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Nonzero temperature effects on antibunched photons emitted by a quantum point contact out of equilibrium

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Electrical current fluctuations in a single-channel quantum point contact can produce photons (at frequency $\omega$ close to the applied voltage $V\times e/\hbar$) which inherit the sub-Poissonian statistics of the electrons. We extend the existing zero-temperature theory of the photostatistics to nonzero temperature $T$. The Fano factor $F$ (the ratio of the variance and the average photocount) is $<1$ for $T<T_c$ (antibunched photons) and $>1$ for $T>T_c$ (bunched photons). The crossover temperature $T_c = \Delta\omega \times \hbar/k_B$ is set by the bandwidth $\Delta\omega$ of the detector, even if $\hbar\Delta\omega \ll eV$. This implies that narrow-band detection of photon antibunching is hindered by thermal fluctuations even in the low-temperature regime where thermal electron noise is negligible relative to shot noise.

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I. INTRODUCTION

It is a celebrated result of Glauber that classical fluctuations of the electrical current $I(t)$ produce photons with Poisson statistics. The variance $Var n$ of the number $n$ of photons detected in a time $t_{\text{det}}$ is then equal to the mean $\langle n \rangle$. The photons produced by a classical current thus behave as independent classical particles. Quantum fluctuations of the current (with $I(t)$ and $I(t')$ noncommuting operators) change the photostatistics. The bosonic nature of the photons would naturally lead to photon bunching, with $Var n > \langle n \rangle$. Photon antibunching, with $Var n < \langle n \rangle$, is also possible, if the photons somehow inherit the sub-Poissonian statistics of the electrons. One then speaks of nonclassical light. While nonclassical light emitted by a quantum mechanical current has not yet been observed, the theory is well developed and there is an active experimental search. Nonclassical light emitted by a quantum conductor such as a quantum point contact would occur in a continuous range of GHz frequencies, in contrast to the discrete frequencies produced by electronic transitions in quantum dots or quantum wells. Various methods of measuring the photostatistics have been developed, such as detection by means of photoassisted tunneling, or by means of the Hanbury-Brown-Twiss effect.

The theoretical prediction is that photons emitted by a single-channel quantum point contact should have a Fano factor $F=Var n/\langle n \rangle$ smaller than unity at zero temperature, for frequencies $\omega$ close to the applied voltage $V\times e/\hbar$. More specifically,

$$ F = 1 - \frac{2}{3} \gamma_0^2 \Delta\omega \tau (1 - \tau), \quad (1.1) $$

for photodetection with efficiency $\gamma_0$ in the frequency interval $\langle eV/h - \Delta\omega, eV/h \rangle$. The transmission probability $\tau$ through the quantum point contact is assumed to be energy independent on the scale of $eV$. Equation (1.1) is derived in the limit of weak coupling $\gamma_0^2 \Delta\omega \ll 1$ of electrons to photons, so that the deviations from Poisson statistics remain small. It is also assumed that the photons can be detected individually, see Ref. 6 for an alternative detection scheme.

It is the purpose of the present paper to extend the theory of Ref. 5 to nonzero temperatures, in order to identify the conditions on the temperature needed to observe the photon antibunching. Clearly, photon bunching should take over when the electrical shot noise drops below the thermal noise, which happens when $k_B T_b$ becomes larger than $eV$. While $k_B T_b < eV$ is the condition for photon antibunching in the case of wide-band detection, a more stringent condition $k_B T_b < \hbar\Delta\omega$ holds for narrow-band detection.

More precisely, we obtain a crossover temperature $T_c = \hbar\Delta\omega/4k_B$ at which $F=1$ for $\Delta\omega \ll eV$. In this low-temperature regime shot noise still dominates over thermal noise, yet the photon antibunching is lost. One qualitative way to understand this is, is to compare the coherence time $t_{\text{coh}} = 1/\Delta\omega$ of the detected radiation with the coherence time $t_r = \hbar/k_B T_r$ of thermally excited electron-hole pairs. For $t_{\text{coh}} > t_r$ the detected photons result from many uncorrelated electron-hole recombination events, and the one-to-one relationship between electron and photon statistics is lost.

