{-\nu}(\nu z) &\sim \left(-\frac{2}{\sqrt{2\pi\nu}}\sin(\nu\pi)\frac{(1+z^2)^{1/4}}{z}e^{-\nu\eta} + \frac{1}{\sqrt{2\pi\nu}}\frac{(1+z^2)^{1/4}}{z}e^{\nu\eta}\right)\left(1+\mathcal{O}(\nu^{-1})\right)
\end{aligned} \tag{S21}$$

Using these asymptotics, we obtain the following relations at $\beta = 2\lambda/\tilde{\gamma} \to \infty$.

$$\begin{aligned}
\lim_{\beta\to\infty}\mathcal{I}(\rho) &\sim 2\left(\frac{\sin\beta\pi}{\beta\pi}\right)\frac{1}{\kappa(0)}\left[e^{\beta(\eta-\eta_0)}\left(1+\frac{1}{S(t)}\right) + e^{-\beta(\eta-\eta_0)}\left(1-\frac{1}{S(t)}\right)\right]\\
\lim_{\beta\to\infty}\mathcal{I}'(\rho) &\sim 2\left(\frac{\sin\beta\pi}{\beta\pi}\right)\frac{1}{\kappa(0)\kappa(t)}\left[e^{\beta(\eta-\eta_0)}\left(S(t)+1\right) - e^{-\beta(\eta-\eta_0)}\left(S(t)-1\right)\right]
\end{aligned} \tag{S22}$$

where $S(t) \equiv (1+\kappa^2(t))^{1/2}$. Therefore

$$\lim_{\beta\to\infty}\frac{\mathcal{I}'(\rho)}{\mathcal{I}(\rho)} = \frac{S(t)}{\kappa(t)}\left[\frac{(S(t)+1)-(S(t)-1)e^{-2\beta(\eta-\eta_0)}}{(S(t)+1)+(S(t)-1)e^{-2\beta(\eta-\eta_0)}}\right] \tag{S23}$$

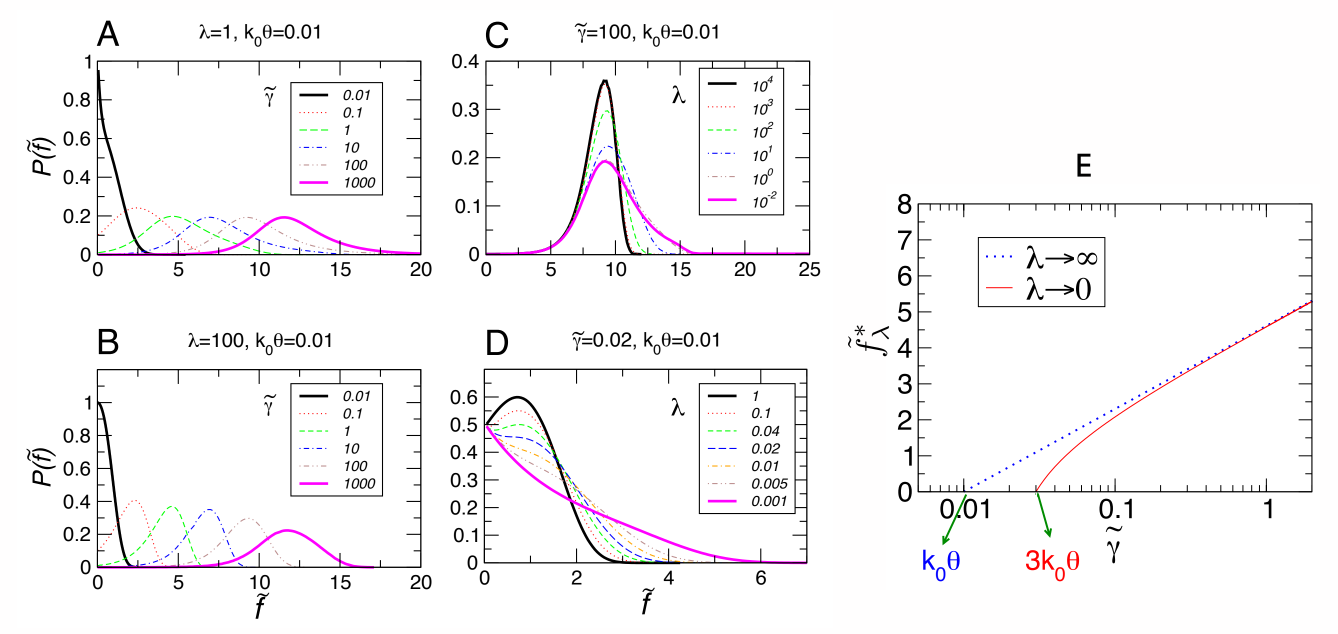

FIG. S1: **A-D** Rupture force distributions, $P(\tilde{f})$, under varying loading rates ($\tilde{\gamma}$) and the gating frequency (λ) characterizing the disorder. **E.** $\tilde{f}^*$ vs $\tilde{\gamma}$ plot under two limiting values of λ.

