

Shot noise and Coulomb blockade of Andreev reflection

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We derive low-energy effective action for a short coherent conductor between normal (N) and superconducting (S) reservoirs. We evaluate interaction correction δG to Andreev conductance and demonstrate a close relation between Coulomb effects and shot noise in NS systems. In the diffusive limit doubling of both shot-noise power and charge in the carriers yields $|\delta G|$ four times larger than in the normal case. Our predictions can be directly tested in future experiments.

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INTRODUCTION

It is well known that low-energy electron transport across the interface between normal metals and superconductors (NS) is provided by the mechanism of Andreev reflection.¹ This mechanism involves conversion of a subgap quasiparticle entering the superconductor from the normal metal into a Cooper pair together with simultaneous creation of a hole that goes back into the normal metal. Each such act of electron-hole reflection corresponds to transferring twice the electron charge $e^*=2e$ across the NS interface and results, e.g., in nonzero conductance of the system at subgap energies.²

Let us assume that two bulk metallic electrodes, one normal and one superconducting, are connected by an arbitrary—though sufficiently short—coherent conductor as it is schematically shown in Fig. 1. This conductor is characterized by the normal-state conductance

$$G_N = \frac{e^2}{h} 2 \sum_n T_n, \quad (1)$$

where T_n define transmissions of all conducting channels and the factor 2 accounts for spin degeneracy. Evaluating the conductance G_A of the NS structure in Fig. 1, at temperatures/voltages well below the superconducting gap Δ one finds²

$$G_A = \frac{(2e)^2}{h} \sum_n T_n, \quad (2)$$

where “Andreev transmissions” \mathcal{T}_n are related to T_n as

$$\mathcal{T}_n = T_n^2 / (2 - T_n)^2. \quad (3)$$

Comparing Eqs. (2) and (3) with the Landauer formula (1) one immediately observes that Andreev conductance G_A formally describes “normal” transport of *spinless* quasiparticles (hence, no extra factor 2 in front of the sum) with charge $e^*=2e$ across some effective coherent scatterer with channel transmissions \mathcal{T}_n , Eq. (3).

Later it was realized that this formal analogy applies not only to electron transport but also to low-frequency shot noise^{3–5} and eventually to full counting statistics (FCS).⁶ Consider, for instance, current fluctuations $\delta I(t) = I(t) - I$ around its average value $I \equiv \langle \hat{I}(t) \rangle$. In normal conductors at

$T \rightarrow 0$ and in the zero frequency limit the correlator for such fluctuations has the well known form⁷

$$\langle |\delta I|^2 \rangle = e |V| G_N \beta_N, \quad \beta_N = \frac{\sum_n T_n (1 - T_n)}{\sum_n T_n}, \quad (4)$$

where V is the average voltage across the conductor. In NS systems Andreev reflection also leads to the current shot noise at energies below the superconducting gap. In this case in the zero-energy/frequency limit and at $T \rightarrow 0$ one obtains⁴

$$\langle |\delta I|^2 \rangle = 2e |V| G_A \beta_A, \quad \beta_A = \frac{\sum_n \mathcal{T}_n (1 - \mathcal{T}_n)}{\sum_n \mathcal{T}_n}, \quad (5)$$

where \mathcal{T}_n is again defined by Eq. (3). Again, a close similarity between Eqs. (4) and (5) is obvious: The result, Eq. (5), just describes shot noise produced by carriers with effective charge $e^*=2e$ in a coherent scatterer with conductance G_A and Fano factor β_A .

In the important case of diffusive NS structures doubling of the carrier charge also implies doubling of the shot noise.⁴ In this case the sums over transmission channels in the above equations can be evaluated in a straightforward manner with the results

$$G_N = G_A, \quad \beta_N = \beta_A = 1/3, \quad (6)$$

which yield $\langle |\delta I|^2 \rangle = 2e |V| G_N / 3$ for NS structures and the two times smaller result in the normal case. This doubling of the shot noise in diffusive NS systems was indeed observed in experiments.^{8,9}

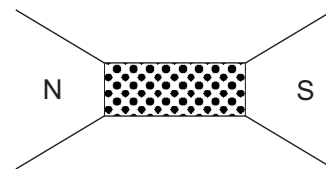


FIG. 1. Short coherent conductor connecting normal and superconducting reservoirs.

