Mesoscopic Conductance Oscillations in a Normal Sample Controlled by the Superconductivity of an NS Boundary

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An arrangement of mesoscopic experiments is considered when the conductance of a mesoscopic normal metal sample is governed not only by the magnetic flux threading the sample but by the supercurrent inside the superconductor coupled to the normal metal sample as well. The conductance of the system is shown to be associated with the SNS Josephson current $j_J$ through the normal part of the system.

By way of example we consider a model when quasi1D wire is connected to a superconductor by its ends and coupled to two reservoirs of electrons by the normal metallic leads. The conductance oscillations caused by change of magnetic field and the supercurrent inside the superconductor have been calculated. The possibility of direct observation of the Josephson current through measurement of the thermal noise in the normal conductor is suggested.

I. INTRODUCTION

Quantum electronic transport in mesoscopic systems has been a subject of many experimental and theoretical investigations, the conductance oscillations and fluctuations being usually governed by the magnetic flux threading the sample (see for review Ref. [1] and references therein). Recently much attention was attracted by the mesoscopic systems containing normal-superconducting interfaces, due to nontrivial properties of normal transport in these systems.

Here we consider a situation where the normal conductance oscillations are controlled not only by the magnetic flux through the normal metal mesoscopic sample, but by the supercurrent inside the superconductor coupled to it. In this case the oscillations are associated with the SNS Josephson current passing through the normal part of the system.

We discuss here the following model of a normal-superconducting mesoscopic system (Fig.1). The normal leads are modelled by ideal normal wires. The scattering processes are confined to NS boundaries and T-junctions. The latter are described by real S-matrices. The system is coupled to a superconductor, S, and to two equilibrium electronic reservoirs, L and R.

The main result of this paper is as follows. In the ballistic limit (no impurity scattering and no mode mixing) and in the absence of the normal reflection at NS-boundaries the normal conductance of the system is given by the formula

$$G = \frac{2e^2}{h} N_\perp \cos \phi_J + 1 + 2\epsilon^2 + o(N_\perp).$$  \hspace{1cm} (1)

Here $\epsilon$ is proportional to the probability of a particle leaving the normal part at a junction connecting it to a reservoir, $N_\perp$ is the number of transverse modes in the lead ADCB (of length $L$ and crossection $\sim N_\perp \lambda_F^2$), and $\phi_J$ is the gauge invariant phase difference between the points A and B of the superconductor. In the limit $\epsilon \to 0$ this equation can be rewritten as follows (cf. [10])

$$G = \epsilon \left[ N_\perp \left( \frac{e^2}{\pi h} + \frac{eL}{\hbar v_F} \frac{dj_J^{(c)}}{d\phi_J} \right) + o(N_\perp) \right],$$  \hspace{1cm} (2)

where $\bar{v}_F = N_\perp^{-1} \sum_{\nu=1}^{N_\perp} v_{F,\nu}^\parallel$ is the average longitudinal Fermi velocity, $j_J^{(c)}$ is the Josephson current through ADCB [weakly connected to the normal reservoirs], controlled by the phase difference $\phi_J$.

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This relation allows us to measure directly the Josephson current vs. phase dependence simply by measuring the equilibrium current fluctuations in a system.

II. BASIC EQUATIONS

In our analysis, we assume that the size of normal part of the system, \( \mathcal{L} \), is less than both the phase breaking length, \( L_\phi = (D \tau_\phi)^{1/2} \), and the normal metal coherence length, \( L_T = (hD/k_B T)^{1/2} \) (but exceeds the superconducting coherence length \( \xi_0 \)). Here \( \tau_\phi^{-1} \) is the inelastic scattering rate, \( D \) is the diffusion constant of quasiparticles, \( k_B \) is the Boltzmann constant, and \( T \) is temperature. For the time being, we put the number of normal modes, \( N_\perp \), to unity. As usually, when discussing a coherent transport at a NS boundary, we use the two-component wave function \( 3 \) of the quasiparticle in normal part of the system:

\[
v = \begin{pmatrix} v^e \\ v^h \end{pmatrix},
\]

where \( v^{(e,h)} \) describes an electron-(hole-)like quasiparticle (with a given energy \( E \)). Due to one-dimensionality of the system, the coordinate dependence of these amplitudes is

\[
v^e(x) = v^e(0) e^{ipx}; \quad v^h(x) = v^h(0) e^{-iqx},
\]

where \( p, q \) are momenta of electron and hole respectively.

In the system we are discussing the scattering occurs only at the junctions (C, D) and NS boundaries (A, B). The latter includes both normal and Andreev scattering, which are characterized by the corresponding reflection amplitudes, \( p_e,h \) and \( \beta_{\pm} \). The quantity \( p_e,h \) is the probability amplitude for an electron-(hole-)like excitation incident from the normal lead to be reflected back as an electron (hole) (normal reflection), while \( \beta_{\pm} \) is the probability amplitude of its conversion to the hole (electron) (Andreev reflection). The latter amplitudes depend explicitly on the superconducting phase in the reflection point.

