Mesoscopic Quadratic Quantum Measurements

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We develop a theory of quadratic quantum measurements by a mesoscopic detector. It is shown that the quadratic measurements should have nontrivial quantum information properties, providing, for instance, a simple way of entangling two noninteracting qubits. We also calculate the output spectrum of a detector with both linear and quadratic response, continuously monitoring two qubits.

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The problem of quantum measurements with mesoscopic solid-state detectors attracts considerable current interest (see, e.g., chapters on quantum measurements in [1]). This interest is motivated in part by the important role of measurement in quantum computing and in part by the possibility, provided by the mesoscopic structures, to study the transition between quantum and classical behavior in systems that are large on the atomic scale. Although mesoscopic detectors can be quite different and include, e.g., quantum point contacts (QPC) [2–8], normal and superconducting single-electron transistors (SETs) [9–15], SQUID magnetometers [16], and generic mesoscopic conductors [17,18], the operating principle of almost all of them is the same. A measured quantum system controls the transmission amplitude t of some particles (electrons, Cooper pairs, or magnetic flux quanta) between the two reservoirs, and their flux provides information on the state of the system [19]. In general, the amplitude t varies together with some control operator x, and for sufficiently weak detector-system coupling, the dependence $t(x)$ can be approximated as linear. The dynamics of such linear measurements is well understood (see, e.g., [20,21]).

At some special bias points, however, the linear response coefficient of the $t(x)$ dependence vanishes and this dependence becomes quadratic. This can happen, for instance, if the amplitude t is formed by two or more interfering tunneling trajectories. Known examples of such situations include dc SQUIDs and superconducting SETs. In this work, we show that a quantum detector operating at such a special point should enable measurements of product operators referring to separate systems and have nontrivial quantum information processing properties, e.g., create a simple entanglement mechanism for noninteracting qubits. Specifically, we consider the measurement of two qubits (Fig. 1), which is the simplest system that reveals the characteristics of quadratic detection. The two qubits are assumed to be coupled to one detector through their basis-forming variables σ_z^j, $j = 1, 2$, i.e., $x = c_1\sigma_z^1 + c_2\sigma_z^2$ so that

$$t(x) = t_0 + \delta_1\sigma_z^1 + \delta_2\sigma_z^2 + \lambda\sigma_z^1\sigma_z^2. \tag{1}$$

The last term in this equation appears due to the nonlinearity of $t(x)$. For the two qubits, Eq. (1) represents the most general dependence of t on σ_z^j, whereas for measurements of other systems, Eq. (1) can be justified as the Taylor's expansion in the weak detector-system coupling. Higher-order terms in this expansion would affect measurements of a larger number of qubits.

The Hamiltonian of the detector-qubit system is

$$H_t = H_0 + H_d + t(\{\sigma_z^j\})\xi + t^\dagger(\{\sigma_z^j\})\xi^\dagger, \tag{2}$$

where $H_0 = -(1/2)\sum_{j=1,2}(\varepsilon_j\sigma_z^j + \Delta_j\sigma_x^j) + (\nu/2)\sigma_z^1\sigma_z^2$. Here Δ_j is the tunnel amplitude and ε_j is the bias of the jth qubit, ν is the qubit interaction energy, H_d is the detector Hamiltonian, and $\xi^\dagger$, ξ are the detector operators that create excitations when a particle is transferred forward or backward between the detector reservoirs. For instance, for the QPC detector, $\xi^\dagger$, ξ describe excitation of electron-hole pairs in the QPC electrodes.