More recently another interesting observation was reported.¹⁰ The authors of this experiment investigated short metallic nanowires attached to bulk superconducting electrodes. In a number of samples superconductivity inside the wire was destroyed due to phase slippage and, hence, such samples effectively represented hybrid normal-superconducting structures, e.g., similar to those depicted in Fig. 1. Remarkably, the authors¹⁰ discovered that as long as the electrodes stayed superconducting the measured I - V curves could be well fitted by the theory of Coulomb blockade in *normal* coherent conductors¹¹ provided the electron charge e was substituted by some effective charge q_{eff} larger than e but smaller than $2e$. If, however, superconductivity in bulk electrodes was suppressed, the I - V curves of *exactly the same form* but with $q_{\text{eff}} \approx e$ were observed. Although these observations strongly indicate that Andreev reflection can be involved, no theoretical explanation of the experiments¹⁰ was offered until now.

Below we develop a theory describing an interplay between Coulomb blockade and Andreev reflection. We will explicitly evaluate the interaction correction to Andreev conductance and demonstrate its direct correspondence to the shot noise in NS structures. Hence, very different experiments in Refs. 8–10 turn out to be closely related. Both measure the same effective charge, i.e., $e^* = q_{\text{eff}}$.

EFFECTIVE ACTION

As it is shown in Fig. 1, we will consider big normal and superconducting reservoirs connected by a rather short normal bridge (conductor) with arbitrary transmission distribution T_n of its conducting modes (for each n the value T_n is the same for spin-up and spin-down electrons). Both phase and energy relaxation of electrons may occur only in the reservoirs and not inside the conductor which length is assumed to be shorter than dephasing and inelastic relaxation lengths. In contrast to¹² (where the Thouless energy ε_{Th} of the normal conductor plays an important role), here ε_{Th} of the bridge is irrelevant as it is supposed to be higher than any other energy scale in our problem. As usually, Coulomb interaction between electrons in the conductor area is accounted for by some effective capacitance C .

In order to analyze electron transport in the presence of interactions we will make use of the effective action formalism combined with the scattering matrix technique. This approach, very successful in the case of normal conductors,^{11,13–15} can be conveniently generalized to the superconducting case. In fact, the structure of the effective action remains the same also in the latter case, one should only replace normal propagators by 2×2 matrix Green's functions which account for superconductivity, as it was done, e.g., in Refs. 16–18.

Following the usual procedure we express the kernel J of the evolution operator on the Keldysh contour in terms of a path integral over the fermionic fields which can be integrated out after the standard Hubbard-Stratonovich decoupling of the interacting term. Then the kernel J takes the form

$$J = \int \mathcal{D}\varphi_1 \mathcal{D}\varphi_2 \exp(iS[\varphi]), \quad (7)$$

where $\varphi_{1,2}$ are fluctuating phases defined on the forward and backward parts of the Keldysh contour and related to fluctuating voltages $V_{1,2}$ across the conductor as $\dot{\varphi}_{1,2}(t) = eV_{1,2}$. Here and below we set $\hbar = 1$.

