

Nonsymmetrized Correlations in Quantum Noninvasive Measurements

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A long-standing problem in quantum mesoscopic physics is which operator order corresponds to noise expressions like $\langle I(-\omega)I(\omega) \rangle$, where $I(\omega)$ is the measured current at frequency ω . Symmetrized order describes a classical measurement while nonsymmetrized order corresponds to a quantum detector, e.g., one sensitive to either emission or absorption of photons. We show that both order schemes can be embedded in quantum weak-measurement theory taking into account measurements with memory, characterized by a memory function which is independent of a particular experimental detection scheme. We discuss the resulting quasiprobabilities for different detector temperatures and how their negativity can be tested on the level of second-order correlation functions already. Experimentally, this negativity can be related to the squeezing of the many-body state of the transported electrons in an ac-driven tunnel junction.

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Although quantum measurement theory has been based on the projection postulate [1], nowadays it includes generalized schemes based on auxiliary detectors, described mathematically by positive operator-valued measures (POVM) [2]. To specify a POVM requires arguments based on physical considerations such as detector efficiency, or the assumption of thermal equilibrium. A real physical interaction generally leads to backaction on the system to be measured, which makes the interpretation of measurements difficult. Hence, all detection schemes are in general invasive as the measured system is perturbed. The disturbance is strongest for projective measurements, as the information in the measurement basis is completely erased. In contrast, other POVM schemes can be much less disturbing, as is often the case in experiments [3–5].

To avoid invasiveness, Aharonov, Albert, and Vaidman [6] studied the limit of a *weak measurement*, in which the system is coupled so weakly to the detector that it remains almost untouched. The price to pay is a large detection noise, which is however completely independent of the system. The gain is that other measurements on a non-compatible observable can be performed. After the subtraction of the detector noise, the statistics of the measurements has a well-defined limit for vanishing coupling, which for incompatible observables turns out to be described by a quasiprobability and not a real probability distribution [7,8].

The most common weak-measurement theories assume that the system-detector interaction is instantaneous [9–14]. Such a Markovian measurement scheme is relevant for many experiments [3] and corresponds to the symmetrized order of operators: $\langle I(-\omega)I(\omega) \rangle \rightarrow \langle \hat{I}(-\omega)\hat{I}(\omega) + \hat{I}(\omega)\hat{I}(-\omega) \rangle/2$, where quantum expectation values are defined as $\langle \hat{X} \rangle = \text{Tr} \hat{X} \hat{\rho}$ for an initial state $\hat{\rho}$. Here,

$\hat{I}(\omega) = \int dt \hat{I}(t) e^{i\omega t}$ is the Fourier transform of the time-dependent current $\hat{I}(t)$ in the Heisenberg picture. However, certain experiments are well described by nonsymmetrized correlators like $\langle \hat{I}(-\omega)\hat{I}(\omega) \rangle$. For $\omega \geq 0$ [4,15–20] this corresponds to noise emitted by the system (emission noise) which is measured, e.g., by an absorptive photodetector. These experiments clearly lie beyond the scope of Markovian weak-measurement theory.

In this Letter, we formulate a general theory of weak detection which allows for the description of nonsymmetrized correlators. We show that emission noise ($\omega \geq 0$) and absorption noise ($\omega < 0$) appear naturally if one takes into account measurements with *memory*. In fact, non-Markovian weak measurements follow just from a few natural assumptions imposed on the POVM in the limit of weak coupling. The results are independent of a particular experimental realization and depend only on a single memory function. By further requiring that no information transfer occurs in thermal equilibrium the scheme is fixed uniquely and contains only the detector temperature as a parameter. Varying the detector temperature interpolates between emission and absorption measurements. As our scheme is independent of other properties of the detector, it applies to circuit QED, mesoscopic current measurements and quantum optical systems equally well. Interestingly, applied to a simple harmonic oscillator, the correlation functions in this scheme are consistent with the Glauber-Sudarshan P function [21] known from quantum optics for absorption detectors. Contrary to the instantaneous measurements, the non-Markovian scheme can violate weak positivity [19]. To test it, we propose a measurement of photon-assisted current fluctuations, which are shown to violate a

Cauchy-Schwarz type inequality, proving the negative quasiprobability of the statistics after deconvolution of the detection noise. Identifying the finite-frequency current operators with quadratures in analogy to quantum optics, we show that the thus created nonequilibrium state of the current is squeezed and therefore has essentially nonclassical correlations.

