Statistics of temperature and potential fluctuations induced by coherent single particle sources

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***Abstract*—The prospect of time controlled information processing with individual electrons in nanoscale systems provides strong motivation for investigations of coupled charge and energy transport properties of single electron sources. Building on our recent work [F. Battista *et al.*, Phys. Rev. Lett. 110, 126602 (2013)] we investigate theoretically the statistical properties of temperature and potential fluctuations in an electronic probe coupled to a generic single electron source. A detailed derivation of the cumulant generating function of the joint probability distribution is presented. Moreover, the probability distribution in stationary phase approximation is analysed.**

I. Introduction

The ongoing miniaturization and close-packing of components in electronic devices has made it increasingly important to minimize and control heat dissipation in electronic systems on the nanometer size. In particular, any time-controlled information processing in nanoscale electronic devices will typically involve clocked application of voltage pulses and generation of time-dependent electrical currents, hence resulting in a time-varying heat flow. Extrapolating the electronics miniaturization trend, the ultimate carriers of information in nanosystems are single electrons. From a fundamental perspective, this motivates a careful investigation of coupled charge and energy transport properties of time-controlled sources for single electrons. In addition, the even more ambitious goal, to perform quantum information processing with single electrons in nanosystems, requires such electron on-demand sources as well as the transport through the system to preserve quantum coherence.

The need for a better understanding of the coupled charge and energy emission properties of electron on-demand sources is further emphasized by the recent rapid experimental progress [1], [2], [3], [4], [5], [6], [7] on fast, accurate single particle emitters, with operation frequencies reaching the GHz regime. Following the experiments, a large number of works investigated the charge transport properties of these coherent electron on-demand sources [8], [9], [10], [11], [12], [13], [14], [15], [16]. Coherent on-demand sources open up for quantum coherent few-electron experiments [17], [18], [19], [20] as well as put in prospect quantum information processing [21], [22], [23] with clocked single and entangled two-particle sources. Until recently, however, the energy emission properties of the sources received little attention [24], [25], [26], [27], [28], [29].

Importantly, although the low-frequency charge emission of ideal on-demand sources is noiseless, the emitted heat fluctuates (see illustration in Fig. I). These heat fluctuations are ubiquitous for quantum coherent sources; particles emitted during a time shorter than the drive period $\mathcal{T}$ have an uncertainty in energy larger than $h/\mathcal{T}$. The properties of the emitted quantum heat fluctuations and the resulting temperature and potential fluctuations induced in a hot-electron probe coupled to the source were recently investigated by us [30]. We considered a system implemented in a coherent conductor in the quantum Hall regime, see Fig. I. Our work provided a compelling illustration of the close connection between quantum thermoelectric transport and fundamental problems of nanoscale electronics. It also provided an important tool to investigate the heat fluctuation properties for different wavepackets and spectral profiles of electrons emitted from on-demand sources [10], [14]. These issues have very recently attracted considerable experimental interest [5], [31], [32], providing strong motivation for further theoretical investigations.

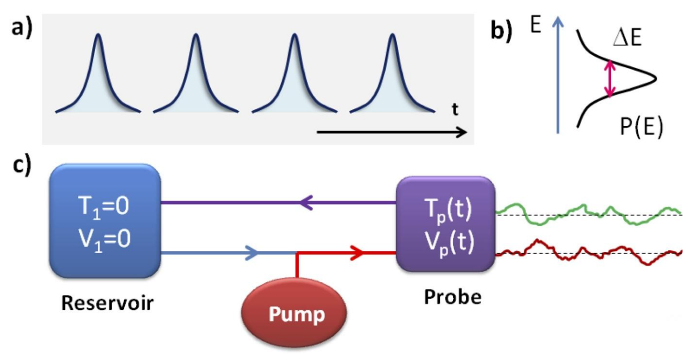

Fig. 1. a) Train of wavepackets emitted from a single electron pump. b) Energy probability distribution of emitted wavepacket. c) Schematic of the system considered in Ref. [30], with a single particle pump coupled to a thermally and electrically floating probe and an electronic reservoir, via edge states.

Here we present a detailed derivation of the joint probability distribution of the potential and temperature fluctuations in a thermally and electrically floating probe connected to the single electron source. We discuss the saddle point solution for the generating function for temperature and chemical potential

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cumulants and present the derivation of the probability distribution in the stationary phase approximation.

