

## Achieving the Threshold Regime with an Overscreened Josephson Junction

Eugene V. Sukhorukov<sup>1</sup> and Andrew N. Jordan<sup>2</sup>

<sup>1</sup>*Département de Physique Théorique, Université de Genève, CH-1211 Genève 4, Switzerland*

<sup>2</sup>*Department of Physics and Astronomy, University of Rochester, Rochester, New York 14627, USA*

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We demonstrate that by utilizing an overscreened Josephson junction as a noise detector it is possible to achieve the threshold regime, whereby the tails of the fluctuating current distribution are measured. This situation is realized by placing the Josephson junction and mesoscopic conductor in an external circuit with very low impedance. In the underdamped limit, overscreening the junction inhibits the energy diffusion in the junction, effectively creating a tunable activation barrier to the dissipative state. As a result, the activation rate is qualitatively different from the Arrhenius form.

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In the course of scientific progress, it is desirable to use recent advances of physical understanding to develop new experimental techniques that will in turn give impetus for further advances. This description is apt for the current state of electron counting statistics in mesoscopic physics, and the aim of the current Letter. Recent experiments in electron counting statistics have measured the asymmetry of the current distribution [1,2] and the effect of the measurement circuit [1,3]. Subsequent researchers have taken several independent approaches to exploring this physics. If the system is transferring electrons on sufficiently long time scales, it is possible to directly count individual electrons [4], look at various higher current cumulants [5], and even examine the conceptually new area of conditional counting statistics [6]. It is also possible to focus on the frequency dependence of the cumulants [7,8]. A further exciting possibility is to measure the tails of the current fluctuation distribution via a threshold detector [9,10].

A natural threshold detector is a Josephson junction (JJ): by measuring the rate of switching out of the metastable supercurrent state into the running dissipative state, information about the statistical properties of the noise driving the system may be extracted. This system has been the subject of recent experiments [11–13]. However, it was noticed that a single JJ typically works near the Gaussian point [9] (a regime that has been well studied in the past; see, e.g., [14,15]), so much so that the third current cumulant makes only a small correction to the escape rate [16–18]. This is because the slow semiclassical dynamics of the JJ averages out the Markovian noise source and becomes effectively Gaussian under typical conditions. The purpose of this Letter is to show that despite this difficulty, the *threshold regime* is realizable for a single underdamped JJ. The threshold regime goes beyond the third cumulant and realizes the full potential of the Josephson detector, where the escape to the running state of the JJ is driven by the tails of the distribution, rather than by relatively small deviations from the average.

Circuit effects are known to be important for electron counting statistics. The measurement circuit can lead to

cascade corrections [19,20] to higher current cumulants, masking the system contribution. A similar effect has been predicted by the authors to occur in JJ detectors [17]. The central idea of this Letter is to have a very small load resistance, resulting in an overscreened junction, while still maintaining the underdamped state (see inset of Fig. 1). The advantage is that not only cascade corrections are suppressed, but also the relative contribution of the third cumulant is enhanced compared to the Gaussian contribution [17], assisting in the measurement of the third cumulant in recent experiments [12,13]. We will demonstrate that this is the case not just for the third cumulant, but that all higher cumulants are also enhanced, leading to non-perturbative activation by rare events.

*Josephson detection circuit.*—The circuit in the inset of Fig. 1 shows the essential part of the detector composed of

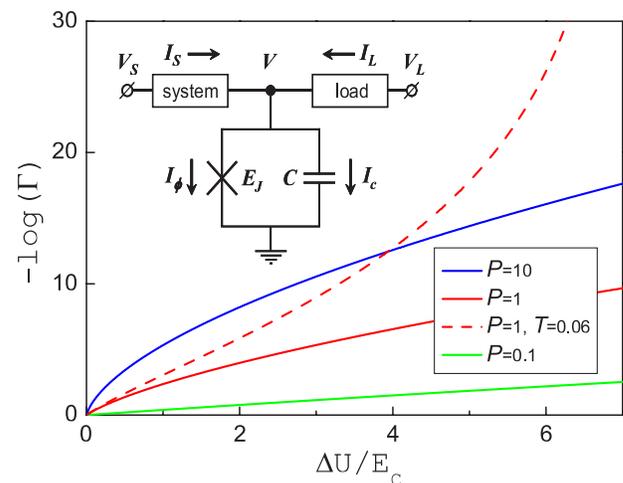


FIG. 1 (color online). The negative logarithm of the escape rate is plotted versus dimensionless barrier height  $\Delta U/E_c$  for various values of the threshold parameter  $\mathcal{P}$ . The solid lines correspond to Poissonian noise, while the dashed line corresponds to a QPC with transparency  $\mathcal{T} = 0.06$ . Inset: Simplified electrical circuit for the Josephson junction (JJ, marked with an X) threshold detector of the noisy mesoscopic system.