In the next section, we give the nonzero temperature generalization of the theory of Ref. 5, and then in Sec. III, we specialize to the shot-noise regime $k_B T_b \ll eV$. General results in both the shot noise and thermal noise regimes are presented in Sec. IV. Technical details are summarized in Appendixes A and B.

II. GENERATING FUNCTION AT NONZERO TEMPERATURE

We seek the non-zero-temperature generalization of the formula

$$ F(\xi) = \sqrt{\prod_{m=1}^{N} \text{Det}(1 + T_m \left[e^\xi e^{\xi'} - 1\right])}, \quad (2.1) $$

for the factorial-moment generating function $F(\xi)$ of the photocount. We first introduce the notation and then present the required generalization.

The photons are produced by time-dependent current fluctuations in a quantum point contact, characterized by trans-
mission eigenvalues $T_1, T_2, \ldots, T_N$, with $N$ the number of propagating electronic modes (counting both orbital and spin degrees of freedom). The current flows between two reservoirs, with Fermi functions

$$f_L(e) = (1 + \exp[(e - eV - E_p)/k_B T])^{-1},$$

$$f_R(e) = (1 + \exp[(e - E_p)/k_B T])^{-1}.$$  \hfill (2.2)

The current fluctuations can be due to thermal noise (at temperature $T$) or due to shot noise (at a voltage $V$ applied over the point contact). We take the transmission eigenvalues $T_{an}$ as energy independent in the range $\max(eV, k_BT)$ near the Fermi energy $E_F$.

The photons are detected during a time $t_{det}$ in a narrow frequency interval $\Delta \omega$ around frequency $\Omega$, as determined by the detection efficiency $\gamma(\omega)$. Antibunching of the photons requires that $\Omega$ is tuned to the applied voltage, $\Omega = eV/h$. (In the following we set $\hbar$ and $e$ equal to unity.)

The average $\langle \cdots \rangle$ in Eq. (2.1) indicates a Gaussian integration over the complex numbers $\xi_{p}$. The photon is detected during a time $t_{det}$ in a narrow frequency interval $\Delta \omega$ around frequency $\Omega$, as determined by the detection efficiency $\gamma(\omega)$. Antibunching of the photons requires that $\Omega$ is tuned to the applied voltage, $\Omega = eV/h$. (In the following we set $\hbar$ and $e$ equal to unity.)

The Fermi function $f_L(e)$ in the left electronic reservoir is contained in the diagonal matrix $f_L$, with elements $f_{L,pp'} = \delta_{pp'}f_L(e_p)$. $e_{p} = e_{p'} \times 2\pi/\nu(t_{det})$. Similarly, the Fermi function $f_R(e)$ in the right reservoir is contained in the diagonal matrix $f_R$.

Following the steps in Appendix A, the expression (2.6) can be reduced to the more compact form

$$F(\xi) = \prod_{n=1}^{N} \text{Det} \left( \begin{array}{c} 1 + T_{mL}(e^{z^2} - 1) \sqrt{T_{mL}(1 - T_{mL})f_L(e^{z^2} - \tilde{e}_{n}^2)} \left( 1 + T_{mL}(e^{z^2} - \tilde{e}_{n}^2) \right) \\ 1 + T_{mL}(e^{z^2} - \tilde{e}_{n}^2) \end{array} \right).$$ \hfill (2.6)

with the definitions $\tilde{e} = 1 - e_L, \tilde{e}_R = 1 - e_R$. $\tilde{M} = e^2 - e^{-Z^2}$. The zero-temperature limit [Eq. (2.1)] follows from Eq. (2.7) by setting $f_{L} = 1, f_{R} = 0$ in the energy interval $E_F < e < E_F + V$. (There are no current fluctuations outside of this energy interval for $T = 0$.)