II. Statistics of isolated source

As a starting point we consider the joint probability distribution of energy and charge emitted from a single particle source under ideal operation conditions. The particles emitted from the source can be decribed by wavepackets, superpositions of plane waves at different energies $\epsilon > 0$, with amplitudes $c(\epsilon)$ ($\epsilon = 0$ at the Fermi surface). For measurements lasting a time t_0, much longer than the drive period $\mathcal{T}$, the joint probability $P_{t_0}(E,Q)$ that a total energy E and a charge Q has been emitted can conveniently be written

$$\begin{aligned} P_{t_0}(E,Q) &= \int d\lambda d\xi e^{i(\xi E+\lambda Q)+N[-i\sigma e\lambda+F(\xi)]}, \\ F(\xi) &= \ln\left[\int d\epsilon p(\epsilon) e^{-i\xi\epsilon}\right] \end{aligned} \tag{1}$$

where $p(\epsilon) = |c(\epsilon)|^2$, normalized as $\int d\epsilon p(\epsilon) = 1$, and $N[-i\sigma e\lambda + F(\xi)]$ the generating function for charge and energy cumulants with $N = t_0/\mathcal{T} \gg 1$ the number of particles emitted within the measurement time. Moreover, $\sigma = 0$ for sources emitting no net charge [1] and $\sigma = 1$ for sources emitting one electron [14] per cycle. The form in Eq. (1) highlights the fact that the charge transfer is noiseless, with only the first cumulant finite, while the transferred energy fluctuates.

III. Source-probe system

Our goal is to investigate how the heat, or energy, fluctuations described by Eq. (1) are manifested in quantities typically accessible via measurements in mesoscopic systems. To this aim we consider a system where the source is coupled to a probe, via the lower edge state of a conductor in the quantum Hall regime, see Fig. I. The probe is in the hot-electron regime, with a floating electron temperature $T_p(t)$ and chemical potential $\mu_p(t)$. Particles emitted from the probe flow along the upper edge into an electronic reservoir electrically grounded and kept at zero temperature.

The potential $\mu_p(t)$ fluctuates on the time scale given by the RC-time, τ_{RC}, while the temperature $T_p(t)$ typically fluctuates on the time scale of the dwell-time in the probe, τ_d. For the hot-electron regime assumed here, the system is in the limit $\tau_{e-e} \ll \tau_{RC}, \tau_d \ll \tau_{e-ph}$. Moreover, the time scales of the fluctuations of the probe, τ_d, τ_{RC}, are considered to be much longer than the drive period $\mathcal{T}$ but much shorter than the measurement time t_0.

IV. Statistics of probe fluctuations

The object of key interest is the joint probability distribution of potential and temperature fluctuations at the probe. To arrive at the distribution we first introduce the dimensionless potentials and temperatures,

$$\mu = \frac{1}{h}\int_0^{t_0} dt \mu_p(t), \qquad T = \frac{1}{h}\int_0^{t_0} dt k_b T_p(t) \tag{2}$$

corresponding to the fluctuating $\mu_p(t), T_p(t)$ integrated over the measurement time t_0. The joint probability distribution $\mathcal{P}_{t_0}(\mu, T)$ can be conveniently written in terms of a cumulant generating function $G(\chi,\theta)$ [33] as

$$\mathcal{P}_{t_0}(\mu,T) = \frac{1}{(2\pi)^2}\int d\chi \int d\theta e^{-i\theta T - i\chi\mu + G(\chi,\theta)}, \tag{3}$$

with χ and θ counting fields for μ and T respectively. From $G(\chi,\theta)$ the cumulants of the joint probability distribution $\mathcal{P}_{t_0}(\mu,T)$ are then, by construction, obtained from successive derivatives with respect to the counting fields

$$t_0\langle\delta T_p^n \delta\mu_p^m\rangle = (-ih)^{n+m} k_b^{-n} \frac{\partial^n\partial^m G(\chi,\theta)}{(\partial\theta)^n(\partial\chi)^m}\bigg|_{\chi,\theta=0}. \tag{4}$$