the JJ with Josephson energy  $E_J$ , and the capacitor  $C$ . The fluctuations in the circuit originate from the combination of the current  $I_L$  through the macroscopic load resistor and by the mesoscopic system current  $I_S$ , which is to be measured. According to Kirchhoff's law, the total current  $I_S + I_L$  is equal to the sum of the Josephson current  $I_\phi = (E_J/\Phi_0) \times \sin\phi$ , where  $\Phi_0 = \hbar/2e$ , and the displacement current  $I_C = C\dot{V}$ . This leads to the equation of motion for the superconducting phase  $\phi$ ,

$$C\Phi_0^2\ddot{\phi} + E_J \sin\phi = \Phi_0(I_S + I_L), \quad (1)$$

where we used the relation  $V = \Phi_0\dot{\phi}$ .

We consider an Ohmic system and load resistor, so that  $\langle I_S \rangle = J_S - G_S V$  and  $\langle I_L \rangle = J_L - G_L V$ , where  $G_S, G_L$  are the system and load conductances, and the constant currents  $J_S = G_S V_S, J_L = G_L V_L$  are tunable parameters that will be shown to control the activation threshold. Equation (1) can be rewritten as a set of Hamiltonian-Langevin equations for the phase variable  $\phi$  and canonically conjugated momentum  $p = \Phi_0 Q$  (where  $Q = CV$  is the total charge on the capacitor),

$$\dot{\phi} = p/m, \quad \dot{p} = -\partial U/\partial\phi + \Phi_0\delta I, \quad (2)$$

with “mass”  $m = \Phi_0^2 C$ , and where  $\delta I = I_S - J_S + I_L - J_L$  is the dissipative part of the system and load current. Equation (2) describes the motion of a “particle” in the tilted periodic potential

$$U(\phi) = -E_J[\cos\phi + \mathcal{J}\phi], \quad (3)$$

stimulated by the dissipative part of the system and load current, where  $\mathcal{J} = \Phi_0(J_S + J_L)/E_J$  is the dimensionless total current bias. Dissipation leads to relaxation of the JJ into one of its metastable supercurrent states, where the phase is localized in one of the potential wells so that  $\langle V \rangle = 0$ . In the dissipative state, the phase drifts along the bias which generates a nonzero voltage drop  $V$ .

*Weak damping threshold limit.*—Here the phase oscillates with the plasma frequency  $\omega_{\text{pl}} = \Omega_J(1 - \mathcal{J}^2)^{1/4}$ , where  $\Omega_J = \sqrt{E_J/m}$ . The energy relaxes slowly with rate  $(G_S + G_L)/C < \omega_{\text{pl}}$  to the local potential minimum. We further assume the separation of time scales,  $\max\{eV_S, T\} > \hbar\omega_{\text{pl}}$ , so that the noise source  $I_S$  is Markovian. According to our previous results [17], the escape rate  $\Gamma$  (predicted to exponential accuracy) is

$$\log\Gamma = -\Phi_0^{-1} \int \lambda dE, \quad \langle \mathcal{H}(\lambda\dot{\phi}) \rangle_E = 0. \quad (4)$$

Here, the function  $\mathcal{H}(z)$ , which we refer to as a Hamiltonian, generates the cumulants (irreducible moments) of  $\delta I$  in Eq. (2), so that  $\partial_z^n \mathcal{H}|_{z=0} = \langle \delta I^n \rangle$ . The function  $\lambda(E)$  is the escape trajectory in the extended energy phase space, or the “instanton line.” The notation  $\langle \dots \rangle_E = (1/T_p) \oint dt(\dots)$  denotes time averaging over a physical trajectory in the  $(q, p)$  phase space at a certain energy  $E$ , with  $T_p(E)$  being the period of the quasiperiodic

motion at that energy. See Ref. [17] for a detailed derivation.