and

$$\lim_{\beta\to\infty}\frac{\mathcal{I}(\rho)}{\mathcal{I}(\rho_0)} = e^{\beta(\eta-\eta_0)}\left[\frac{(S(t)+1)-(S(t)-1)e^{-2\beta(\eta-\eta_0)}}{2S(t)}\right]. \tag{S24}$$

With $\lim_{\tilde{\gamma}\to 0} S(t) = S$ and $\lim_{\tilde{\gamma}\to 0}\beta(\eta-\eta_0) = \lambda St$ where $S \equiv \left(1+\frac{4k_0\theta}{\lambda}\right)^{1/2}$, it is now easy to show

$$\lim_{\tilde{\gamma}\to 0}\frac{\mu(t)}{\mu(0)} = \frac{1}{2}\left[1 + S\frac{(S+1)-(S-1)e^{-2\lambda St}}{(S+1)+(S-1)e^{-2\lambda St}}\right] \tag{S25}$$

and

$$\lim_{\tilde{\gamma}\to 0}\nu(t) = -\frac{\lambda t}{2}(S-1) + \log\left[\frac{(S+1)-(S-1)e^{-2\lambda St}}{2S}\right]^{-1/2}. \tag{S26}$$

Thus, substituting Eq.S25 and S26 into $\Sigma(t) = \int_0^\infty dr\overline{C}(r,t) = \frac{1}{\sqrt{2\theta}}\frac{e^{\nu(t)}}{\sqrt{\mu(t)}}$ recovers the previous result for survival probability in Zwanzig's FB model [1]

$$\lim_{\tilde{\gamma}\to 0}\Sigma(t) = \exp\left(-\frac{\lambda}{2}(S-1)t\right)\left[\frac{(S+1)^2-(S-1)^2E}{4S}\right]^{-1/2}. \tag{S27}$$

For $\lambda \to \infty$ and $\lambda \to 0$, $\lim_{\lambda\to\infty}\lim_{\tilde{\gamma}\to 0}\Sigma(t) = \exp(-k\theta t)$ and $\lim_{\lambda\to 0}\lim_{\tilde{\gamma}\to 0}\Sigma(t) = (1+2k\theta t)^{-1/2}$, respectively.

Survival probability ($\Sigma(\tilde{f})$) and rupture force distribution ($P(\tilde{f})$) for $\lambda\to\infty$ and $\lambda\to 0$: For $\lambda\to\infty$, taking $\int_0^\infty dr(\cdots)$ on Eq.S6 with pre-averaged rate constant $k(t)\theta$ and transforming t into $\tilde{f}$, we obtain $\tilde{\gamma}\partial_{\tilde{f}}\Sigma_{\lambda\to\infty}(\tilde{f}) = -k(\tilde{f})\theta\Sigma_{\lambda\to\infty}(\tilde{f})$, which leads to

$$\Sigma_{\lambda\to\infty}(\tilde{f}) = \exp\left[-\frac{1}{\tilde{\gamma}}\int_0^{\tilde{f}} d\tilde{f}k(\tilde{f})\theta\right] \tag{S28}$$

and the rupture force distribution ($P(\tilde{f}) = -d\Sigma(\tilde{f})/d\tilde{f}$)

$$P_{\lambda\to\infty}(\tilde{f}) = \frac{1}{\tilde{\gamma}}k(\tilde{f})\theta\Sigma_{\lambda\to\infty}(\tilde{f}) \quad \text{(S29)}$$

The most probable force ($\tilde{f}^*$) is obtained using $[\partial_{\tilde{f}}P_{\lambda\to\infty}]_{\tilde{f}=\tilde{f}^*} = 0$, which is equivalent to $\tilde{\gamma}[\partial_{\tilde{f}}k(\tilde{f})]_{\tilde{f}=\tilde{f}^*} = [k(\tilde{f})]^2_{\tilde{f}=\tilde{f}^*}\theta$. Using $k(\tilde{f}) = k_0e^{\tilde{f}}$, one can easily show that

$$\tilde{f}^*_{\lambda\to\infty} = \log\left[\tilde{\gamma}/(k_0\theta)\right]. \quad \text{(S30)}$$

This expression is equivalent to the standard DFS theory except for the presence of the θ term. The fast variation of r-coordinate effectively modifies the reactivity k_0r^2 into $k_0\theta$.