### III. SHOT NOISE REGIME

The result [Eq. (2.7)] holds for any temperature, provided that the energy dependence of the transmission eigenvalues may be neglected. In particular, it describes both thermal noise and shot noise. A simpler formula is obtained in the shot noise regime $k_BT \ll V$. Thermal noise can then be neglected and only the finite temperature effects on the shot noise are retained. We assume $\Delta \omega \ll \Omega = V$, so even if $k_BT \ll V$, the relative magnitude of $\Delta \omega$ and $k_BT$ is still arbitrary.

#### A. Generating function

The first simplification in this regime is that we may set $f_{L} e^{-\tilde{f}_{L} L} = 0$, since $f_{L}(e) f_{L}(e') \to 0$ for $e' \ll e$. Equation (2.7) reduces to

$$F(\xi) = \prod_{n=1}^{N} \text{Det} \left( e^{-Z^2} + T_{mL}M_{L} \right),$$ \hfill (3.1)

where we have multiplied by $\text{Det} e^{-Z^2} = 1$.

The second simplification is that we can ignore energies separated by $pV$ with $p \geq 2$, because $V$ is the largest energy scale in the problem. Since $Z^2$ and $Z'^2$ connect energies separated by $p\Omega = pV$, we may set $Z^2 = Z'^2 \to 0$ for $p \geq 2$. From Eq. (3.1) we arrive at

$$F(\xi) = \prod_{m=1}^{N} \text{Det} \left( 1 - Z^2 + T_{mL}(Z + Z^2) \right).$$ \hfill (3.2)

Following the steps in Appendix B, the determinant may be rewritten in the more convenient form (bilinear in $Z, Z'$),
\[ F(\xi) = \left\{ \prod_{m=1}^{N} \text{Det}(1 + T_m(1 - T_m)f_z Z_{kR} Z^\dagger) \right\}. \quad (3.3) \]

**B. Moment expansion**

The generating function (3.3) is of the form \( F(\xi) = \Pi_m \text{Det}(1 + X_m) \) with \( X_m \) of order \( \xi \). An expansion in powers of \( \xi \) can be obtained by starting from the identity

\[ \prod_m \text{Det}(1 + X_m) = \exp \left[ \sum_m \text{Tr} \ln(1 + X_m) \right]. \quad (3.4) \]

and expanding in turn, the logarithm and the exponential. Up to second order in \( \xi \) we have the expansion

\[ F(\xi) = 1 + \left( \sum_m \text{Tr} X_m \right) - \frac{1}{2} \left( \sum_m \text{Tr} X_m^2 \right) + \frac{1}{2} \left( \left( \sum_m \text{Tr} X_m \right)^2 \right) + O(\xi^3), \quad (3.5) \]

from which we can extract the first two factorial moments,

\[ F(\xi) = 1 + \xi \langle n \rangle + \frac{1}{2} \xi^2 (\langle n^2 \rangle - \langle n \rangle^2) + O(\xi^3). \quad (3.6) \]

We perform the Gaussian averages and obtain the average photocount \( \langle n \rangle \) and the variance \( \text{Var} n = \langle n^2 \rangle - \langle n \rangle^2 \) in the shot noise regime,

\[ \langle n \rangle = \frac{\text{det} S_1}{2\pi} \int d\omega \gamma(\omega) \int d\epsilon f_L(\epsilon + \omega)f_R(\epsilon), \quad (3.7) \]

\[ \text{Var} n = \langle n \rangle + \frac{\text{det} S_1}{2\pi} \int d\omega \left[ \gamma(\omega) \int d\epsilon f_L(\epsilon + \omega)f_R(\epsilon) \right]^2 - \frac{\text{det} S_1}{2\pi} \int d\omega \left[ f_L(\epsilon) \int d\omega \gamma(\omega)f_R(\epsilon - \omega) \right]^2 - \frac{\text{det} S_2}{2\pi} \int d\omega \left[ f_R(\epsilon) \int d\omega \gamma(\omega)f_L(\epsilon + \omega) \right]^2. \quad (3.8) \]

We have defined

\[ S_p = \sum_m \left[ T_m(1 - T_m) \right]^p. \quad (3.9) \]