The effective action consists of two terms, $S[\varphi] = S_c[\varphi] + S_l[\varphi]$, where

$$iS_c[V] = \frac{C}{2e^2} \int_0^t dt' (\dot{\varphi}_1^2 - \dot{\varphi}_2^2) \equiv \frac{C}{e^2} \int_0^t dt \dot{\varphi}^+ \dot{\varphi}^- \quad (8)$$

describes charging effects and the term $S_l[V]$ accounts for electron transfer between normal and superconducting reservoirs. It reads¹⁸

$$S_l[\varphi] = -\frac{i}{2} \sum_n \text{Tr} \ln \left[1 + \frac{T_n}{4} (\{\check{G}_N, \check{G}_S\} - 2) \right], \quad (9)$$

where \check{G}_N and \check{G}_S are 4×4 Green-Keldysh matrices of normal and superconducting electrodes which product implies time convolution and which anticommutator is denoted by curly brackets. In Eq. (8) we also introduced ‘‘classical’’ and ‘‘quantum’’ parts of the phase, respectively, $\varphi_+ = (\varphi_1 + \varphi_2)/2$ and $\varphi_- = \varphi_1 - \varphi_2$.

Without loss of generality we can set the electric potential (and, hence, fluctuating phases) of the superconducting terminal equal to zero. Then the Green-Keldysh matrix of this electrode can be written in a simple form $\check{G}_S = \begin{pmatrix} \hat{G}_R & \hat{G}_K \\ 0 & \hat{G}_A \end{pmatrix}$ with

$$\hat{G}_{R/A}(t) = \pm \delta(t) \hat{\tau}_3 - \theta(\pm t) [\Delta \hat{\tau}_3 J_1(\Delta t) + i \hat{\Delta} J_0(\Delta t)]$$

and $\hat{G}_K = \hat{G}_R F - F \hat{G}_A$, where $F(t) = -iT / \sinh[\pi T t]$ is the Fourier transform of $1 - 2n(\epsilon)$ and $n(\epsilon) = 1 / (1 + e^{\epsilon/T})$ is the Fermi function. Here $J_{0,1}$ are the Bessel functions, $\hat{\tau}_i$ are the Pauli matrices, $\theta(t)$ is the Heaviside step function, and $\hat{\Delta} = i\Delta \hat{\tau}_2$, where Δ is chosen real.

The Green-Keldysh matrix of the normal terminal is defined as

$$\check{G}_N(t, t') = \frac{1}{2} \begin{pmatrix} \hat{1} & \hat{1} \\ \hat{1} & -\hat{1} \end{pmatrix} \check{Q}_N(t, t') \begin{pmatrix} \hat{1} & \hat{1} \\ \hat{1} & -\hat{1} \end{pmatrix}, \quad (10)$$

where

$$\begin{aligned} \check{Q}_N(t, t') &= \int \frac{d\epsilon}{2\pi} e^{-i\epsilon(t-t')} \begin{pmatrix} e^{i\varphi_1(t)\hat{\tau}_3} & 0 \\ 0 & e^{i\varphi_2(t)\hat{\tau}_3} \end{pmatrix} \\ &\times \begin{pmatrix} [1 - 2n(\epsilon)]\hat{\tau}_3 & 2n(\epsilon)\hat{\tau}_3 \\ 2[1 - n(\epsilon)]\hat{\tau}_3 & [2n(\epsilon) - 1]\hat{\tau}_3 \end{pmatrix} \\ &\times \begin{pmatrix} e^{-i\varphi_1(t')\hat{\tau}_3} & 0 \\ 0 & e^{-i\varphi_2(t')\hat{\tau}_3} \end{pmatrix}. \end{aligned} \quad (11)$$

Substituting the above expressions for \check{G}_S and \check{G}_N into Eq. (9) we arrive at the action which fully describes transfer of electrons between N terminal and S terminal to all orders in T_n .

In the limit of small channel transmissions one can expand S_t in powers of T_n . Keeping the terms up to $\sim T_n^2$ one recovers the contribution from Andreev reflection. At low energies this part of the action reduces to the same form¹⁷ as that for normal tunnel barriers¹⁶ in which one substitutes e by $2e$ and G_N by G_A . Here, however, we are aiming at a more general description which includes arbitrary transmission values T_n . For this reason we will proceed differently.