We note that after replacing the node voltage  $V$  with  $\Phi_0\dot{\phi}$ , the first two coefficients in the expansion of  $\mathcal{H}$  are (i)  $\langle \delta I \rangle = -(G_S + G_L)\Phi_0\dot{\phi}$  and (ii)  $\langle \delta I^2 \rangle = 2G_L T + \langle I_S^2 \rangle$ . We now focus on the limit of small load resistance,  $G_L \gg G_S$ , so circuit backaction can be neglected [17]. In this case, the load controls coefficient (i). However, we observe that for a large system voltage, the system shot noise will dominate the load resistor noise in coefficient (ii) if  $eV_S > G_L T/G_S$ . The circuit load, being a macroscopic resistor, has no higher current cumulants, so the Hamiltonian takes the form  $\langle \mathcal{H}(\lambda\dot{\phi}) \rangle_E = -G_L \Phi_0 \lambda \langle \dot{\phi}^2 \rangle_E + J_S \langle F(\lambda\dot{\phi}) \rangle_E$ , where  $F$  generates the generalized Fano factors of the system,

$$F(z) = \sum_{n=2}^{\infty} \frac{z^n \langle I_S^n \rangle}{n! \langle I_S \rangle}. \quad (5)$$

The instanton equation (4) may be brought to dimensionless form

$$\frac{E_C \lambda \langle F(\lambda\dot{\phi}) \rangle_E}{\hbar \langle (\lambda\dot{\phi})^2 \rangle_E} = \mathcal{P} \equiv \frac{eG_L}{4CJ_S}, \quad (6)$$

where  $E_C = e^2/2C$  is the capacitor's charging energy, and we define the threshold parameter  $\mathcal{P}$ , which is the ratio of the JJ energy relaxation rate  $G_L/C$  to the mesoscopic system's excitation rate  $J_S/e$ . If relaxation is weak,  $\mathcal{P} \ll 1$ , then Eqs. (5) and (6) may be truncated at the first (Gaussian) term, giving the escape rate (4) as

$$\log\Gamma_G = -\frac{\Delta U}{T_{\text{eff}}}, \quad (7)$$

where  $T_{\text{eff}} = \langle I_S^2 \rangle / 2G_L$  is an effective noise temperature, so the Arrhenius form is recovered in agreement with [17]. However, if the relaxation rate dominates the excitation rate so  $\mathcal{P} \gg 1$ , then we enter the *threshold regime* where the solution to (6) is nonperturbative in the cumulant expansion.

The physical reason for this is that while the Gaussian contribution to the system noise cannot compensate for the fast energy relaxation, rare current events (where many electrons are sequentially transmitted in a short amount of time) may be able to excite the JJ strongly enough to overcome damping. The power of these current kicks is proportional to the velocity  $\dot{\phi}$ , and therefore creates positive feedback for rare-event activation in the weak damping regime, because the velocity continues to grow as the particle gradually ascends the potential well to the escape point.

By introducing the JJ quality factor  $\mathcal{Q} = \omega_{\text{pl}} C/G_L > 1$  and the separation of time scales parameter  $\mathcal{R} = eV_S/\hbar\omega_{\text{pl}} > 1$  we reformulate the condition for the threshold regime as  $\mathcal{P} = (1/\mathcal{Q}\mathcal{R})(e^2/4\hbar G_S) > 1$ ; therefore,  $G_S < e^2/\hbar$ . This implies that the system is a tunnel junction, which is known to create Poissonian noise (however,

see the discussion below). The circuit backaction can be neglected if  $eV_S > eV = e\Phi_0\dot{\phi} \sim \sqrt{E_C E}$ . We will show below that  $E \sim E_C$  at the threshold, so  $eV_S/E_C = Q\mathcal{R}(\hbar G_L/e^2) > 1$ . Together with  $\mathcal{P} > 1$  this gives the overscreening condition  $G_L > G_S$ .