Since the two reservoirs are at the same temperature, we can write \( f_L(\epsilon) = f(\epsilon - V - E) \) and \( f_R = f(E - \epsilon) \) in terms of a single Fermi function

\[ f(\epsilon) = (1 + e^{\epsilon/k_BT})^{-1}. \quad (3.10) \]

We abbreviate \( \Gamma(\epsilon, \omega) = \gamma(\omega)f(\epsilon)f(\omega - V - E) \) and can then write Eqs. (3.7) and (3.8) in the compact form

\[ \langle n \rangle = \frac{\text{det} S_1}{2\pi} \int d\omega \int d\epsilon \Gamma(\epsilon, \omega), \quad (3.11) \]

\[ \text{Var} n = \langle n \rangle + \frac{\text{det} S_1}{2\pi} \int d\omega \int d\epsilon \Gamma(\epsilon, \omega) \times \left[ S_1^2 \int d\epsilon' \Gamma(\epsilon, \omega) - 2S_2 \int d\epsilon' \Gamma(\epsilon, \omega') \right]. \quad (3.12) \]

The difference \( \text{Var} n - \langle n \rangle \) contains a positive term \( \propto S_1^2 \) and a negative term \( \propto S_2 \). The sign of this difference determines whether there is bunching or antibunching of the detected photons.

**C. Crossover from antibunching to bunching**

To investigate the crossover from antibunching to bunching with increasing temperature, we take a block-shaped response function

\[ \gamma(\omega) = \begin{cases} \gamma_0 & \text{if } V - \Delta \omega < \omega < V \\ 0 & \text{otherwise.} \end{cases} \quad (3.13) \]

In the low-temperature regime \( k_BT \ll \Delta \omega \) the function \( \Gamma(\epsilon, \omega) \) then has a block shape as well and we recover the results

\[ \langle n \rangle = \frac{\text{det} \Delta \omega}{2\pi} \gamma_0 \Delta \omega \frac{1}{3} S_1, \quad (3.14) \]

\[ \text{Var} n - \langle n \rangle = \frac{\text{det} \Delta \omega}{2\pi} (\gamma_0 \Delta \omega)^2 \frac{1}{3} (S_1^2 - 2S_2). \quad (3.15) \]

of Ref. 5. These correspond to a Fano factor

\[ F = 1 + \frac{2}{3} \gamma_0 \Delta \omega (S_1^2 - 2S_2/S_1). \quad (3.16) \]

For a single-channel conductor \( S_2 = S_1^2 \), so there is antibunching \((F<1)\) at low temperatures.

At high temperatures \( k_BT \gg \Delta \omega \), but still in the shot-noise regime \( k_BT \ll V \), we may substitute \( \Gamma(\epsilon, \omega) \rightarrow -\gamma(\omega)k_BTdf(\epsilon)/d\epsilon \) into Eqs. (3.11) and (3.12), which gives

\[ \langle n \rangle = \frac{\text{det} \Delta \omega}{2\pi} \gamma_0 k_BT \frac{1}{3} S_1, \quad (3.17) \]

\[ \text{Var} n - \langle n \rangle = \frac{\text{det} \Delta \omega}{2\pi} (\gamma_0 k_BT)^2 S_1^2. \quad (3.18) \]

The Fano factor

\[ F = 1 + \gamma_0 k_BT S_1 \quad (3.19) \]

is now \( \gg 1 \)—hence, there is photon bunching.

The crossover temperature \( T_c \), at which \( F=1 \), can be calculated numerically from Eqs. (3.11) and (3.12). In the single-channel case, when \( S_2 = S_1^2 \), we find

\[ k_BT_c = 0.25 \Delta \omega. \quad (3.20) \]

The crossover is shown graphically in Fig. 1.
Antibunching crosses over into bunching as a result of thermal noise. The temperature 
\( T_c \) at which antibunching crosses over into bunching, so when \( F=1 \), follows the shot-noise limit [Eq. (3.20)] for narrow-band detection

\[ \Delta \omega \ll V \]

With increasing bandwidth, \( T_c \) drops below the shot noise limit, in particular for small transmission probability \( \tau \). For \( \tau=0.5 \) the shot-noise limit remains accurate even for bandwidths \( \Delta \omega \) as large as \( V/2 \).