We make two assumptions about the detector: the tunneling between reservoirs is weak and can be described in the lowest nonvanishing order in t, and the characteristic time scale of tunneling is much shorter than that of the qubit evolution due to H_0. For the QPC detector this means that the QPC is in the tunneling regime and the voltage across it is much larger than the

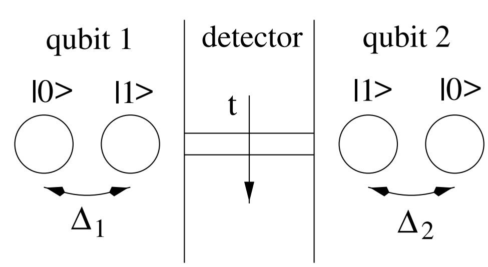

FIG. 1. Diagram of a mesoscopic detector measuring two qubits. The qubits modulate amplitude t of the tunneling of detector particles between the two reservoirs.

qubit energies. Under these assumptions, the dynamics of measurement depends only on the correlators

$$\gamma_+ = \int_0^\infty dt \langle \xi(t)\xi^\dagger \rangle, \qquad \gamma_- = \int_0^\infty dt \langle \xi^\dagger(t)\xi \rangle, \quad (3)$$

which set the scale $\Gamma_\pm \equiv 2\,\mathrm{Re}\,\gamma_\pm$ of the forward and backward detector tunneling rates. In Eq. (3), the angled brackets denote averaging over internal degrees of freedom of the detector reservoirs taken to be in a stationary state. The correlators $\langle \xi(t)\xi \rangle$, $\langle \xi^\dagger(t)\xi^\dagger \rangle$ that do not conserve the number of particles are assumed to vanish.

The measurement contribution to the evolution of the qubit density matrix ρ is obtained by standard lowest-order perturbation theory in tunneling. To describe qubit dynamics conditioned on a particular outcome of the measurement, we keep in the evolution equation the number of particles n transferred through the detector. Since the correlators that do not conserve n vanish, only the terms diagonal in n contribute to the evolution. In the "measurement" basis of eigenstates of the σ_z^j operators, $|\uparrow\uparrow\rangle$, $|\uparrow\downarrow\rangle$, $|\downarrow\uparrow\rangle$, and $|\downarrow\downarrow\rangle$, in which each state $|k\rangle$ is characterized by the value t_k of the transmission amplitude (1), $t_1 = t_0 + \delta_1 + \delta_2 + \lambda$, $t_2 = t_0 + \delta_1 - \delta_2 - \lambda$, etc., the measurement contribution to $\dot\rho$ is

$$\dot\rho_{kl}^n = -(1/2)(\Gamma_+ + \Gamma_-)(|t_k|^2 + |t_l|^2)\rho_{kl}^n + \Gamma_- t_k t_l^* \rho_{kl}^{n+1} + \Gamma_+ t_k^* t_l \rho_{kl}^{n-1} - i[\delta H, \rho^n]_{kl}. \quad (4)$$

Here $\delta H = \sum_j \delta\varepsilon_j \sigma_z^j + \delta\nu \sigma_z^1 \sigma_z^2$ is the renormalization of the qubit Hamiltonian due to coupling to the detector: $\delta\varepsilon_j = \mathrm{Re}(\delta_j t_0^* + \delta_{j'}\lambda^*)\,\mathrm{Im}(\gamma_- + \gamma_+)$ and $\delta\nu = \mathrm{Re}(\delta_1\delta_2^* + t_0\lambda^*)\,\mathrm{Im}(\gamma_- + \gamma_+)$, where $j, j' = 1, 2$, $j' \neq j$. Equation (4) is the basis for our quantitative discussion of quadratic measurements. It generalizes to an arbitrary detector and two qubits, the equation obtained in [5] for a qubit measured with the QPC in the tunnel regime.

Disregarding the index n in Eq. (4), we obtain the equation for the measurement-induced evolution of the qubit density matrix averaged over different measurement outcomes. Together with the evolution due to the qubit Hamiltonian H_0 this equation is

$$\dot\rho_{kl} = -\gamma_{kl}\rho_{kl} - i[H_0, \rho]_{kl}. \quad (5)$$

Here $\gamma_{kl} \equiv (1/2)(\Gamma_+ + \Gamma_-)|t_k - t_l|^2$, and we included in H_0 two renormalization terms: δH (4) and $\delta H'$ due to phases $\varphi_{kl} \equiv \arg(t_k t_l^*)$ of the transfer amplitudes in Eq. (4) defined by $[\delta H', \rho]_{kl} = (\Gamma_+ - \Gamma_-)|t_k t_l| \sin\varphi_{kl}\rho_{kl}$.