We start by developing a general framework of weak quantum measurement based upon the POVM formalism including non-Markovian features. We consider a set of n independent detectors continuously recording n time-dependent signals $a_j(t)$ for $j = 1, \dots, n$. Each detector is related to an observable \hat{A}_j . For example, n ammeters are inserted in a complex circuit: $a_j(t)$ is the recorded current in the branch j and $\hat{I}_j(t)$ the current operator in that branch. Note that in general $\hat{A}_j(t)$ and $\hat{A}_k(t')$ do not commute even if $j \neq k$ since \hat{A}_j and \hat{A}_k may not commute with the Hamiltonian. We want to relate classical correlators of measured quantities like $\langle a_1(t_1) \cdots a_n(t_n) \rangle$ to their equivalent for weak quantum measurements $\langle \cdots \rangle_w$. These should involve linear correlators of the \hat{A}_j , which can be taken at different times to allow for memory effects of the detectors, while preserving causality. The requirements of linearity and causality are fulfilled by replacing $a_j(t)$ in the correlator by a superoperator $\int dt' \check{A}_j^{t-t'}(t')$ and perform time order, i.e.,

$$\langle a_1(t_1) \cdots a_n(t_n) \rangle_w = \text{Tr} \int d^n t' \mathcal{T} [\check{A}_n^{t_n-t'_n}(t'_n) \cdots \check{A}_1^{t_1-t'_1}(t'_1)] \hat{\rho}. \quad (1)$$

Here, \mathcal{T} denotes time order with respect to the arguments in brackets, $\hat{\rho}$ is the density matrix, and \check{A} are superoperators defined as

$$\check{A}_j^{t-t'}(t') = g_j(t-t') \check{A}_j^c(t') + f_j(t-t') \check{A}_j^q(t')/2. \quad (2)$$

The superoperators $\check{A}_j^{c/q}$ [22] act on any operator \hat{X} like an anticommutator or commutator: $\check{A}_j^c \hat{X} = \{\hat{A}_j, \hat{X}\}/2$ and $\check{A}_j^q \hat{X} = [\hat{A}_j, \hat{X}]/i\hbar$. In the above expressions we assumed for simplicity that the detectors are in a stationary state so that only time differences $t_j - t'_j$ matter.

We will also assume that the average of single measurements coincides with the usual average for projective measurements, i.e., $\langle a_j(t) \rangle_w = \langle \hat{A}_j(t) \rangle$. This implies $g_j(t-t') = \delta(t-t')$. Other choices of g simply mimic the effect of classical frequency filters. Thus, the only freedom left is the choice of the real function f_j that multiplies \check{A}_j^q . Note that $f_j(t)$ can be nonzero for $t > 0$ without violating causality, since it is accompanied by \check{A}_j^q and only future measurements are affected. For the last measurement, future effects disappear because the leftmost \check{A}^q vanishes under the trace in Eq. (1). For simplicity, we will assume a single

$f = f_j$, independent of j . The limit $f = 0$ corresponds to the Markovian case.

Now we want to show that correlations obeying these requirements can be obtained from the general quantum measurement formalism. Based on Kraus operators \hat{K} [23], the probability distribution of the measurement results is $\rho = \langle \check{K} \rangle$ for $\check{K} \hat{X} = \hat{K} \hat{X} \hat{K}^\dagger$, where the only condition on \hat{K} is that the outcome probability is normalized regardless of the input state $\hat{\rho}$. Here we need \hat{K} to be time dependent. In general, we assume that $\hat{K}[\hat{A}, a]$ is a functional of the whole time history of observables $\hat{A}(t)$ and outcomes $a(t)$. We shall assume that the functional \hat{K} is stationary so it depends only on relative time arguments.

The essential step to satisfy Eq. (1) is to take the limit $\hat{K} \sim \hat{I}$, which corresponds to a noninvasive measurement. This can be obtained from an arbitrary initial POVM by rescaling $\hat{K}[\hat{A}, a] \rightarrow \hat{K}_\eta = C(\eta) \hat{K}[\eta \hat{A}, \eta a]$ with $\eta \rightarrow 0$, which defines $\rho_\eta = \langle \check{K}_\eta \rangle$. Here, $C(\eta)$ is a normalization factor.

