Detecting Charge Noise with a Josephson Junction: A Problem of Thermal Escape in Presence of Non-Gaussian Fluctuations

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Motivated by several experimental activities to detect charge noise produced by a mesoscopic conductor with a Josephson junction as on-chip detector, the switching rate out of its zero-voltage state is studied. This process is related to the problem of thermal escape in presence of non-Gaussian fluctuations. In the relevant case of weak higher than second order cumulants, an effective Fokker-Planck equation is derived, which is then used to obtain an explicit expression for the escape rate. Specific results for the rate asymmetry due to the third moment of current noise allow to analyze experimental data and to optimize detection circuits.

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A complete understanding of electronic transport through mesoscopic conductors necessitates the knowledge of all noise properties of the corresponding current. This is the goal of full counting statistics, which has attracted considerable activity in recent years [1,2]. Theoretically, generating functionals of current noise in the low frequency limit have been calculated for a variety of systems comprising tunnel contacts, diffusive wires, and ballistic cavities [2]. Experimentally, in a pioneering measurement the third moment of current noise produced in a tunnel junction was detected by analog amplifiers and filtering techniques in [3] and later also in [4]. Since then strong efforts have been made towards on-chip detection schemes, first because they are faster and second because they give access to finite frequency noise properties. Lately, the distribution of charges flowing through a quantum dot have been extracted on-chip in the low frequency regime [5]. The goal now is to push devices into higher frequency ranges (GHz), where quantum effects, electron-electron interactions, and plasmon dynamics are relevant.

Based on an idea proposed in [6] and studied later in various scenarios [7,8], currently, several experiments are aiming to set up circuits with Josephson junctions (JJs) as detection elements [9–11]. JJs can be fabricated and manipulated in a very controlled way, offer a large bandwidth depending on their plasma frequency, and contain an intrinsic amplification mechanism. Namely, the switching out of the zero-voltage state is exponentially sensitive to variations of the barrier potential and to the noise strength. For electrical noise large positive and negative fluctuations from the mean occur with different probabilities. The idea is thus to probe this asymmetry by measuring switching rates for mean mesoscopic currents flowing forward and backward, respectively. Indeed, recent experimental results [11] indicate a sufficient sensitivity of such a circuit to retrieve the third cumulant.

A typical setup [10] consists of a JJ, on which two currents I_b and I_m are injected. Current I_b is a standard bias current coming from a source in parallel to the JJ, while I_m runs through a noise generating mesoscopic conductor in series with the JJ in such a way that no dc component of I_m passes through the detector. Because of substantial heating from the additional electrical noise the JJ operates in the regime of classical escape, where quantum effects are negligible. Then, according to the resistively and capacitively shunted junction model, the phase φ of the JJ moves in a tilted washboard potential $-E_J\cos(\varphi) - (\hbar/2e)\langle I_b\rangle\varphi$ with Josephson energy E_J and is subject to Johnson-Nyquist noise $\delta I_b = I_b - \langle I_b\rangle$ and stationary non-Gaussian current fluctuations $\delta I_m = I_m - \langle I_m\rangle$. If the phase is initially trapped in one of the wells (zero-voltage state), it may for sufficiently large $\langle I_b\rangle < I_c = (2e/\hbar)E_J$ escape so that the JJ switches to a finite voltage state. Further, since the third cumulant vanishes in equilibrium due to time-reversal symmetry, experimentally, the mesoscopic conductor is at low temperatures driven far from equilibrium into the shot noise regime, where no fluctuation-dissipation theorem applies. Hence, the switching of the JJ can be visualized as the diffusive dynamics of a fictitious particle in a metastable well with non-Gaussian continuous fluctuations acting as external random driving force.

In the past JJs have been used to thoroughly confirm various theories for thermal escape [12,13] including cases of external driving either by time-periodic forces or Gaussian noise sources related to phenomena such as resonant activation [14,15] and stochastic resonance [16]. When putting the experimental situation described above into this context, one arrives at a new type of rate problem: thermal escape driven by non-Gaussian continuous noise. Theoretically, the challenge here is that externally driven escape necessitates a full dynamical description because the stationary well state from which particles are ejected over the barrier is only known *a posteriori*. In the classical domain and for Gaussian noise processes a dynamical formulation is based on a Fokker-Planck equation (FPE) for the phase-space distribution as shown, e.g., in [13,17]. Thus, the first goal is to derive a generalized FPE for weak non-Gaussian noise, which then serves as a basis for the rate calculation. A solution to this problem is of general interest for rate processes in complex media and crucial for present electrical noise measurements because non-Gaussian components are typically very small compared with a prevailing Gaussian background. In this Letter we develop the framework for such a rate theory and give explicit expressions in case of a JJ as detector.