*Poissonian noise.*—In order to illustrate the threshold behavior, we consider a simple harmonic potential  $U(\phi) = (1/2)m\omega_{\text{pl}}^2\phi^2$  with a sharp cutoff at  $\phi = \phi_0$ . In this simplified model  $\phi(t) = \sqrt{2E/E_J} \cos(\omega_{\text{pl}}t)$ . The averages in Eq. (6) over the periodic orbit at constant energy may now be done exactly. The Poissonian noise generator is  $F(z) = e^z - z - 1$ . Using the integral representation of the modified Bessel function of order 0,  $I_0(z) = (1/2\pi) \times \int_0^{2\pi} d\theta \exp(-z \sin\theta)$ , we find Eq. (6) simplifies to

$$\frac{I_0(z) - 1}{z} = 2\mathcal{P} \sqrt{\frac{E}{E_C}}, \quad z \equiv \lambda \sqrt{\frac{2E}{m}}. \quad (8)$$

This equation is numerically solved for  $\lambda$  and integrated over the energy  $E$ . The result for the logarithm of the escape rate is plotted in Fig. 1 as a function of  $\Delta U$  for different values of  $\mathcal{P}$ .

For  $\mathcal{P} \ll 1$ , the Bessel function may be expanded as  $I_0(z) = 1 + z^2/4 + \dots$ , which gives  $\lambda = 2\hbar\mathcal{P}/E_C$ . When  $\lambda$  is substituted into (4), the result (7) is recovered, because for Poissonian noise  $\langle I_S^2 \rangle = G_S V_S$ . In the opposite limit,  $\mathcal{P} \gg 1$ , the right-hand side of Eq. (8) is large, and we utilize the asymptotic form of the Bessel function,  $I_0(z) \rightarrow e^z/\sqrt{2\pi z}$  for  $z$  large. Thus the variable  $z$  must be only a logarithmically growing solution of Eq. (8):  $z = L_H(E) \equiv \log(2\mathcal{P}\sqrt{2\pi E/E_C})$ , up to the double-logarithmic correction. This gives  $\lambda = L_H(E)\sqrt{m/2E}$ , so the escape rate (4) of activation due to Poissonian noise may be approximated by

$$\log\Gamma_P = -\sqrt{\frac{\Delta U}{E_C}} \log\left(2\mathcal{P}\sqrt{\frac{2\pi\Delta U}{E_C}}\right). \quad (9)$$

The dominant square-root behavior is clearly seen in Fig. 1, and is a sign of the break down of Gaussian noise activation.

For the escape to be achievable, we require that  $\Delta U$  is not much larger than  $E_C$ . More rigorously, for strong bias the action scales as  $\sqrt{\Delta U/E_C} \sim (1 - \mathcal{J})^{3/4} \sqrt{E_J/E_C}$ . The total number of states in the quantum well,  $N_{\text{tot}} \sim \Delta U/\hbar\omega_{\text{pl}}$ , also scales down as  $N_{\text{tot}} \sim (1 - \mathcal{J})^{5/4} \times \sqrt{E_J/E_C}$ , but must be large in the semiclassical limit. Therefore, we estimate  $-\log\Gamma_P \sim N_{\text{tot}} \log\mathcal{P}/\sqrt{1 - \mathcal{J}_{\text{st}}}$ . In Refs. [12,13] this number varies from 10 to 17. This indicates that it may be hard but nevertheless feasible to observe the threshold behavior.

*Anharmonic correction.*—Here we show that a realistic potential  $U(\phi)$  leads to only a small anharmonic correction to the escape rate plotted in Fig. 1. In the case of a general potential, the instanton line Eq. (6) may be found by noting

that for Poissonian noise the integral in  $\langle F(\lambda\dot{\phi}) \rangle_E$  in the threshold regime is exponentially dominated by the largest value of  $\dot{\phi}$ . Energy conservation,  $(m/2)\dot{\phi}^2 + U(\phi) = E$ , indicates that this will be near the bottom of the potential, where the potential is approximately harmonic with frequency  $\omega_{\text{pl}}$ . Therefore, the previous Bessel function result will hold to excellent approximation. The other average,  $\langle \dot{\phi}^2 \rangle_E$ , is generalized by noting that by using  $dt = d\phi/|\dot{\phi}|$  and conservation of energy, the equation  $\partial_E[T_p(E) \times \langle \dot{\phi}^2 \rangle_E] = T_p(E)/m$  holds. This equation may be integrated to find  $\langle \dot{\phi}^2 \rangle_E = 2\pi\hbar N(E)/mT_p(E)$ , where  $N(E) = (1/2\pi\hbar) \int_0^E dE' T_p(E')$  is the number of quantum states in the cavity below the energy  $E$ .