Let us define the matrix $\check{X}_0[\varphi_+] = 1 - T_n/2 + (T_n/4)\{\check{G}_N, \check{G}_S\}|_{\varphi_+=0}$. As the action S_t vanishes for $\varphi_-(t) = 0$ one has $\text{Tr} \ln \check{X}_0 = 0$. Making use of this property we can identically transform the action, Eq. (9), to

$$S_t = -\frac{i}{2} \sum_n \text{Tr} \ln [1 + \check{X}_0^{-1} \circ \check{X}'], \quad (12)$$

where $\check{X}' = 1 + (T_n/4)(\{\check{G}_N, \check{G}_S\} - 2) - \check{X}_0$. At temperatures and voltages well below the superconducting gap Andreev con-

tribution to the action dominates. Hence, it suffices to consider the limit of low energies $\epsilon \ll \Delta$ and set $\check{G}_S \rightarrow \begin{pmatrix} \hat{\tau}_2 & 0 \\ 0 & \hat{\tau}_2 \end{pmatrix}$. Then we obtain

$$\check{X}_0^{-1}(t, t') = \frac{2}{2 - T_n} \times \begin{pmatrix} \delta(t, t') \hat{1} & -\frac{2T_n}{2 - T_n} \sin[\varphi_+(t) - \varphi_+(t')] F(t, t') i \hat{\tau}_2 \\ 0 & \delta(t, t') \hat{1} \end{pmatrix}, \quad (13)$$

and

$$\check{X}'(t, t') = \frac{T_n}{2} \delta(t, t') \begin{pmatrix} 0 & -\sin \varphi_-(t) i \hat{\tau}_2 \\ \sin \varphi_-(t) i \hat{\tau}_2 & 0 \end{pmatrix} + T_n F(t, t') \times \begin{pmatrix} -\cos \frac{\varphi_-(t)}{2} \sin \frac{\varphi_-(t')}{2} \cos[\varphi_+(t) - \varphi_+(t')] i \hat{\tau}_2 & \left[\cos \frac{\varphi_-(t)}{2} \cos \frac{\varphi_-(t')}{2} - 1 \right] \sin[\varphi_+(t) - \varphi_+(t')] i \hat{\tau}_2 \\ \sin \frac{\varphi_-(t)}{2} \sin \frac{\varphi_-(t')}{2} \sin[\varphi_+(t) - \varphi_+(t')] i \hat{\tau}_2 & \sin \frac{\varphi_-(t)}{2} \cos \frac{\varphi_-(t')}{2} \cos[\varphi_+(t) - \varphi_+(t')] i \hat{\tau}_2 \end{pmatrix}. \quad (14)$$

Now let us assume that either dimensionless Andreev conductance $g_A = 4 \sum_n T_n$ is large, $g_A \gg 1$, or temperature is sufficiently high (though still smaller than Δ). In either case one can describe quantum dynamics of the phase variable φ within the quasiclassical approximation^{11,13} which amounts to expanding S_t in powers of (small) ‘‘quantum’’ part of the phase $\varphi_-(t)$. Employing Eqs. (12)–(14) and expanding S_t up to terms $\sim \varphi_-^2$ we arrive at the Andreev effective action

$$iS_t = iS_R - S_I, \quad (15)$$

where

$$iS_R = -\frac{ig_A}{2\pi} \int_0^t dt' \varphi^-(t') \dot{\varphi}^+(t'), \quad (16)$$

$$S_I = -\frac{g_A}{4} \int_0^t dt' \int_0^t dt'' \frac{T^2}{\sinh^2[\pi T(t' - t'')]} \varphi^-(t') \varphi^-(t'') \times \{\beta_A \cos[2\varphi^+(t') - 2\varphi^+(t'')] + 1 - \beta_A\}. \quad (17)$$

Equations (15)–(17) represent the central result of our work. It is remarkable that the action S_t is expressed in exactly the same form as that for normal conductors^{11,13} derived within the same quasiclassical approximation for the

phase variable $\varphi(t)$. In order to observe the correspondence between the action^{11,13} and that defined in Eqs. (15)–(17) one only needs to interchange

$$G_N \leftrightarrow G_A, \quad \beta_N \leftrightarrow \beta_A \quad (18)$$

and to account for an extra factor 2 in front of the phase φ_+ under \cos in Eq. (17). This extra factor implies doubling of the charge during Andreev reflection.