V. CONCLUSION

In conclusion, we have investigated the effects of a non-zero temperature on the degree of antibunching of photons produced by current fluctuations in a quantum point contact. Antibunching crosses over into bunching as a result of thermal noise in the point contact, but this is not the dominant effect in the case of narrow-band detection. In that case, the finite coherence time of electron-hole pairs governs the transition from photon antibunching to photon bunching, which occurs at a temperature \( k_B T_c = \Delta \omega \) even if \( k_B T_c \ll V \) (so even if thermal noise is negligible relative to shot noise).

The optimal conditions for the observation of antibunched photons are reached for a bandwidth \( \Delta \omega = V/2 \) and a transmission probability \( \tau = 1/2 \) through a (spin-resolved) single-channel quantum point contact. In that case \( k_B T_c = V/8 \) has the largest value at any given applied voltage. The currently available detection range \( 4 \text{ kHz} < \omega < 8 \text{ GHz} \Rightarrow \Delta \omega = 4 \text{ GHz} = V/2 \), should make it possible to detect antibunching at temperatures below 1 GHz \( \approx 50 \) mK.

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APPENDIX A: DERIVATION OF THE GENERATING FUNCTION AT NONZERO TEMPERATURE

We briefly describe how the analysis of Ref. 5 can be generalized to nonzero temperatures, in order to arrive at Eq. (2.6). Referring to the equations in that paper, the first equation which changes is Eq. (5) of Ref. 5, which now reads...
\[ F(\xi) = (e^{-\alpha^* DZ a} e^{b^* DZ b} e^{b DZ^j} e^{-\alpha DZ^j a}). \]  
(\text{A1})

The four factors correspond to the four current operators that need to be taken into account: \( I_{in}^l, I_{out}^l, I_{out}^r, \) and \( I_{in}^r \).

The operator \( a^\dagger \) creates an incoming electron, while \( b^\dagger \) creates an outgoing electron. The matrix \( D \) projects on the right lead, where the current is evaluated. (Since \( D \) commutes with \( Z \), we can write \( DZ \) instead of \( DDZ \).) One can relate \( b=sa \), with \( S \) the unitary scattering matrix, so one can write the entire generating function in terms of the operators \( a \). The expectation value \( \langle \cdots \rangle \) is both an expectation value over the fermion operators \( a \), as well as the average over the Gaussian variables \( Z, Z^j \).

Following the steps of Ref. 5, we calculate the expectation value of the fermion operators by means of the identity
\[
\left\langle \prod_n e^{iA_n a_n} \right\rangle = \det(1+AB),
\]
(\text{A2})

\[
A = \left( \prod_n e^{A_n} \right) - 1, \quad B_{ij} = \langle a_i^\dagger a_j \rangle.
\]
(\text{A3})

We have \( B_{ij} = \delta_{ij} f_i \), with \( f_i \) the Fermi occupation number in channel \( i \). The matrix \( A \) is given by \( A = e^X e^Y e^X - 1 \), with \( X = -DZ \) and \( Y = S^j DZ^j S \). Notice that \( X^p = D(-Z)^p \) and \( Y^p = S^j DZ^j S \).

We now make the assumption of an energy independent scattering matrix, so \( S, S^j \) commute with \( Z, Z^j \). The determinant is invariant under a change of basis, and by working in the eigenchannel basis we can reduce \( S \) to a \( 2 \times 2 \) matrix \( S_m \) for each eigenchannel,
\[
S_m = \left( \sqrt{1-T_m} \quad \sqrt{T_m} \right) \left( \sqrt{T_m} \quad \sqrt{1-T_m} \right),
\]
(\text{A4})

with \( T_m, m = 1, 2, \ldots N \) the transmission eigenvalue. The matrix structure of \( f, D, \) and \( Z \) in this basis is
\[
f = \begin{pmatrix} f_L & 0 \\ 0 & f_R \end{pmatrix}, \quad D = \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix}, \quad Z = \begin{pmatrix} Z & 0 \\ 0 & Z \end{pmatrix}.
\]
(\text{A5})

Substitution of Eqs. (A2)–(A5) into Eq. (A1) leads after some algebraic manipulations to the result [Eq. (2.6)].