Combining results for  $\langle F \rangle_E$  and  $\langle \dot{\phi}^2 \rangle_E$  and going to the asymptotic threshold regime, we conclude that the instanton solution takes the previously found form  $\lambda = L(E) \times \sqrt{m/2E}$  with the logarithm replaced by  $L(E) = L_H(E) - \log[A(E)]$ , where the function

$$A(E) = \frac{T_p E}{2\pi\hbar N} = E \partial_E \log N \quad (10)$$

characterizes the anharmonicity. It takes the value  $A = 1$  at the bottom of the potential well and diverges logarithmically at the top of the barrier as  $A \sim \log[\Delta U/(\Delta U - E)]$ . Thus this factor may compensate the large parameter  $\mathcal{P}$  in close vicinity of the barrier top,  $\Delta U - E \sim \Delta U e^{-\mathcal{P}}$ , so the Gaussian noise there dominates. However, the overall anharmonic contribution to the large logarithm  $L(E)$  is relatively small and can be neglected.

*Stabilization effects.*—In what follows we wish to illustrate an important fact: statistics of rare current fluctuations carries complementary information about a Markovian process which is not contained in any finite current cumulant. As a first example, we consider the shot noise from a quantum point contact (QPC), which is known to be a binomial process. The noise Fano factors are generated by  $F(z) = (1/\mathcal{T}) \log[1 + \mathcal{T}(e^z - 1)] - z$ , where  $\mathcal{T}$  is the transmission of the QPC. The cumulants are obtained by the expansion (5) at  $z = 0$ . Therefore, in the tunneling regime,  $\mathcal{T} \ll 1$ , the logarithm can be expanded to lowest order in  $\mathcal{T}$  so the noise of the QPC is Poissonian,  $F(z) = e^z - z - 1$ , as in the example considered above. However, our results indicate that the Poissonian process can always provide a strong enough current fluctuation for the JJ to escape from the metastable state with some small probability. This is not the case for a QPC in the tunneling regime, as we demonstrate below.

Indeed, the rare current events of the QPC are determined by the asymptotic of  $F(z)$  at large  $z$ :

$$F(z) + z = \frac{1}{\mathcal{T}} \times \begin{cases} z + \log\mathcal{T}, & z \rightarrow +\infty, \\ \log(1 - \mathcal{T}), & z \rightarrow -\infty. \end{cases} \quad (11)$$

This result implies that with a small probability  $\mathcal{T}^M$  the current acquires its maximum value  $J_{\text{max}} = J_S/\mathcal{T} = e^2 V_S/\pi\hbar$ , so that all  $M = eV_S \Delta t/\pi\hbar$  electrons arriving

at the QPC during time interval  $\Delta t$  are transmitted. Similarly, with the probability  $(1 - \mathcal{T})^M$  all the electrons are reflected giving zero current [21]. If  $\Delta U$  exceeds some critical value, the maximum or minimum current fluctuation creates insufficient bias for the JJ to escape from the supercurrent state [9]. We named this phenomenon the *Pauli stabilization effect* because it originates from the Pauli principle of electron occupation [10].

To estimate the value of the bias,  $\mathcal{J}_{\text{st}}$ , and escape rate,  $\Gamma_{\text{st}}$ , at the stabilization point, we average  $F$  in (11) over the energy-conserving trajectories, which yields  $\langle F \rangle_E = (1/\mathcal{T})\{\lambda\Delta\phi/T_p + (1/2)\log[\mathcal{T}(1 - \mathcal{T})]\}$ , where  $\Delta\phi$  is the distance between the two turning points at energy  $E$ . The denominator of (6) is the same as before, so the instanton line is given by

$$\lambda_{\text{st}}(E) = \frac{aT_p}{\Delta\phi - bN}, \quad (12)$$

with the coefficients  $a = -(1/2)\log[\mathcal{T}(1 - \mathcal{T})]$  and  $b = 16\pi\mathcal{P}\mathcal{T}$ . Integration of this instanton line leads to singular behavior when  $\Delta\phi = 16\pi\mathcal{P}\mathcal{T}N$ . Evaluated at  $E = \Delta U$  this equality determines the stabilization point  $\mathcal{J}_{\text{st}}$  below which the rate  $\Gamma$  vanishes. We find this point by assuming the strong bias limit  $1 - \mathcal{J} \ll 1$  in (3) giving a cubic potential:  $U/E_J = (1 - \mathcal{J})\phi - \phi^3/6$ . Skipping a number of steps, we present the result:

$$(1 - \mathcal{J}_{\text{st}})^{3/4} = \frac{5\sqrt{E_C/E_J}}{32 \times 2^{1/4}\mathcal{P}\mathcal{T}}. \quad (13)$$