The desired correlation function (1) can be derived by the following limiting procedure for an almost general POVM,

$$\langle a_1(t_1) \cdots a_n(t_n) \rangle_w = \lim_{\eta \rightarrow 0} \langle a_1(t_1) \cdots a_n(t_n) \rangle_\eta, \quad (3)$$

where the average on the right-hand side is with respect to ρ_η . We assume the absence of internal correlations between different detectors, namely, $\hat{K}[\hat{A}, a] = \mathcal{T} \prod_j \hat{K}[\hat{A}_j, a_j]$, where \mathcal{T} applies to the time arguments of \hat{A} .

Expanding $\hat{K}[\hat{A}, a]/k[a] = 1 + \int dt' F[a, t'] \hat{A}(t') + \mathcal{O}(\hat{A}^2)$, we find, up to $\mathcal{O}(\hat{A}^2)$,

$$\check{K}/|k[a]|^2 \simeq 1 + \int dt' (2\text{Re}F\check{A}^c(t') - \hbar\text{Im}F\check{A}^q(t')). \quad (4)$$

Here, $|k[a]|^2$ is a functional probability of time-resolved outcomes independent of the properties of the system which represents the detection noise. As we want the measurement to be noninvasive to lowest order, we impose the condition that $\int F[a, t'] |k[a]|^2 Da$ vanishes; Da is the functional measure. Our conditions imply that $\int 2a(t) \text{Re}F[a, t'] |k[a]|^2 Da = \delta(t-t')$, and we get $f(t-t') = -\int 2a(t) \hbar \text{Im}F[a, t'] |k[a]|^2 Da$. Thus, the most general weak Kraus operator takes the form given in Eq. (4), which is our main result. A particular Gaussian example of a POVM realizing this scheme is presented in the Supplemental Material A [24]. We emphasize that our measurement scheme is not limited to any particular model of a detector, rather it captures generic properties of a general weakly invasive detector, whose property is encoded in the choice of the real function $f(t)$.

To discuss the consequences of different forms of f , we now calculate the noise spectral density,

$$S_{ab}(\omega) = \int dt e^{i\omega t} \langle a(t)b(0) \rangle_w. \quad (5)$$

An important special case is a system in a thermal equilibrium state, $\hat{\rho} \sim \exp(-\hat{H}/k_B T)$. We further assume that the averages of \hat{A} and \hat{B} vanish. If the detector temperature T_d is equal to T and in the absence of other nonequilibrium effects (like a bias voltage, or special initial conditions), we expect that no information transfer from the system to the detector occurs, i.e., that $S_{ab}(\omega) = 0$. This requirement leads to a necessary condition on the form of f (see Supplemental Material B [24]): $f(\omega) = i\hbar(2n_B(\omega) - 1) = i\hbar \coth(\hbar\omega/2k_B T_d)$, where $n_B(\omega)$ is the Bose distribution at temperature T_d . Equivalently, $f(t) = |k_B T_d| \coth(\pi t k_B T_d / \hbar)$ (at zero temperature $f(t) = \hbar/\pi t$). We use the name *equilibrium order* for this special choice of f . The zero temperature case has been also called *time normal* [25]. It is relevant for experimental situations like in [4] and consistent with the quantum tape [19] or photodetection model [20] if the temperature of the tape (or the photons) is T_d .

The necessary form of f is also sufficient. Indeed, the property $S_{ab}(\omega) = 0$ follows from the fluctuation-dissipation theorem [26] ($\langle \langle \dots \rangle \rangle_T$ denotes the equilibrium average)

$$\int dt e^{i\omega t} \langle \hat{A}(t) \hat{B}(0) \rangle_T = \int dt e^{i\omega t + \hbar\omega/k_B T} \langle \hat{B}(0) \hat{A}(t) \rangle_T, \quad (6)$$

because for an arbitrary stationary state we get

$$S_{ab}(\omega) = \int dt e^{i\omega t} \langle e^{\hbar\omega/2k_B T_d} \hat{B}(0) \hat{A}(t) - e^{-\hbar\omega/2k_B T_d} \hat{A}(t) \hat{B}(0) \rangle / \sinh(\hbar\omega/2k_B T_d). \quad (7)$$

For zero detector temperature, this reduces to

$$\int e^{i\omega t} dt \langle \theta(-\omega) \hat{A}(t) \hat{B}(0) + \theta(\omega) \hat{B}(0) \hat{A}(t) \rangle, \quad (8)$$

which corresponds to emission noise for $\hat{A} = \hat{B}$ and $\omega > 0$ [18]. Note that in that case (7) and (8) are even functions of ω but the operators do not appear in symmetrized form. Thus, for $T_d \neq T$, $S_{ab}(\omega)$ is in general not equal to zero and contains information about the system. It is interesting to note that reversing the sign of f transforms $S_{ab}(\omega)$ into absorption noise for $\omega \leq 0$. Hence, measuring absorption noise requires a detector formally described by a negative temperature T_d in f and Eq. (7).