For this purpose let us consider the Langevin equation for the diffusive motion of a particle of mass m

$$m\ddot{\varphi}(t) + V'(\varphi) - \eta(t) + m\gamma\dot{\varphi}(t) = \xi(t), \qquad (1)$$

where φ denotes a generalized coordinate, the overdot represents d/dt, and the prime represents $d/d\varphi$. The barrier potential $V(\varphi)$ is assumed to be sufficiently smooth with a well located at $\varphi = 0$ with frequency $\omega_0 = \sqrt{V''(0)/m}$ and a barrier top at $\varphi = \varphi_b$ with frequency $\omega_b = \sqrt{|V''(\varphi_b)|/m}$ and height $V_b = V(\varphi_b) - V(0)$. The white thermal Gaussian noise obeys $\langle\xi(t)\rangle_\beta = 0$ and is related to the damping γ via the dissipation fluctuation theorem: $\langle\xi(t)\xi(t')\rangle_\beta = (2m\gamma/\beta)\delta(t-t')$ with inverse temperature $\beta = 1/k_BT$. The statistical properties of the stationary non-Gaussian noise are determined by the generating functional

$$e^{-S[w(s)]} = \left\langle \mathcal{T}\exp\left[i\int_0^t ds\,\eta(s)w(s)\right]\right\rangle, \qquad (2)$$

where $\mathcal{T}$ is the time ordering operator and the cumulants are gained from the functional derivatives of $S[w]$. In particular, we assume $S_1(t) = \langle\eta(t)\rangle = i\partial S[w]/\partial w(t)|_{w=0} = 0$, while the autocorrelation function reads

$$S_2(t) = \langle\eta(t)\eta(0)\rangle = \partial^2 S[w]/\partial w(t)\partial w(0)|_{w=0}$$

and the third cumulant follows accordingly as

$$S_3(t,t') = \langle\eta(t+t')\eta(t')\eta(0)\rangle.$$

Now, particles initially confined in the well region may escape such that for sufficiently high barriers ($\beta V_b \gg 1$) a stationary flux appears related to an escape rate $\Gamma = \langle v\delta(\varphi - \varphi_b)\rangle_{\rm flux}/N_{\rm well}$. The denominator is the population in the well and the expectation value of $v = \dot{\varphi}$ is taken with respect to a quasistationary nonequilibrium state, the flux state, which may be cast into the form

$$P_{\rm flux}(v,\varphi) = P_{\rm equi}(v,\varphi)f_{\rm flux}(v,\varphi). \qquad (3)$$

Here $P_{\rm equi}$ denotes the equilibrium distribution and $f_{\rm flux}$ describes deviations from it such that it tends to 0 towards the continuum and to 1 towards the well. The flux state, as a phase-space density, is obtained as a stationary solution to the Fokker-Planck equation (FPE) corresponding to (1). In case of vanishing non-Gaussian noise, $\eta(t) = 0$, the latter one reads $\partial_t P(v,\varphi,t) = L_0 P(v,\varphi,t)$ where

$$L_0 = -v\partial_\varphi + \partial_v[V'(\varphi)/m + \gamma v] + \gamma/(m\beta)\partial_v^2. \qquad (4)$$

Then, from $L_0 P_{\rm flux} = 0$ one finds for moderate to strong friction the known expression [13] $\Gamma = (\omega_0\Omega/2\pi)\times\exp(-\beta V_b)$ with the scaled Grote-Hynes frequency $\Omega = -\gamma/(2\omega_b) + \sqrt{(\gamma/2\omega_b)^2 + 1}$.

For finite non-Gaussian noise the translation of (1) into an equivalent FPE for the averaged phase-space distribution leads to only formal expressions [17,18]. Here, in accordance with the experimental situation we proceed by assuming that (i) non-Gaussian fluctuations are weak and (ii) sufficiently fast compared to the bare dynamics of the system (detailed conditions will be given below). Then, one considers $V_{\rm eff} = V - \varphi\eta(t)$ as an effective time-dependent potential so that for each realization of the non-Gaussian random force a FPE for $P_\eta(v,\varphi,t)$ exists where L_0 is replaced by $L_0 - \eta(t)/m\partial_v$. To gain a FPE for the averaged distribution $\langle P_\eta(v,\varphi,t)\rangle$, we switch to the interaction picture $P_\eta(t) = \exp[L_0 t]Q_\eta(t)$ with $\partial_v(t) = \exp(-L_0 t)\partial_v\exp(L_0 t)$. By averaging the equation for $Q_\eta(t)$ over the η noise, the exact solution is expressed in terms of the generating functional with the counting field $w(s)$ substituted by $(i/m)\partial_v(s)$, i.e.,

$$\begin{aligned}\langle Q_\eta(t)\rangle &= e^{-S[(i/m)\partial_v(s)]}Q(0)\\ &= \mathcal{T}\exp\left[\int_0^t ds\sum_{k\geq 2}\frac{C_k(s)}{m^k}\right]Q(0). \qquad (5)\end{aligned}$$