For $\lambda\to 0$ the bottleneck radius is quenched to a single value, say, r_0. In this case the noise-averaged probability of the molecule found at the configuration of r_0 at force $\tilde{f}$, $\overline{C}(r_0,\tilde{f}) = \exp\left(-\frac{1}{\tilde{\gamma}}\int_0^{\tilde{f}}d\tilde{f}k(\tilde{f})r_0^2\right)$, should be weighted by $\phi(r_0)\left[=\sqrt{\frac{2}{\pi\theta}}e^{-r_0^2/2\theta}\right]$ as $\Sigma_{\lambda\to 0}(\tilde{f}) = \int_0^\infty dr_0\overline{C}(r_0,\tilde{f})\phi(r_0)$ to give the survival probability,

$$\Sigma_{\lambda\to 0}(\tilde{f}) = \left[1+\frac{2\theta}{\tilde{\gamma}}\int_0^{\tilde{f}}d\tilde{f}k(\tilde{f})\right]^{-1/2}. \quad \text{(S31)}$$

A similar procedure as in Eqs. S29 and S30 leads to

$$P_{\lambda\to 0}(\tilde{f}) = \frac{1}{\tilde{\gamma}}k(\tilde{f})\theta\left[\Sigma_{\lambda\to 0}(\tilde{f})\right]^3 \quad \text{(S32)}$$

and

$$\tilde{f}^*_{\lambda\to 0} = \log\{(\tilde{\gamma}/k_0\theta)\left(1-2k_0\theta/\tilde{\gamma}\right)\}. \quad \text{(S33)}$$

Comparison between $P_{\lambda\to\infty}(\tilde{f})$ and $P_{\lambda\to 0}(\tilde{f})$ one-dimensional models : Asymptotic behaviors of $P(f)$ with two limiting λ values at large $\tilde{f}\gg\tilde{f}^*$, $P_{\lambda\to\infty}(\tilde{f})$ and $P_{\lambda\to 0}(\tilde{f})$ are obtained by using the Bell model for $k(\tilde{f})$. Comparison between $P_{\lambda\to\infty}(\tilde{f})$ and $P_{\lambda\to 0}(\tilde{f})$ can be made by using the explicit form of $k(\tilde{f}) = k_0e^{\tilde{f}}$.

$$P_{\lambda\to\infty}(\tilde{f}) = \frac{k_0\theta}{\tilde{\gamma}}\exp\left[\tilde{f}-\frac{k_0\theta}{\tilde{\gamma}}(e^{\tilde{f}}-1)\right] \quad \text{(S34)}$$

and

$$P_{\lambda\to 0}(\tilde{f}) = \frac{k_0\theta}{\tilde{\gamma}}\exp(\tilde{f})\left[1+2\frac{k_0\theta}{\tilde{\gamma}}(e^{\tilde{f}}-1)\right]^{-3/2}. \quad \text{(S35)}$$

For $\tilde{f}\to\infty$, $P(\tilde{f})$ behaves as

$$\lim_{\tilde{f}\to\infty}\log P_{\lambda\to\infty}(\tilde{f}) \sim \tilde{f}-\frac{k_0\theta}{\tilde{\gamma}}\exp(\tilde{f})$$
$$\lim_{\tilde{f}\to\infty}\log P_{\lambda\to 0}(\tilde{f}) \sim -\tilde{f}/2. \quad \text{(S36)}$$

It is worth noting that depending on the λ value ($\lambda\to\infty$ or 0) $P_\lambda(\tilde{f})$ differs in its asymptotic behavior with respect to $\tilde{f}$ (see Eqs.S35 and S36).

The asymptotic behavior of the so-called microscopic model [5, 6], whose force range is limited by the critical force ($f < f_c = \Delta G^\ddagger/\nu\Delta x^\ddagger$), is reduced to that of Gumbel distribution only if $f^* < f \ll f_c$. If $f^* < f\to f_c$ then the unbinding force distribution decays precipitously to zero as $\sim(1-f/f_c)^{1/\nu-1}$($\nu = 2/3$: cubic potential) and linearly ($\nu = 1/2$: harmonic cusp potential) ($\lambda\to\infty$ corresponds to the Bell model). Note that the model in [5, 6] corresponds to $\lambda\to\infty$.

In contrast, for $\tilde{f}\to 0$,

$$\lim_{\tilde{f}\to 0}P_{\lambda\to\infty}(\tilde{f}) \sim \frac{k_0\theta}{\tilde{\gamma}}\times\exp\left[\left(1-\frac{k_0\theta}{\tilde{\gamma}}\right)\tilde{f}\right]$$
$$\lim_{\tilde{f}\to 0}P_{\lambda\to 0}(\tilde{f}) \sim \frac{k_0\theta}{\tilde{\gamma}}\times\left[1+\left(1-3\frac{k_0\theta}{\tilde{\gamma}}\right)\tilde{f}+\mathcal{O}(\tilde{f}^2)\right]. \quad \text{(S37)}$$

The initial slope of $P(\tilde{f})$ is determined by the value of $k_0\theta/\tilde{\gamma}$.

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