It is interesting to note that for this special choice of f the higher-order fluctuations also vanish if $\hat{\rho} \propto \exp(-\hat{H}/k_B T)$ and $T = T_d$. We can write the Fourier transform of (1) as

$$\int d^n t e^{i \sum_k \omega_k t_k} \text{Tr} \mathcal{T} \prod_k \sum_{\pm} \frac{e^{\pm \hbar\omega_k/2k_B T} \check{A}_k^{\pm}(t_k)}{2 \sinh(\hbar\omega_k/2k_B T)} \hat{\rho}, \quad (9)$$

with $\check{A}^+ \hat{X} = \hat{A} \hat{X}$ and $\check{A}^- \hat{X} = \hat{X} \hat{A}$. Now, we can split $\hat{\rho} = \hat{\rho}^{1/2} \hat{\rho}^{1/2}$, expand the above expression as a sum of operator products and move one factor $\hat{\rho}^{1/2}$ leftwards and the other rightwards so that they meet again at the trace sign, which gives (1) in the form

$$\int d^n t e^{i \sum_k \omega_k t_k} \text{Tr} \hat{\rho} \mathcal{T} \prod_k \sum_{\pm} \frac{e^{\pm \hbar\omega_k/2k_B T} \check{A}_k^{\pm}(t_k \mp i\hbar/2k_B T)}{2 \sinh(\hbar\omega_k/2k_B T)}. \quad (10)$$

Shifting $t \rightarrow t \pm i\hbar/2k_B T$ and using $\text{Tr} \check{A}^q \dots = 0$ leads to

$$\int d^n t e^{i \sum_k \omega_k t_k} \text{Tr} \hat{\rho} \mathcal{T} \prod_k \check{A}_k^q(t_k) / 2i \sinh(\hbar\omega_k/2k_B T) = 0. \quad (11)$$

The vanishing of all correlations means that at any temperature $T = T_d$ in equilibrium no information is transferred, not even in higher-order correlators. This is a property of the memory function f in the idealized limit of our weak detection scheme. In the photoabsorption scheme at zero temperature, these properties are intuitively clear [20], because no photons are emitted.

Subtracting the (large) detection noise [by deconvoluting the $|k[a]|^2$ term in Eq. (4)], the correlations defined in Eq. (1) can be described by a quasiprobability. Surprisingly, the equilibrium order differs qualitatively from the symmetrized one, when one considers weak positivity, i.e., second-order correlations can be obtained from a positive probability [8,27]. The symmetrized correlation matrix $C_{ab} = \langle ab \rangle = \langle \hat{A} \hat{B} + \hat{B} \hat{A} \rangle / 2$ is positive definite, and the Gaussian probability distribution $\propto \exp(-\sum_{ab} C_{ab}^{-1} ab / 2)$ reproduces all first- and second-order symmetrized quantum correlations. This is not the case for equilibrium order. If $T_d > T$ the autocorrelation is negative meaning that also the underlying quasiprobability is negative. For example, at $T = 0$ and $\omega > 0$ Eq. (7) gives $\langle a(-\omega) a(\omega) \rangle_w = -\langle \hat{A}(\omega) \hat{A}(-\omega) \rangle e^{-\hbar\omega/2k_B T_d} / 2 \sinh(\hbar\omega/2k_B T_d)$. In a stationary situation and for $T_d = 0$, the weak positivity holds, because the correlation matrix (8) is positive definite. However, it can be violated in nonstationary situations. This can be demonstrated using a two level system with the Hamiltonian $\hat{H} = \hbar\Omega \hat{\sigma}_z / 2$, with observables $\hat{A} = \hat{B} = \hat{\sigma}_x + \hat{\sigma}_z$ and the initial state $\hat{\rho}(0) = (\hat{1} + \hat{\sigma}_y) / 2$. By direct calculation (see Supplemental Material C [24]) we find $\langle a(0) b(0) \rangle_w = -(2/\pi) \ln \Omega t_\infty$, where t_∞ is a cutoff set by intrinsic decoherence or detector backaction. Since the observables measured by the two detectors are the same, $\langle a(0) b(0) \rangle_w = \langle a^2(0) \rangle_w$ and weak positivity is obviously violated.