SHOT NOISE AND INTERACTION CORRECTION

Further analysis is formally similar to that of Ref. 11. Hence, we can immediately proceed to our final results. Let us define the average current and the current-current correlator as

$$\langle \hat{I}(t) \rangle = ie \int \mathcal{D}\varphi_{\pm} \frac{\delta}{\delta \varphi_-(t)} e^{iS[\varphi]}, \quad (19)$$

$$\frac{1}{2} \langle \hat{I}\hat{I} \rangle_+ = -e^2 \int \mathcal{D}\varphi_{\pm} \frac{\delta^2}{\delta \varphi_-(t) \delta \varphi_-(t')} e^{iS[\varphi]}, \quad (20)$$

where $\langle \hat{I}\hat{I} \rangle_+ = \langle \hat{I}(t) \hat{I}(t') + \hat{I}(t') \hat{I}(t) \rangle$. In the absence of interactions we set $\dot{\varphi}_+ = eV$ and trivially recover the standard result $I = G_A V$. For the current fluctuations $\delta I(t)$ from Eqs. (15)–(20) analogously to Ref. 11 we obtain

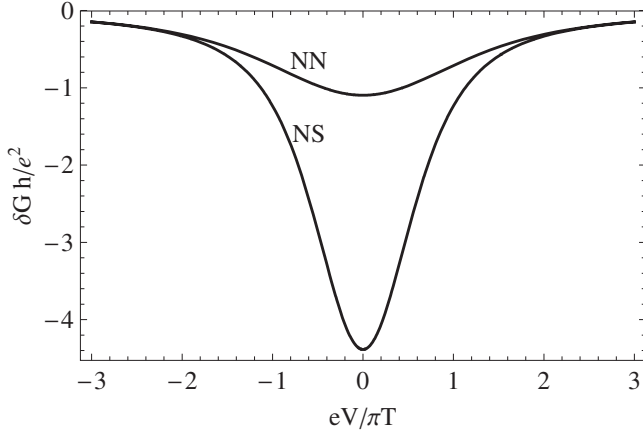


FIG. 2. The interaction correction $\delta G = dI/dV - G_N$ for short diffusive conductors at $T = G_N/2\pi C$. The upper and lower curves correspond to normal and NS structures, respectively.

$$\frac{\langle |\delta I_\omega|^2 \rangle}{G_A} = (1 - \beta_A) \omega \coth \frac{\omega}{2T} + \frac{\beta_A}{2} \sum_{\pm} (\omega \pm 2eV) \coth \frac{\omega \pm 2eV}{2T}. \quad (21)$$

This equation fully describes current noise in NS structures at energies well below the superconducting gap. For $eV \gg T, \omega$ Eq. (21) reduces to the result,⁴ Eq. (5), while in the diffusive regime the correlator, Eq. (21)—together with Eq. (6)—matches with the semiclassical result.¹⁹

Let us now turn on interactions. In this case one should add the charging term, Eq. (8), to the action and account for phase fluctuations. Proceeding along the same lines as in Ref. 11, for $g_A \gg 1$ or $\max(T, eV) \gg E_C = e^2/2C$ we get

$$I = G_A V - 2e\beta_A T \operatorname{Im} \left[w\Psi \left(1 + \frac{w}{2} \right) - iv\Psi \left(1 + \frac{iv}{2} \right) \right], \quad (22)$$

where $\Psi(x)$ is the digamma function, $w = g_A E_C / \pi^2 T + iv$ and $v = 2eV / \pi T$. This result is plotted in Fig. 2.

The last term in Eq. (22) is the interaction correction to the I - V curve which scales with Andreev Fano factor β_A in exactly the same way as the shot noise, Eq. (5). Thus, we arrive at an important conclusion: *interaction correction to Andreev conductance of NS structures is proportional to the shot-noise power in such structures*. This fundamental relation between interaction effects and shot noise goes along with that established earlier for normal conductors^{11,20} extending it to superconducting systems. In both cases this relation is due to discrete nature of the charge carriers passing through the conductor.