Next, we note that at the critical point  $\mathcal{J} = \mathcal{J}_{\text{st}}$  the integral (4) with  $\lambda = \lambda_{\text{st}}(E)$  from Eq. (12) is convergent because  $\lambda_{\text{st}}(E)$  has only an inverse square-root divergence at  $E = \Delta U$ . Therefore,  $\log\Gamma_{\text{st}}$  can be estimated by dropping the least singular term,  $bN$ , in the denominator of  $\lambda_{\text{st}}(E)$  and replacing the potential with the harmonic approximation  $U = (1/2)m\omega_{\text{pl}}^2\phi^2$ . Straightforward evaluation then gives  $\log\Gamma_{\text{st}} \sim \sqrt{\Delta U/E_C}\log[\mathcal{T}(1 - \mathcal{T})]$ . At strong bias  $\Delta U \sim (1 - \mathcal{J}_{\text{st}})^{3/2}E_J$ , and using the result (13) we find the escape rate right before stabilization:

$$\log\Gamma_{\text{st}} \sim \frac{\log[\mathcal{T}(1 - \mathcal{T})]}{\mathcal{P}\mathcal{T}}. \quad (14)$$

We estimate (14) in terms of the quality factor  $\mathcal{Q} > 1$  and the separation of time scales parameter  $\mathcal{R} > 1$  as  $\log\Gamma_{\text{st}} \sim \mathcal{Q}\mathcal{R}\log[\mathcal{T}(1 - \mathcal{T})]$ . Alternatively, using the total number of states  $N_{\text{tot}} > 1$  in the quantum well, and Eq. (13), we find that  $\log\Gamma_{\text{st}} \sim N_{\text{tot}}\log[\mathcal{T}(1 - \mathcal{T})]/\sqrt{1 - \mathcal{J}_{\text{st}}}$ , which agrees with the previous estimate for Poissonian noise. To compare with the Poissonian case, the escape rate for a QPC with the transmission  $\mathcal{T} = 0.06$  is plotted in Fig. 1 in the case of the simple harmonic potential with a sharp cutoff. The sharp potential leads to a logarithmic divergence at the stabilization point  $\Delta U/E_C = 6.86$ , in contrast with the rate discontinuity discussed above for a smooth potential.

Our second example is the telegraph process: The system current switches randomly from the value  $I_1$  to the value  $I_2$  and back with the rate  $\gamma_1$  and  $\gamma_2$ , respectively. We have found in Ref. [22] that the cumulant generator of the telegraph process is given by the formula

$$\mathcal{H}_S(z) = \bar{I}z - \bar{\gamma} + \sqrt{(\Delta I z - \Delta\gamma)^2/4 + \gamma_1\gamma_2}, \quad (15)$$

where  $\Delta I = I_2 - I_1$ ,  $\Delta\gamma = \gamma_2 - \gamma_1$ ,  $\bar{I} = (I_1 + I_2)/2$ , and  $\bar{\gamma} = (\gamma_1 + \gamma_2)/2$ . In the case of slow switching,  $\gamma_1, \gamma_2 \ll I_1, I_2$ , the noise becomes super-Poissonian with the current distributed between  $I_1$  and  $I_2$ . The sharp cutoff of the distribution function at these values results in the stabilization effect. The asymptotic form  $\mathcal{H}_S = \bar{I}z - \bar{\gamma} + |\Delta I z - \Delta\gamma|/2$  at  $|z| \rightarrow \infty$  leads to the result (12) with the new coefficients  $a = \bar{\gamma}/\Delta I$  and  $b = 8\pi G_L E_C/e\Delta I$ . Therefore, at the stabilization point the results (13) and (14) hold after the replacement  $\mathcal{P}\mathcal{T} \rightarrow G_L E_C/2e\Delta I$  and  $\log[\mathcal{T}(1 - \mathcal{T})] \rightarrow -2\bar{\gamma}/\Delta I$ . With the new separation of time scales requirement,  $\mathcal{R}' = \bar{\gamma}/\omega_{\text{pl}} > 1$ , the action at the stabilization point can be estimated as  $-\log\Gamma_{\text{st}} \sim \mathcal{Q}\mathcal{R}'$ , so the telegraph stabilization effect should also be observable.

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