To determine $G(\chi,\theta)$ we first introduce an intermediate time scale τ, such that $\mathcal{T} \ll \tau \ll \tau_d, \tau_{RC}$. On the time scale τ the statistics of net transferred energy E_p and charge Q_p in the probe can be described by the source generating function $\tau h_s(\lambda,\xi) = (\tau/\mathcal{T})[-ie\sigma\lambda + F(\xi)]$ where $F(\xi)$ is given in Eq. (1). The probe generating function $\tau h_p(\lambda,\xi,E_p,Q_p)$ is given by [33], [34]

$$\begin{aligned} h_p = \frac{1}{h}\int d\epsilon \Big[&\ln[1+f_p(\epsilon)(e^{ie\lambda+i\epsilon\xi}-1)] + \\ &\ln[1+f(\epsilon)(e^{-ie\lambda-i\epsilon\xi}-1)]\Big], \end{aligned} \tag{5}$$

where $f_p(\epsilon) = f_p(\epsilon,\mu_p,T_p)$ and $f(\epsilon)$ are the probe and the reservoir Fermi distribution functions and ξ and λ are the counting fields for E_p and Q_p respectively. The energy integral in Eq. (5) can be carried out [35], giving

$$h_p = \frac{1}{2h}\frac{2\mu_p(ei\lambda) + k_b T_p(ei\lambda)^2 + i\xi\left[\pi^2(k_bT_p)^2/3 + \mu_p^2\right]}{1-(k_bT_p)i\xi}. \tag{6}$$

Importantly, the energy E_p and charge Q_p are related to T_p and μ_p as

$$E_p = \nu\left[\mu_p^2/2 + \pi k_b T_p)^2/6\right], \quad Q_p = \nu e\mu_p, \tag{7}$$

where ν is the density of states in the probe. Working within the framework of the stochastic path integral formalism [35], [36], we can then express $G(\chi,\theta)$ as a path integral over all configurations of E_p and charge Q_p during the measurement. In the long time limit we have

$$\begin{aligned} e^{G(\chi,\theta)} &= \int dQ_p dE_p d\lambda d\xi e^{S(Q_p,E_p,\lambda,\xi)} \\ S &= t_0[i(\theta k_b T_p + \chi\mu_p)/h + h_p + h_s] \end{aligned} \tag{8}$$

Along similar lines as in Refs. [25], [34], [35], the integral in Eq. (8) can be solved in the saddle point approximation. The leading corrections to the saddle point solution are typically an order $\mathcal{T}/\tau_d, \mathcal{T}/\tau_{RC} \ll 1$ smaller and hence negligible [36]. The saddle point equations are given by

$$\frac{\partial S}{\partial\xi} = 0, \quad \frac{\partial S}{\partial\lambda} = 0, \quad \frac{\partial S}{\partial E_p} = 0, \quad \frac{\partial S}{\partial Q_p} = 0 \tag{9}$$

These relations together with Eqs. (7) and (8) allow us to write

$$\begin{aligned} i\frac{\chi}{h}+\frac{\partial h_p}{\partial \mu_p} &= 0, \quad \frac{\partial (h_p+h_s)}{\partial \lambda}=0 \\ i\frac{k_b\theta}{h}+\frac{\partial h_p}{\partial T_p} &= 0 \quad \frac{\partial (h_p+h_s)}{\partial \xi}=0. \end{aligned} \tag{10}$$

Denoting the solutions to the saddle point equations $\lambda^*, \xi^*, \mu_p^*$ and T_p^*, functions of χ and θ, we first find from Eqs. (10) the relations

$$e\lambda^* = -\chi - \sigma\hbar\omega\xi^*, \qquad \mu_p^* = i\chi k_b T_p^* + \sigma\hbar\omega. \tag{11}$$

Inserting these results back into Eqs. (10) we can further write

$$k_b T_p^* = \frac{6q}{\pi^2 g + 3q}\frac{1}{i\xi^*} \tag{12}$$

where we introduced $g = g(\chi,\theta)$ and q as

$$g = \frac{3}{\pi^2}\left[\frac{(i\chi)^2}{2}+i\theta\right], \qquad q = -\frac{\pi^2}{6}\left(1-\sqrt{1-2g}\right)^2. \tag{13}$$

Finally, the remaining variable ξ^* is obtained from the relation

$$q = z^2\left[\frac{dF}{dz}+\frac{\sigma}{2}\right], \tag{14}$$

where we introduced $z = i\xi^*\hbar\omega$. Note that this equation does not have a general solution for z, since it depends on the specific properties of the single particle source via $F(z)$. By inserting the saddle point solutions back into S we then obtain the generating function, $G(\chi,\theta) = S(\mu_p^*, T_p^*, \lambda^*, \xi^*)$, as

$$G(\chi,\theta) = N\left[\frac{d[zF(z)]}{dz}+\sigma(z+i\chi)\right] \tag{15}$$

recalling that $N = t_0/\mathcal{T}$.