Another important observation is that the interaction correction to Andreev conductance defined in Eq. (22) has exactly the same functional form as that for normal conductors,

cf. Eq. (25) in Ref. 11. Furthermore, in a special case of diffusive systems due to Eq. (6) the only difference between the interaction corrections to the I - V curve in normal and NS systems is the charge doubling in the latter case. As a result, the Coulomb dip on the I - V curve of a diffusive NS system at any given T is exactly *two times narrower* than that in the normal case. We believe that this narrowing effect was detected in normal wires attached to superconducting electrodes,¹⁰ cf. Fig. 3c in that paper.²¹

It should also be pointed out that the experiment¹⁰ was carried out with SNS structures rather than NS systems considered here. Therefore the relation between our results and those¹⁰ requires extra analysis. Making use of the sample parameters¹⁰ we estimate an effective Thouless energy of the N wire to be ~ 0.08 K which is much smaller than temperature $T \gtrsim 0.3$ K employed in this experiment. Hence, for the sake of our comparison SNS systems¹⁰ can be viewed as two connected in series NS junctions. For this reason our results remain applicable to the systems¹⁰ as well. We also point out that with the aid of our theoretical expressions and experimental parameters¹⁰ we can estimate the effective capacitance as $C \approx 1.7$ fF which agrees well with the estimate¹⁰ ~ 1 fF. Also the magnitude of the interaction correction measured in Ref. 10 is consistent with our predictions.²²

The above discussion demonstrates that seemingly different experiments in Refs. 8–10 are actually closely related: doubling of the shot noise found in NS structures^{8,9} corresponds to narrowing of the I - V curves observed in Ref. 10, i.e., $e^* = q_{\text{eff}} = 2e$. The key reason behind this correspondence is the relation between shot noise and interaction correction to conductance in NS systems established above. The absolute value of this interaction correction is proportional to (effective charge) \times (shot-noise power), i.e., doubling of the shot noise in diffusive NS structures implies *four times bigger* interaction correction to conductance than in the normal case, see Fig. 2. The above predictions can be verified by independently measuring shot-noise and Coulomb-blockade effects in the same NS structure, e.g., as it was already done in normal conductors.²³

CONCLUSIONS

In summary, we theoretically described the interplay between Coulomb blockade and Andreev reflection and demonstrated a direct relation between shot-noise and interaction effects in NS systems. The fundamental physical reason behind this relation lies in discrete nature of the charge carriers—electrons and Cooper pairs—passing through NS interfaces. Further extension of our theory will include the impact of interactions on FCS.

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- ²¹We would like to point out that actual values of q_{eff} observed in the experiment (Ref. 10) were somewhat smaller than $2e$. Although at the moment we do not have a detailed quantitative explanation for this observation, we speculate that it can be attributed to additional effect of normal quasiparticles. Indeed, the experiment (Ref. 10) was performed at not too low temperatures $T \gtrsim 0.3$ K, in which case normal quasiparticles are always present above the superconducting gap thus providing an additional contribution to the interaction correction in Eq. (22). Under certain conditions this additional contribution can have the same functional form as in Eq. (22) but with the effective charge equal to e (Ref. 11). The combination of these two different interactions corrections may effectively yield values of q_{eff} bigger than e but smaller than $2e$, exactly as it was observed in Ref. 10.
- ²²Note that the ratio between the absolute values of the interaction correction in superconducting and normal cases displayed in Fig. 3c of Ref. 10 should be influenced by the fact that in the latter case superconducting electrodes become normal thus playing a role of an extra shunt resistance (of unknown value) which should yield relative increase in the interaction correction in the normal state. This effect would lead to a *decrease* in the above ratio below 4 predicted here, in a qualitative agreement with the observations (Ref. 10).
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