The determinant in Eq. (2.6) can be reduced by means of the folding identity
\[
\det \left( \begin{array}{cc} M_{11} & M_{12} \\ M_{21} & M_{22} \end{array} \right) = \det M_{11} \det (M_{22}-M_{21}M_{11}^{-1}M_{12}).
\]
(\text{A6})

leading to
\[
F(\xi) = \prod_{m=1}^N \left[ 1 + T_m f_L(e^{2\varepsilon}e^{-\varepsilon^*}) - T_m(1-T_m)f_R(e^{-\varepsilon^*}f_L(e^{2\varepsilon}-1)]^{-1}f_L(e^{2\varepsilon}-1) \right].
\]
(\text{A7})

We continue the reduction of the determinant, using first the identity
\[
[1 + T_m f_L(e^{2\varepsilon}e^{-\varepsilon^*})]^{-1} f_L(e^{2\varepsilon}-1) = - f_L(e^{2\varepsilon}-1)]^{1 + T_m f_L(e^{2\varepsilon}e^{-\varepsilon^*})} e^{-\varepsilon^*},
\]
(\text{A8})

then multiplying the determinant by \( \det e^{2\varepsilon} = 1 \), and finally combining the product of three determinants into a single determinant. In this way we eliminate the matrix inversion, arriving at
\[
F(\xi) = \prod_{m=1}^N \left[ \det \left(1 + T_m f_L(e^{2\varepsilon}e^{-\varepsilon^*}) - 1 \right) + T_m(1-T_m)f_R(e^{-\varepsilon^*}f_L(e^{2\varepsilon}-1) \right) \right]
\]
\[
\prod_{m=1}^N \left[ \det \left(1 + T_m f_R(e^{2\varepsilon}e^{-\varepsilon^*}) - 1 \right) + T_m(1-T_m)f_L(e^{2\varepsilon}-1) \right] \right).
\]
(\text{A9})

This is Eq. (2.7) in the main text.

**APPENDIX B: DERIVATION OF THE GENERATING FUNCTION IN THE SHOT NOISE REGIME**

Starting from the expression (3.2) for the generating function in the shot noise regime \( k_B T \ll V \), we give the steps required to arrive at the bilinear form (3.3). We group terms with \( Z \) and with \( Z^j \) in the matrices \( A_p = T_m f_L f_R e^{2\varepsilon} \) and \( B_m = T_m f_L f_R e^{2\varepsilon} \), so that Eq. (3.2) can be written as
\[
F(\xi) = \prod_{m=1}^N \left[ \det \left(1 + A_m + B_m \right) \right].
\]
(\text{B1})

Because energies separated by \( V^p \) with \( p \geq 2 \) can be discarded, we may set \( A_p \to 0, B_p \to 0 \). For any pair of matrices \( A, B \) which square to zero, one has the identity
\[
\text{Det}(1 + A + B) = \text{Det}(1 - AB). \quad (B2)
\]

This leads to

\[
F(\xi) = \left( \prod_{m=1}^{N} \text{Det}(1 + T_m Z \bar{f}_R \bar{f}_L - T_m Z \bar{f}_R \bar{f}_L Z^2 \bar{f}_R \bar{f}_L) \right)^N. \quad (B3)
\]

Eq. (3.3) follows by noting that \( \bar{f}_R \bar{f}_L \rightarrow \bar{f}_L \) for \( k_B T \ll \Omega = \Omega', \) since the Fermi function \( f_L \) in this term is evaluated at energies near \( E_F, \) where it can be replaced by unity. Similarly \( Z \bar{f}_R \bar{f}_L \rightarrow Z \bar{f}_L, \) since \( f_R \) is evaluated at energies near \( E_F + V \) where it can be replaced by unity.

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