From the relation in Eq. (4) we obtain the lowest order cumulants of the joint probability distribution. The first order derivatives lead to the temperature and chemical potential average values,

$$\bar{\mu}_p = \sigma\hbar\omega, \quad \bar{T}_p = \sqrt{\frac{1}{g_0 l_0 \mathcal{T}}[2\langle\epsilon\rangle - \sigma\hbar\omega]}, \tag{16}$$

where $g_0 = e^2/h$ is the (single spin) conductance quantum and $l_0 = (\pi k_B/e)^2/3$ the Lorentz number. We stress that $\bar{\mu}_p$ and $\bar{T}_p$ depend only on the source properties ω and $\langle\epsilon\rangle$ and fundamental constants [5]. Moreover, we note that $\langle\epsilon\rangle > \hbar\omega/2$ follows along the lines of Ref. [37].

From the second order derivatives we obtain the temperature and chemical potential fluctuations

$$\begin{aligned} \langle(\delta\mu_p)^2\rangle &= hk_b\bar{T}_p, \quad \langle\delta\mu_p\delta T_p\rangle = 0 \\ \langle(\delta T_p)^2\rangle &= \frac{1}{g_0 l_0}\left[k_b\bar{T}_p + \frac{1}{2}\frac{(\Delta\epsilon)^2}{\langle\epsilon\rangle - \sigma\hbar\omega/2}\right]. \end{aligned} \tag{17}$$

Several interesting conclusions can be drawn from this result:

• The potential fluctuations $\langle(\delta\mu_p)^2\rangle$ are proportional to the average temperature $\bar{T}_p$. This fluctuation-dissipation type relation forms the basis for noise thermometry, i.e. to determine the system temperature from noise measurements [38].

• In contrast, the temperature fluctuations $\langle(\delta T_p)^2\rangle$ are a sum of two physically distinct terms. The first term, proportional to $\bar{T}_p$, results from the finite temperature of the probe. It is a classical contribution, i.e. it would be present even if the injected particles had a well defined energy, $\Delta\epsilon = 0$. The second term is proportional to $(\Delta\epsilon)^2$. It is a quantum contribution, i.e. a direct result of the uncertainty of the energy of the injected particle.

• There are no correlation between the voltage and the temperature fluctuations. However, higher order cumulants of the fluctuations are typically non-zero.

Importantly, Eqs. (16-17) are in agreement with the first and second cumulants of temperature and potential fluctuations obtained within a Boltzmann-Langevin approach in Ref. [30].

For the higher order fluctuations an important relation can be derived, namely that all the even chemical potential cumulants can be expressed in terms of the temperature cumulants as

$$\langle(\delta\mu_p)^{2n}\rangle = (2n-1)!!(k_b h)^n\langle(\delta T_p)^n\rangle. \tag{18}$$

Of particular importance for an experimental realization of our proposal, the lowest order potential fluctuation cumulant that provides direct information about the quantum heat fluctuations of the source is $\langle(\delta\mu_p)^4\rangle$, since it is directly proportional to $\langle(\delta T_p)^2\rangle$. Formally, the relation in Eq. (18) is a consequence of the counting fields in the cumulant generating function entering only via $g(\chi,\theta)$, Eqs. (13). Moreover, the factorial growth of the cumulants with order n is in agreement with general predictions in Ref. [39].