There is an interesting connection between equilibrium order and the Glauber-Sudarshan P function [21]. Let us take the harmonic oscillator $\hat{H} = \hbar\Omega(\hat{p}^2 + \hat{x}^2)/2$, with $[\hat{x}, \hat{p}] = i$ and $\Omega > 0$, and consider correlations with respect to the quasiprobability (1) with antisymmetric $f(t)$. In this case the time order is irrelevant as shown in Supplemental Material D [24]. Let us define $\check{A}(t) = \check{A}^c(t) + \int dt' f(t-t') \check{A}^q(t') / 2$. The evolution of $\check{x}(t)$ and $\check{p}(t)$ is just classical,

$$\begin{aligned}\check{x}(t) &= \check{x} \cos(\Omega t) + \check{p} \sin(\Omega t), \\ \check{p}(t) &= \check{p} \cos(\Omega t) - \check{x} \sin(\Omega t),\end{aligned}\quad (12)$$

while for $t = 0$ [since $f(t)$ is antisymmetric and real] we have $\check{x} = \check{x}^c + i\check{p}^q f(\Omega)/2$, $\check{p} = \check{p}^c - i\check{x}^q f(\Omega)/2$. Equation (12) means that all correlators in this case undergo the classical time evolution. We now define the ladder operators through $\hat{a}^\dagger = (\hat{x} + i\hat{p})/\sqrt{2}$ with $[\hat{a}, \hat{a}^\dagger] = \hat{1}$. This leads to $\check{a} = \check{a}^c - f(\Omega)\check{a}^q/2$ and $\check{a}^\dagger = \check{a}^{\dagger c} + f(\Omega)\check{a}^{\dagger q}/2$. In the zero-temperature case, $f(\Omega) = i\hbar$ (a perfect photodetector), and defining $\alpha = (x + ip)/\sqrt{2}$ we get the single-time quasiprobabilistic average $\langle \alpha^n \alpha^{*k} \rangle = \text{Tr} \hat{\rho} \hat{a}^{\dagger k} \hat{a}^n$. On the other hand, this is a property of the Glauber-Sudarshan function $P(\alpha)$, defined by $\hat{\rho} = \int d^2\alpha P(\alpha) |\alpha\rangle\langle\alpha|$ for normalized coherent states $\hat{a}|\alpha\rangle = \alpha|\alpha\rangle$, $\langle\alpha|\alpha\rangle = 1$ [21]. Since $\langle \alpha^n \alpha^{*k} \rangle_P = \int d^2\alpha \alpha^n \alpha^{*k} P(\alpha) = \text{Tr} \hat{\rho} \hat{a}^{\dagger k} \hat{a}^n$, we find that the quasiprobability for a zero-temperature detector is identical to $P(\alpha)$. It is interesting to note that reversing the sign of f leads to the Husimi-Kano Q function instead of P [2], while $f = 0$ gives the Wigner function [7,28].

The fact that we obtain the P function shows the deep connection between the non-Markovian weak measurement formalism and the quantum-optical detector theory. One of the interesting consequences is that zero-temperature equilibrium order is consistent with photo-absorptive detection schemes, in which the P -function appears naturally [2]. It is also interesting to draw a link between the violation of weak positivity in equilibrium order and the properties of squeezed states. The ground state of a harmonic oscillator fulfills $\langle \hat{x}^2 \rangle = 1/2$, which corresponds to $\langle x^2 \rangle_P = 0$. A squeezed state can be such that $\langle \hat{x}^2 \rangle < 1/2$, still minimizing the Heisenberg uncertainty principle. This translates into a negative variance of the position described by the (quasiprobability) P function, i.e., $\langle x^2 \rangle_P < 0$ [29] and is therefore equivalent to a violation of weak positivity.

Let us now consider how our results apply to the case of current fluctuations in mesoscopic conductors. The quantum description of the noise in the junction, $S_I(\omega) = \int dt e^{i\omega t} \langle \delta I(t) \delta I(0) \rangle$, where $\delta I(t) = I(t) - \langle I(t) \rangle$, will depend on the choice of f in (1). For $f = 0$, we get symmetrized noise $S_I^s = G\hbar \sum_{\pm} w(\omega \pm eV/\hbar, T)/2$, where G is the conductance, V is the constant bias voltage, and $w(\alpha, T) = \alpha \coth(\hbar\alpha/k_B T)$ [19]. For f given by equilibrium order with an arbitrary T_d , we obtain $S_I = S_I^s - G\hbar w(\omega, T_d)$. Hence, the detection schemes differ by a term that is independent of the voltage and the temperature of the system, making it impossible to detect nonclassicality in this scheme.