Evolution (5) of the qubit density matrix is reflected in the detector current. The form of the current I operator in the qubit space is obtained by the same lowest-order perturbation theory in tunneling that leads to Eq. (5):

$$I = (\Gamma_+ - \Gamma_-)t^\dagger t. \quad (6)$$

This equation can be used to calculate both the dc current $\langle I \rangle = \mathrm{Tr}\{I\rho_0\}$, where ρ_0 is the stationary solution of Eq. (5), and the current spectral density

$$S_I = S_0 + 2\int_0^\infty d\tau \cos\omega\tau (\mathrm{Tr}\{I e^{L\tau}[I\rho_0]\} - \langle I \rangle^2). \quad (7)$$

Here $S_0 = (\Gamma_+ + \Gamma_-)\,\mathrm{Tr}\{t^\dagger t \rho_0\}$, and $e^{L\tau}[A]$ denotes the evolution of the matrix A during time interval τ governed by Eq. (5).

Decay of the off-diagonal matrix elements Eq. (5) is the result of averaging over the measurement outcomes. However, since n is the classical detector output, it is legitimate to ask a question about the qubit evolution for a specific measurement outcome n. Such a "conditional" description of the measurement dynamics (see, e.g., [21]) is convenient for the calculation of more complicated correlators involved, for instance, in problems of feedback control of the measured system. In our case, it is obtained by first solving Eq. (4) in terms of n. Noticing that Eq. (4) coincides in essence with the recurrence relations for the modified Bessel functions I_n and assuming the initial condition $\rho_{kl}^n(0) = \rho_{kl}(0)\delta_{n,0}$ we get

$$\rho_{kl}^n(\tau) = \rho_{kl}(0)(\Gamma_+/\Gamma_-)^{n/2} I_n(2\tau|t_k t_l|\sqrt{\Gamma_+\Gamma_-}) \times \exp\{-(1/2)(\Gamma_+ + \Gamma_-)(|t_k|^2 + |t_l|^2)\tau - in\varphi_{kl}\}. \quad (8)$$

Following the same steps as in [21], the qubit density matrix conditioned on the particular "observed" number n of transferred particles is obtained then by selecting the term with this n in Eq. (8) and normalizing the resulting reduced density matrix. For weak detector-qubit coupling, $|\delta_j|, |\lambda| \ll |t_0|$, when individual tunneling events do not provide significant information on the qubit state, it is convenient to condition the evolution on the quasi-continuous current $I(t)$ in the detector. Then, the "Bayesian" equation for the qubit density matrix is

$$\dot\rho_{kl} = -i[H_0, \rho]_{kl} - \gamma_{kl}\rho_{kl} + I_f(t)\rho_{kl}\left[\frac{1}{2S_0}\sum_j \rho_{jj}(I_k + I_l - 2I_j) - i\varphi_{kl}\right], \quad (9)$$

where S_0 is the background current noise [see Eq. (7)], variation of which with the qubit state can be neglected in the weak-coupling limit, $S_0 = (\Gamma_+ + \Gamma_-)|t_0|^2$. Also, $I_k = (\Gamma_+ - \Gamma_-)|t_k|^2$ is the average detector current in the qubit state k, and $I_f(t) = I(t) - \sum_k \rho_{kk} I_k$ is the noise of the detector current. Equation (9) is written in the Itô form, in which averaging over $I_f(t)$ can be done by simply omitting the terms with it. In the weak-coupling regime, $\gamma_{kl} = (1/2)(\Gamma_+ + \Gamma_-)[(|t_k| - |t_l|)^2 + \varphi_{kl}^2|t_0|^2]$. It is the same ensemble-averaged decoherence rate as in Eq. (5), but in Eq. (9) it leads to decoherence only after averaging over $I_f(t)$.