V. Probability distribution, stationary phase approximation

During the measurement time t_0 a large number of particles are pumped, $N \gg 1$. Since N enters as a prefactor in the cumulant generating function $G(\chi,\theta)$ in Eq. (15), the integral in the probability distribution $\mathcal{P}_{t_0}(\mu, T)$ can be evaluated in the saddle point, or stationary phase, approximation. The probability distribution is then obtained with exponential accuracy. The saddle point equations for θ and χ, similar to Eqs. (10), are given by

$$-iT + \frac{\partial G(\chi,\theta)}{\partial\theta} = 0, \qquad -i\mu + \frac{\partial G(\chi,\theta)}{\partial\chi} = 0. \tag{19}$$

Making use of that in $F(z)$, determining the first term in $G(\chi,\theta)$, the counting fields χ,θ enter only via $g(\chi,\theta)$, and denoting the saddle point solutions χ^*, θ^*, we directly find

$$\chi^* = i\left[\sigma N - \mu\right]/T. \tag{20}$$

Further, substituting Eqs. (13-14) into Eq. (19) we find

$$\left.\frac{dF}{dz}\right|_{z=z^*} = -\frac{\pi^2}{6}\left(\frac{Tq^*}{N}\right)^2 - \frac{\sigma}{2}, \tag{21}$$

where we introduced

$$z^* = N\frac{q^*-1}{q^*T}, \qquad q^* = \sqrt{1-i\frac{6\theta^*}{\pi^2}}. \tag{22}$$

Eq. (21) constitutes an implicit relation for θ^*. Now, inserting the saddle point solutions χ^* and θ^* back into the exponent of the joint probability distribution in Eq. (3) we arrive, after some algebra, at the logarithmic probability

$$\ln \mathcal{P}_{t_0}(\mu, T) = -iT\theta^* + G(0, \theta^*) - (\mu - \bar{\mu})^2/2T \qquad (23)$$

where $\bar{\mu} = t_0\bar{\mu}_p/h = \sigma N$ is the average chemical potential. Since θ^* is a function of temperature T only, Eq. (23) shows that the potential μ displays Gaussian fluctuations, of width $\sqrt{T}$, around the average $\bar{\mu}$ for any given temperature T. Importantly, this holds for an arbitrary single particle source. As a consequence, the marginal potential distribution $\mathcal{P}_{t_0}(\mu) = \int dT\mathcal{P}_{t_0}(\mu, T)$ is symmetric around $\bar{\mu}$, *i.e.* odd cumulants are zero. It should however be noted that $\mathcal{P}_{t_0}(\mu)$ is not Gaussian, it is a weighted sum of Gaussian distributions with different widths.

The marginal temperature distribution $\mathcal{P}_{t_0}(T) = \int d\mu\mathcal{P}_{t_0}(\mu, T)$ is given by

$$\mathcal{P}_{t_0}(T) = -iT\theta^* + G(0, \theta^*). \qquad (24)$$

As is clear from Eq. (23), the marginal temperature distribution contains the two non-trivial, source dependent terms in the logarithmic joint probability distribution. Accordingly, it is a function of θ^* only, defined from Eq. (21). It is therefore natural to focus the investigations on $\mathcal{P}_{t_0}(T)$. As a basis for further investigations, we discuss here the simplest possible case, a source where the energy of the particles is not fluctuating, i.e. $p(\epsilon) = \delta(\epsilon - \langle\epsilon\rangle)$. This classical case corresponds to a source generating function $F(z) = z\langle\epsilon\rangle/\hbar\omega$. From Eqs. (21) and (24) we then, after some algebra, arrive at the simple result

$$\ln \mathcal{P}_{t_0}(T) = -\frac{\pi^2}{6}T\left(1 - \frac{\bar{T}}{T}\right)^2 \qquad (25)$$

where we defined, similarly to $\bar{\mu}$, the average temperature $\bar{T} = t_0k_b\bar{T}_p/h$. We see from Eq. (25) that for small fluctuations $T - \bar{T} \ll \bar{T}$ the distribution is Gaussian while for $T \ll \bar{T}$ the probability is suppressed $\mathcal{P}_{t_0}(T) \propto e^{-\pi^2\bar{T}^2/(6T)}$, guaranteeing $\mathcal{P}_{t_0}(T) \to 0$ for $T \to 0$. The probability for large fluctuations $T \gg \bar{T}$ is suppressed as $\mathcal{P}_{t_0}(T) \propto e^{-T\pi^2/6}$. Some examples of probability distributions which deviate from this classical result were presented in our paper [30].

In conclusion, we have presented a detailed derivation of the full statistical distribution of potential and temperature fluctuations induced in a hot electron probe coupled to a single particle source in the quantum Hall regime. This derivation extends on the results in our earlier work [30] and provides additional insight into the properties and nature of the fluctuations.

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