An experimentally feasible test of squeezing and violation of weak positivity is possible using a coherent conductor (e.g., a tunnel junction for the sake of simplicity) subject to an ac voltage bias $V(t) = V_{ac} \cos \Omega t$ [30]. Consider the classical inequality

$$|\delta I(\omega) - \delta I(-\omega)|^2 \geq 0 \Rightarrow \langle |\delta I(\omega)|^2 \rangle \geq \text{Re} \langle \delta I^2(\omega) \rangle. \quad (13)$$

For symmetrized order one gets [31]

$$\begin{aligned}\langle \{\delta \hat{I}(\omega), \delta \hat{I}(\omega')\} \rangle / 2 &= 2\pi\hbar G \sum_m \delta(\omega + \omega' - 2m\Omega), \\ \sum_n J_n(eV_{ac}/\hbar\Omega) J_{n-2m}(eV_{ac}/\hbar\Omega) w(\omega - n\Omega),\end{aligned}\quad (14)$$

where J_n are the Bessel functions. In the case of equilibrium order at $T_d = 0$ one only has to subtract $2\pi\hbar G |\omega| \delta(\omega + \omega')$ from the above result. As shown in Fig. 1, the classical inequality is violated for $\omega = \Omega$ in a certain range of $eV_{ac}/\hbar\Omega$, but only in equilibrium order. This can be reinterpreted in terms of the existence of squeezing in the quantum shot noise: consider the two quadratures associated with the finite-frequency current operator: $\hat{A} = i[\delta \hat{I}(\omega) - \delta \hat{I}(-\omega)]/2$ and $\hat{B} = [\delta \hat{I}(\omega) + \delta \hat{I}(-\omega)]/2$. Using $\langle [\hat{I}(\omega), \hat{I}(-\omega)] \rangle = 2t_0 G \hbar \omega$, we find [32] $\langle [\hat{A}, \hat{B}] \rangle = it_0 G \hbar \omega$ (with the total detection time t_0). Thus the squeezing condition [2],

$$\langle \hat{A}^2 \rangle < |\langle [\hat{A}, \hat{B}] \rangle| / 2, \quad (15)$$

is related to the violation of weak positivity, $\langle A^2 \rangle_w < 0$ in equilibrium order with $T_d = 0$ and allows the violation of Eq. (13). Hence, according to Fig. 1, quantum shot noise with ac driving creates current states, which resemble squeezed light for a certain range of the ac voltage.

In conclusion, we have presented a theory of a generic weak-measurement scheme that includes emission noise. It requires a non-Markovian POVM with a specially chosen memory function f , which has no analog in the Markovian picture. The scheme is consistent with the absence of information flow between system and detector in equilibrium at a given temperature. Hence, any detection requires a nonequilibrium situation. Another direct consequence is that even the simple Markovian detection process must involve a nonequilibrium detector state. Finally,

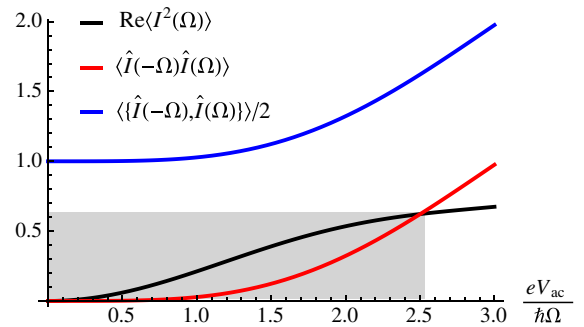


FIG. 1 (color online). Quantum correlation functions (in units of $2\pi G\hbar\Omega t_0$) for a tunnel junction at zero temperature $T = 0$. The emission noise $\langle \hat{I}(-\Omega)\hat{I}(\Omega) \rangle$ (red line) violates the classical inequality (13) for a certain range of eV_{ac} (shaded region). This violation is equivalent to the squeezing condition (15) for the symmetrized noise $\langle \{\hat{I}(-\Omega), \hat{I}(\Omega)\} \rangle / 2$ (blue line).

nonsymmetrized order leads to a violation of weak positivity, which can be tested experimentally by violation of suitable inequalities, equivalent to the squeezing condition in some cases.

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Note added.—Our prediction about ac-driven squeezing has been recently confirmed experimentally [33].

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