We now use the equations obtained above to discuss several quantitative characteristics of quadratic measurements. We start with the purely quadratic detectors, when $\delta_j = 0$, so that $I_1 = I_4 = (\Gamma_+ - \Gamma_-)|t_0 + \lambda|^2 \equiv I_{\uparrow\uparrow}$ and

$I_2 = I_3 = (\Gamma_+ - \Gamma_-)|t_0 - \lambda|^2 \equiv I_{\uparrow\downarrow}$. In this case, if the qubits are *stationary,* $H_0 = 0$, the detector effectively measures the product operator $\sigma_z^1\sigma_z^2$ of the two qubits. That is, on the time scale of measurement time $\tau_m = 4S_0/(I_{\uparrow\uparrow} - I_{\uparrow\downarrow})^2$, the subspace $\{|1\rangle, |4\rangle\}$, in which the states of the two qubits are the same and the average detector current is $I_{\uparrow\uparrow}$, is distinguished from the subspace $\{|2\rangle, |3\rangle\}$ in which the states of the two qubits are opposite and the current is $I_{\uparrow\downarrow}$, while the states within these subspaces are not distinguished. This property of quadratic measurements can be used to design a simple error-correction scheme for dephasing errors [22].

Next, we consider the case of *identical, unbiased, noninteracting* qubits with nonvanishing Hamiltonian, $H_0 = -(\Delta/2)\sum_j \sigma_x^j$. In this case the two degenerate zero-energy eigenstates of H_0 can be chosen as $\{|\uparrow\uparrow\rangle - |\downarrow\downarrow\rangle, |\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle\}$. In the remaining subspace that is denoted D_+, in the basis $\{|\uparrow\uparrow\rangle + |\downarrow\downarrow\rangle, |\uparrow\downarrow\rangle + |\downarrow\uparrow\rangle\}$, H_0 reduces to $-\Delta\sigma_x$ and mixes the states with similar and opposite states of the two qubits. Accordingly, there are three possible measurement outcomes characterized by the different dc currents $\langle I\rangle$ in the detector, $I_{\uparrow\uparrow}$, $I_{\uparrow\downarrow}$, and $(I_{\uparrow\uparrow} + I_{\uparrow\downarrow})/2$. These outcomes can be interpreted as a measurement of the operator $\sigma_y^1\sigma_y^2 + \sigma_z^1\sigma_z^2$. Conditional Eq. (9) can be used to simulate how the qubits, on the time scale $\simeq 4S_0/I_a^2$, $I_a \equiv (I_{\uparrow\uparrow} - I_{\uparrow\downarrow})/2$, are driven into one of the three outcomes driven by the specific realization of the detector current. The probabilities of different outcomes depend on the initial state. In the first two outcomes, the initial state is projected on one of the fully entangled states of the two qubits, e.g.,

$$\langle I\rangle = I_{\uparrow\downarrow} \Leftrightarrow |\psi\rangle = (|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle)/\sqrt{2}. \qquad (10)$$

Thus, quadratic measurements of two symmetric qubits provide a simple way of generating entangled states of qubits that in contrast to linear measurements [23] is based only on monitoring the dc current.