Here, we have used the cumulant expansion of $S[(i/m)\partial_v(s)]$ with $C_k(s) = \exp[-L_0 s]\hat{C}_k(s)\exp[L_0 s]$, where the two lowest order cumulant operators read

$$\hat{C}_2(s) = \int_0^s du S_2(u)\partial_v\partial_v(-u)$$

$$\hat{C}_3(s) = -\int_0^s du\int_0^{s-u} du' S_3(u,u')\partial_v\partial_v(-u)\partial_v(-u-u').$$

The above result reveals that the generator for the averaged dynamics is directly given by the generating functional of the non-Gaussian noise. Its cumulant expansion now allows for a systematic approximation.

To do so, we take only the two leading order terms into account in (5), which in turn leads to $\partial_t\langle P_\eta\rangle = [L_0 + \hat{C}_2(t)/m^2 + \hat{C}_3(t)/m^3]\langle P_\eta\rangle$. While this FPE applies to weak non-Gaussian noise of an arbitrary type, following assumption (ii), the noise correlation functions vanish on a sufficiently short time scale so that the non-Markovian dynamics reduces effectively to a Markovian one. In this way, one arrives at a generalized FPE of the form $\partial_t\langle P\rangle = L_{\rm eff}\langle P\rangle$ with

$$L_{\rm eff} = L_0 + (c_2/m^2)\partial_v^2 - (c_3/m^3)\partial_v^3 \qquad (6)$$

where

$$c_2 = \int_0^\infty du S_2(u), \qquad c_3 = \int_0^\infty du\int_0^\infty du' S_3(u,u').$$

Note that this expression generalizes previous results obtained in the overdamped limit [18]. If the external noise were purely Gaussian, i.e., $c_3 = 0$, the operator $L_{\rm eff}$ were exact. Hence, what we really have to assume are small higher than second order cumulants, while c_2 may be large. Thus, one introduces an effective temperature

$$T_{\rm eff} = T + c_2/(k_B m\gamma) \tag{7}$$

and incorporates the Gaussian components of the non-Gaussian noise into a renormalized diffusion term according to $L_0(\beta) + (c_2/m^2)\partial_v^2 \to L_0(\beta_{\rm eff})$. We note in passing that this expression coincides with the one derived for resonant activation [14]. Experimentally, the heating due to c_2 is substantial so that (7) is required to capture the actual temperature of the JJ.

With the generalized FPE at hand we now attack the rate calculation. For this purpose, it is convenient to work with dimensionless quantities $\tau = \omega_0 t$, $x = \varphi\sqrt{\beta_{\rm eff} m\omega_0^2}$, $p = v\sqrt{\beta_{\rm eff} m}$, $U = \beta_{\rm eff} V$, and $\rho = \gamma/\omega_0$, $\bar{c}_3 = c_3(\beta_{\rm eff}/m)^{3/2}$; then x_b, $U_b \gg 1$ and $\bar{c}_3$ serves as a small parameter. We start with the equilibrium state (no flux) which determines the dominating exponential activation factor in (3) and write $P_{\rm equi} \propto P_{\beta_{\rm eff}} \exp(-\bar{c}_3 G)$ with the Boltzmann distribution $P_{\beta_{\rm eff}}$. Upon inserting this ansatz into $L_{\rm eff} P_{\rm equi} = 0$ we find perturbatively a distribution of the form

$$G(x, p) = \phi_0(x) + \sum_{n=1}^{3} \phi_n(x) p^n/n, \tag{8}$$

where coordinate dependent functions can be expressed in terms of only ϕ_2 as

$$\begin{aligned} \phi_1(x) &= -\rho\phi_2(x)/U'(x) = -\phi_3(x), \\ \phi_0(x) &= 3x - 3\rho\int_0^x dy\phi_1(y) + \int_0^x dy U'(y)\phi_2(y). \end{aligned} \tag{9}$$