In the third scenario, when $\langle I\rangle = (I_{\uparrow\uparrow} + I_{\uparrow\downarrow})/2$, the two qubits are confined to the subspace D_+ and perform coherent quantum oscillations. Equation for the density matrix (5) reduced to D_+ is

$$\dot{\rho}_{kl} = i\Delta[\sigma_x, \rho]_{kl} - \gamma\begin{pmatrix} 0 & \rho_{12} \\ \rho_{21} & 0 \end{pmatrix}, \qquad (11)$$

where $\gamma \equiv 2(\Gamma_+ + \Gamma_-)|\lambda|^2$. Solving this equation and using the fact that the current operator (6) is reduced in D_+ to $I_a\sigma_z$, we find the current spectral density (7):

$$S_I(\omega) = S_0 + \frac{8\Delta^2 I_a^2 \gamma}{(\omega^2 - 4\Delta^2)^2 + \gamma^2\omega^2}, \qquad (12)$$

where $S_0 = (\Gamma_+ + \Gamma_-)(|t_0|^2 + |\lambda|^2)$. Qualitatively, for $\gamma \ll \Delta$, spectral density (12) describes coherent oscillations of the two qubits with the frequency 2Δ, twice the oscillation frequency in one qubit. Similar to the case of linear measurements [24], the maximum of the ratio of the oscillation peak to noise S_0 is 4. As one can see from Eq. (12), this maximum is reached in the case of weak measurement $|\lambda| \ll |t_0|$ by the "ideal" detector for which $\arg(t_0\lambda^*) = 0$ and only Γ_+ or Γ_- is nonvanishing. An interesting feature of the regime of continuous measurement required to observe the oscillation spectra similar to Eq. (12) is that it is sensitive to the detector backaction. This means that, in contrast to the dc "single-shot" measurements [as, e.g., in Eq. (10)], continuous measurements provide information about the quantum properties of the detector itself.

For *different qubit tunnel amplitudes*, transitions [with the rate $(\Delta_1 - \Delta_2)^2/2\gamma$ for small $\Delta_1 - \Delta_2$] between the states $|\uparrow\uparrow\rangle - |\downarrow\downarrow\rangle$ and $|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle$ mix the measurement outcomes $I_{\uparrow\uparrow}$ and $I_{\uparrow\downarrow}$. This means that for $\Delta_1 \neq \Delta_2$ there are only two outcomes that have the same dc current and differ by current spectral densities. In one, the qubits are again in the subspace D_+ and the spectral density is given by Eq. (12) where now $2\Delta \to \Delta_1 + \Delta_2$. In the other, the qubits are confined to the subspace orthogonal to D_+, and both the qubit dynamics and the current spectral density are given by the same Eqs. (11) and (12) with $\Delta \to (\Delta_1 - \Delta_2)/2$.

As the last application of the general theory, we consider two identical qubits measured by a weakly and symmetrically coupled detector with *arbitrary nonlinearity.* It is convenient to discuss this situation in terms of the total effective spin S of the two qubits, which determines the amplitude (1) of detector tunneling:

$$t = t_0 + 2\delta S_z + \lambda(2S_z^2 - 1). \qquad (13)$$

The $S = 0$ state (10) does not evolve in time under the qubit Hamiltonian and represents one of the measurement outcomes characterized by the dc detector current $I_{\uparrow\downarrow}$ and flat current spectral density $S_I(\omega) = (\Gamma_+ + \Gamma_-)|t_0 - \lambda|^2$. Three other, $S = 1$, states are mixed by measurement and represent the second measurement outcome. We take the basis of the $S = 1$ subspace as three energy eigenstates $S_x = -1, 0, 1$ with energies $\{\Delta, 0, -\Delta\}$. The detector induces transitions between these states with the rate independent of the transition's direction, so that the stationary qubit density matrix in this subspace is $\rho_0 = 1/3$. Equations (6) and (13) show then that the dc detector current for this outcome is $\langle I\rangle = (\Gamma_+ - \Gamma_-) \times [2(|t_0|^2 + |\lambda|^2) + |t_0 + \lambda|^2 + 8|\delta|^2]/3$, and can be written as $\langle I\rangle = (I_{\uparrow\uparrow} + I_{\downarrow\downarrow} + I_{\uparrow\downarrow})/3$, where the currents $I_{\uparrow\uparrow}, \ldots$ are introduced in the same way as before, e.g., $I_{\uparrow\uparrow} = (\Gamma_+ - \Gamma_-)|t_0 + 2\delta + \lambda|^2$. The background current noise S_0 coincides with $\langle I\rangle$ with $\Gamma_+ - \Gamma_-$ replaced by $\Gamma_+ + \Gamma_-$.

The system performs oscillations at frequencies Δ and 2Δ whose spectral peaks should have Lorentzian form for weak detector-qubit coupling. Evaluating the current matrix elements from Eqs. (6) and (13), and evolution of the density matrix from Eq. (5) reduced to the $S = 1$ subspace, we obtain parameters of these Lorentzians in the spectral density (7):

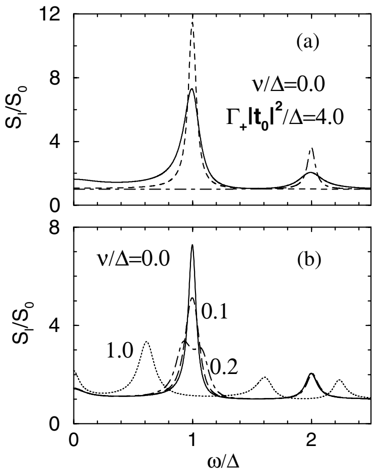

FIG. 2. Spectra of coherent quantum oscillation in two qubits. (a) Transition from linear to quadratic measurement: dashed line, solid line, and dot-dashed line correspond to linear, ($\delta = 0.1t_0$, $\lambda = 0$), intermediate ($\delta = \lambda = 0.1t_0$), and quadratic ($\delta = 0$, $\lambda = 0.1t_0$) cases. (b) Effect of qubit-qubit interaction on spectrum for $\Gamma_+|t_0|^2 = 2.0\Delta$ and $\delta = \lambda = 0.1t_0$.

$$\omega \simeq j\Delta, \qquad S_I(\omega) = S_0 + \frac{2}{3}\,\frac{a_j^2\gamma_j}{(\omega - j\Delta)^2 + \gamma_j^2}, \qquad (14)$$

$$j = 1, 2, \qquad \gamma_j = (\Gamma_+ + \Gamma_-)(j|\delta|^2 + |\lambda|^2),$$

$$a_1 = 4(\Gamma_+ - \Gamma_-)\,\mathrm{Re}[(t_0 + \lambda)\delta^*] = (I_{\uparrow\uparrow} - I_{\downarrow\downarrow})/2,$$

$$a_2 = 2(\Gamma_+ - \Gamma_-)(\mathrm{Re}[t_0\lambda^*] + |\delta|^2) = (I_{\uparrow\uparrow} + I_{\downarrow\downarrow} - 2I_{\uparrow\downarrow})/4.$$

Note that condition $\rho_0 = 1/3$, used in Eq. (14) and Fig. 2, implies that the coefficient δ of linear measurement mixing all three $S = 1$ states, does not vanish identically. Otherwise, only two states spanning the subspace D_+ defined above are mixed, and $\rho_0 = 1/2$ in this subspace, as assumed in Eq. (12). In general, there is also a spectral peak at $\omega = 0$ caused by switching between states with different average currents. For small but finite δ, $\delta \ll \lambda$, this peak can be very high, e.g., in the case of an "ideal" detector ($\Gamma_- = \arg t_0\lambda^* = \arg t_0\delta^* = 0$) its height and half-width are, respectively, $(8\lambda^2/27\delta^2)S_0$ and $6\delta^2\Gamma_+$.