Note that the property $\phi_1 = -\phi_3$ ensures that $\langle p\rangle_{\rm equi} = 0$ in leading order. While the above expressions hold in general, the solution of ϕ_2 depends on the specific form of the potential U. To better understand the deviations from the thermal equilibrium we thus first study a harmonic potential $U(x) = x^2/2$, which approximates a metastable potential around its well minimum. The result is $\phi_2(x) = -2x/(1 + 2\rho^2)$ and the corresponding distribution $P_{\rm equi}$ has a simple interpretation: it coincides with the thermal distribution driven by the external force $\eta(t)$ and averaged in the long time limit over the non-Gaussian noise. Specifically, one has $P_{\rm equi}(p, x) = \lim_{\tau\to\infty}\langle P_{\beta_{\rm eff}}(p_\tau, x_\tau)\rangle$, where $p_\tau = p - \dot{f}(\tau)$ and $x_\tau = x - f(\tau)$ with $f(\tau) = \sqrt{\beta_{\rm eff}/m\omega_0^2}\int_0^\infty d\tau' \chi(\tau - \tau')\eta(\tau')$ and $\chi(\tau)$ the response function for a damped harmonic oscillator. As a consequence, $P_{\rm equi}(p, x)$ shows an asymmetry even for a symmetric potential depending linearly on the third cumulant S_3. In contrast, in a setup where for the η noise a dissipation fluctuation theorem applies, a linear dependence on odd cumulants occurs only on a transient time scale [8].

A generic metastable potential is of the form $U(x) = (x^2/2)(1 - 2x/3x_b)$ describing particularly a well-barrier segment of a tilted washboard potential of a JJ. Then,

$$\phi_2(x) = -\frac{2x_b z(1-z)^a}{1+a}\,{}_2F_1(a, 1+a, 2+a, z), \tag{10}$$

where $z = x/x_b$, $a = 2\rho^2$, and $P_{\rm equi}$ follows together with (9). There is a little subtlety here in that the solution (10) applies for $2\rho^2 < 1$ only outside a narrow range around x_b. The global solution is then constructed by properly matching the local solution around x_b onto the latter one. Since the vicinity around the top affects only the prefactor and not the experimentally dominant exponential factor of the rate, we will give further details elsewhere. This exponential activation factor can now be inferred from $P_{\rm equi}(p = 0, x = x_b)$ and is obtained as $\Gamma \propto \exp[-U_b(1-g)]$ with the third cumulant correction $g = -\bar{c}_3\phi_0(x_b)/U_b$ explicitly evaluated as

$$g(\rho) \approx \frac{6\bar{c}_3}{5\rho^2 + 1}\frac{U_b}{x_b}. \tag{11}$$

The remaining prefactor to the rate is determined by $f_{\rm flux}$ in (3) and the well population $N_{\rm well}$. The latter one is easily calculated from $P_{\rm equi}(p, x)$ around the well bottom, while the former one is derived from the local dynamics around the barrier top. Eventually, the rate is found as

$$\Gamma = \frac{\omega_0\Omega}{2\pi}\left(1 - \sqrt{\frac{\pi}{2}}\bar{c}_3\phi_2'(x_b)\frac{\Omega^2\rho\kappa}{9 + 2\rho^2}\right)e^{-U_b(1-g)}, \tag{12}$$

where $\kappa = 2\Omega(9 + 4\rho^2) + \rho[8\rho^2 + \sqrt{\rho\Omega}(7 + 2\rho^2)] + 4$ and $\phi_2'(x_b)$ results for $2\rho^2 \le 1$ from the matching procedure leading, e.g., for $2\rho^2 \ll 1$ to $\phi_2'(x_b) = 1/\rho^2$. The above expression is the main finding of this Letter, namely, the thermal escape rate out of a metastable well in the presence of weak ($\bar{c}_3 \ll 1/U_b$) and fast (correlation time $\tau_c \ll 1/\omega_0$) non-Gaussian continuous noise. This result indeed verifies that the detector transforms the noise asymmetry into a rate asymmetry depending on the sign of the third cumulant $\bar{c}_3$. For $\bar{c}_3 > 0[<0]$ the noise distribution favors large positive (negative) fluctuations so that the barrier is effectively reduced (enhanced) and the distribution in the metastable potential develops a tail towards large positive (negative) momenta and coordinates causing a rate increase (decrease).