Figure 2 shows current spectral density calculated from Eqs. (5) and (7) for an ideal detector without making a weak-coupling approximation. Figure 2(a) illustrates the transition between "single-qubit" oscillations at $\omega \simeq \Delta$ in the case of linear measurement and oscillations at $\omega \simeq 2\Delta$ for the quadratic measurement. One can see that in agreement with Eq. (14) the $\omega \simeq \Delta$ peak is typically higher than the quadratic peak at double frequency. It is at the same time more sensitive to qubit-qubit interaction as illustrated in Fig. 2(b). Even weak interaction $\nu \ll \Delta$ first suppresses and then splits the peak at $\omega \simeq \Delta$ in two while affecting the quadratic peak only slightly.

In summary, we discussed quadratic quantum measurements and demonstrated their nontrivial quantum-information properties. We also calculated output spectra of the quadratic detector measuring coherent oscillations in two qubits. Consistent with the case of classical oscillations, quadratic measurement results in the spectral peak at frequency that is twice the frequency of individual qubit oscillations. Quadratic measurements should be an interesting and potentially useful tool in solid-state quantum devices.

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[1] *Quantum Noise in Mesoscopic Physics*, edited by Yu. V. Nazarov (Kluwer, Dordrecht, 2003).
[2] M. Field *et al.*, Phys. Rev. Lett. **70**, 1311 (1993).
[3] E. Buks, R. Schuster, M. Heiblum, D. Mahalu, and V. Umansky, Nature (London) **391**, 871 (1998).
[4] J. M. Elzerman *et al.*, Phys. Rev. B **67**, 161308 (2003).
[5] S. A. Gurvitz, Phys. Rev. B **56**, 15 215 (1997).
[6] I. L. Aleiner, N. S. Wingreen, and Y. Meir, Phys. Rev. Lett. **79**, 3740 (1997).
[7] Y. Levinson, Europhys. Lett. **39**, 299 (1997).
[8] L. Stodolsky, Phys. Rep. **320**, 51 (1999).
[9] D. V. Averin and K. K. Likharev, J. Low Temp. Phys. **62**, 345 (1986).
[10] Yu. Makhlin, G. Schön, and A. Shnirman, Phys. Rev. Lett. **85**, 4578 (2000).
[11] M. H. Devoret and R. J. Schoelkopf, Nature (London) **406**, 1039 (2000).
[12] G. Johansson, A. Käck, and G. Wendin, Phys. Rev. Lett. **88**, 046802 (2002).
[13] A. B. Zorin, Phys. Rev. Lett. **76**, 4408 (1996).
[14] A. A. Clerk, S. M. Girvin, A. K. Nguyen, and A. D. Stone, Phys. Rev. Lett. **89**, 176804 (2002).
[15] J. Delahaye *et al.*, Science **299**, 1045 (2003).
[16] O. Buisson, F. Balestro, J. P. Pekola, and F. W. J. Hekking, Phys. Rev. Lett. **90**, 238304 (2003).
[17] S. Pilgram and M. Büttiker, Phys. Rev. Lett. **89**, 200401 (2002).
[18] A. A. Clerk, S. M. Girvin, and A. D. Stone, Phys. Rev. B **67**, 165324 (2003).
[19] In some regimes, the particle transfer through the detector takes place in more than one step and cannot be characterized by one transmission amplitude t. In this work, we concentrate on one-step transfer, since only in this case can the detector be quantum limited.
[20] D. V. Averin, in Ref. [1].
[21] A. N. Korotkov, in Ref. [1]; Phys. Rev. B **67**, 235408 (2003).
[22] D. V. Averin and R. Fazio, JETP Lett. **78**, 1162 (2003).
[23] R. Ruskov and A. N. Korotkov, Phys. Rev. B **67**, 241305 (2003).
[24] A. N. Korotkov and D. V. Averin, Phys. Rev. B **64**, 165310 (2001); L. N. Bulaevskii and G. Ortiz, Phys. Rev. Lett. **90**, 040401 (2003); E. Il'ichev *et al.*, Phys. Rev. Lett. **91**, 097906 (2003); A. Shnirman, D. Mozyrsky, and I. Martin, cond-mat/0311325.