In the sequel, the general result (12) is applied to a circuit, where a JJ acts as detector for the current noise produced by a normal tunnel junction subject to a voltage V in the shot noise regime $eV \gg k_B T$. Further, $eV/\hbar \gg \omega_0$ guarantees that the noise is much faster than the plasma frequency $\omega_0 = (\sqrt{2E_J E_C}/\hbar)(1 - s_b^2)^{1/4}$ ($s_b = I_b/I_c$) of the JJ with charging energy $E_C = 2e^2/C$ (capacitance C) and that its back action onto the tunnel junction is negligible. Then, the current statistics is purely Poissonian: during a (scaled) time interval τ_p an average number of N charges passes the conductor so that $s_m \equiv \langle I_m/I_c\rangle = Ne\omega_0/(\tau_p I_c)$ and $S_3(\tau, \tau') = \delta(\tau + \tau')\delta(\tau)(e\omega_0/I_c)^2 s_m$. The procedure to extract the third cumulant is to measure switching rates with bias currents s_b (or equivalently mesoscopic currents s_m) flowing forward and backwards, re-

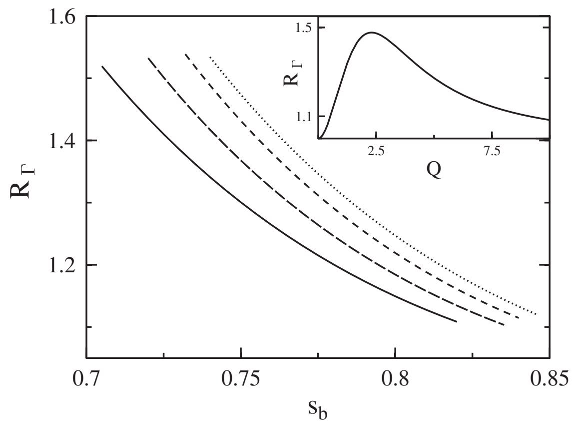

FIG. 1. Rate asymmetry R_Γ according to (13) for $\omega_0\hbar\beta = 0.5$, $\beta E_J = 200$, $Q = 2.5$, and various s_m vs the bias current s_b: $s_m = 3$ (solid), $s_m = 2.3$ (long-dashed), $s_m = 1.7$ (short-dashed), $s_m = 1.1$ (dotted). The inset shows R_Γ for the same parameters and $s_b = 0.75$, $s_m = 1.1$ as a function of Q.

spectively [11]. This way, one finds for the rate asymmetry $R_\Gamma = \Gamma(|s_b|)/\Gamma(-|s_b|)$ as the dominant contribution $R_\Gamma \approx \exp[2U_b g(|s_b|)]$. In the case considered here the result for this asymmetry expressed in junction parameters reads

$$R_\Gamma = \exp\left[\frac{4\sqrt{2}\beta_{\rm eff}^3 E_C E_J^2 Q^2}{3(5+Q^2)} \frac{s_m(1-|s_b|)^2}{\sqrt{1+|s_b|}}\right], \qquad (13)$$

with the quality factor $Q = 1/\rho$ and the effective temperature $T_{\rm eff} = T + Q\hbar\omega_0 s_m/(4k_B\sqrt{1-s_b^2})$. Note that in the exponent the ratio between third cumulant and Gaussian noise basically appears via $E_J^2\beta_{\rm eff}^2 s_m$ so that $R_\Gamma \to 1$ for increasing s_m. The above finding not only allows to understand experimental data, but may be used to optimize the detection circuit as well. Namely, as a function of Q, the asymmetry R_Γ exhibits a maximum for intermediate Q values, but decreases towards higher (underdamped) and lower (overdamped) ones (Fig. 1). This also reveals that a rate calculation in either of these limiting cases is not sufficient. For typical experimental parameters $s_m = 2\omega_0\hbar\beta = 1$, $\beta E_J = 200$, and $s_b = 0.75$ one gains a maximal $R_\Gamma \approx 1.45$ at $Q \approx 2.5$. Asymmetry ratios for various s_m, see Fig. 1, lie around this value in agreement with recent experimental findings [11]. Note that for these parameters $T_{\rm eff}/T \approx 1.5$. Precise numerical simulations confirm the result (13) and will be discussed elsewhere [19]. The general results (11) and (12) determine the rate asymmetry due to the third cumulant also for other mesoscopic conductors such as diffusive wires or chaotic cavities. Since they apply for fast fluctuations, they provide a tool to analyze current noise measurements in the interesting high frequency regime [20].

To summarize, we have developed a formalism to describe the switching process of a JJ in the presence of weak and fast non-Gaussian fluctuations. From the corresponding rate expression the rate asymmetry due to the third moment of current noise has been obtained. Our results for the generalized FPE (6) and the escape rate (12) are applicable to other decay processes in physics and chemistry, where complex noise sources are present.

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