The scattering matrices in the junctions relate the amplitudes of quasiparticles in the three connected leads (see Fig.2a):

\[
w_\alpha = \sum_{\beta=1}^{3} S_{\alpha \beta} v_\beta; \quad \alpha, \beta = 1, 2, 3.
\]

Here \( v_\alpha \) (\( w_\alpha \)) is the in(out)going wave in the \( \alpha \)-th lead respectively, described by a two-component vector (Eq.(3)); \( S_{\alpha \beta} \) is a \( 3 \times 3 \) unitary matrix, each element of it being a \( 2 \times 2 \) diagonal matrix (since no electron-hole mixing occurs in the normal part of the system):

\[
S_{\alpha \beta} = \begin{pmatrix} S_{e \beta} & 0 \\ 0 & S_{h \beta} \end{pmatrix}.
\]

Now we can easily exclude the "stubs" AC and BD from consideration, using the boundary conditions at the NS boundary, which in matrix form read (see Fig.2a)

\[
v_3 = A w_3; \quad A = \begin{pmatrix} \rho_- & \gamma_- \\ \gamma_+ & \rho_+ \end{pmatrix}.
\]

Here the dot denotes the matrix multiplication in \( (e,h) \)-space, \( \rho_- = \rho_e \exp(2i\pi l) \), \( \rho_+ = \rho_h \exp(-2i\pi l) \), \( \gamma_- = \beta_-^* \exp(i(p - q)l) \), \( \gamma_+ = \beta_+ \exp(i(p - q)l) \), and \( l \) is the length of the "stub" (AC or BD).

As a result, instead of initial \( 3 \times 3 \) "supermatrix", \( S \), we obtain a \( 2 \times 2 \) one, \( \mathbf{S} \), which relates only the wave amplitudes in the lead directly connecting the normal reservoirs. The matrix elements of the two are related as follows:

\[
S_{ab} = S_{ab} + S_{a3} \cdot B \cdot S_{3b}; \quad a, b = 1, 2,
\]

where the \( 2 \times 2 \)-matrix in the \( (e,h) \)-space, \( B \), is given by

\[
B = \left( \begin{array}{cc} \left( S_{33} - \frac{\rho_+}{\tau} \right) & -\frac{\gamma_+}{\tau^2} \\ \frac{\gamma_-}{\tau^2} & \left( S_{33} - \frac{\rho_-}{\tau} \right) \end{array} \right)^{-1} \left( \begin{array}{cc} S_{33} - \frac{\rho_+}{\tau} & -\frac{\gamma_+}{\tau} \\ -\frac{\gamma_-}{\tau} & S_{33} - \frac{\rho_-}{\tau} \end{array} \right).
\]
Here

$$\Gamma = \gamma_+ \gamma_- - \rho_+ \rho_-.$$  \hspace{1cm} (10)

We will regard the case of weak coupling of the system to the normal reservoirs. Moreover, we choose the initial scattering matrices, $S$, to be real and one-parametric \[8\], and satisfying electron-hole symmetry, so that $S^e = S^h$. Then the "left" and "right" initial scattering matrix, in the junctions C and D respectively, have the form:

$$S_L = \begin{pmatrix} -\sqrt{1-2\i} & \sqrt{1} & \sqrt{1} \\ \sqrt{1} & \i & \i \\ \i & \i & \i \end{pmatrix} \quad ; \quad S_R = \begin{pmatrix} \i & \sqrt{1} & \i \\ \sqrt{1} & -\sqrt{1-2\i} & \sqrt{1} \\ \i & \i & \i \end{pmatrix}. \hspace{1cm} (11)$$

Here $0 \leq \epsilon \leq 1/2$ is a coupling parameter,

$$r = \frac{1}{2} \left[ \{1 - 2\epsilon\}^{1/2} - 1 \right], \quad t = \frac{1}{2} \left[ \{1 - 2\epsilon\}^{1/2} + 1 \right]$$

are the reflection (transmission) amplitudes between the "stub" and the lead CD,

$I$ is the unit $2 \times 2$-matrix in $(e, h)$-space.

In the absence of normal scattering at the NS boundary ($\rho_{e, h} = 0$) this leads to the following expressions for the effective scattering matrices:

$$S_L = \begin{pmatrix} Q_L & T_L \\ T_L & R_L \end{pmatrix} \quad ; \quad S_R = \begin{pmatrix} R_R & T_R \\ T_R & Q_R \end{pmatrix}.$$ \hspace{1cm} (12)

Here the submatrices have the following form:

$$T = \sqrt{1} + \frac{\sqrt{1}}{r^2 - \gamma_+^2 \gamma_-^2} \begin{pmatrix} r & -\gamma_+^* \\ -\gamma_-^* & r \end{pmatrix}; \hspace{1cm} (13)$$

$$Q = -\sqrt{1-2\i} + \frac{\epsilon}{r^2 - \gamma_+^2 \gamma_-^2} \begin{pmatrix} r & -\gamma_+^* \\ -\gamma_-^* & r \end{pmatrix}; \hspace{1cm} (14)$$

$$R = \i + \frac{t^2}{r^2 - \gamma_+^2 \gamma_-^2} \begin{pmatrix} r & -\gamma_+^* \\ -\gamma_-^* & r \end{pmatrix}. \hspace{1cm} (15)$$

Notice that the effective scattering matrices do mix electron- and holon-like parts of the quasiparticle wave function, since they include the effects of Andreev scattering at NS boundaries.

### III. NORMAL CONDUCTANCE OF THE SYSTEM

After substituting the initial scattering matrices by the effective ones (12), we reduce the initial problem to a generalized Landauer one \[12\], where a 1D normal lead with scatterers connecting two equilibrium electronic reservoirs (Fig\[3\]). The difference from the generic Landauer case is due to the necessity to account for electrons and holes separately \[3\]. The normal two probe conductance between the reservoirs at temperature $T \to 0$ is thus given by the Lambert formula \[3\]

$$G = \frac{2e^2}{h} \left\{ \langle T_{0}^{(\gamma)} \rangle + \langle T_{0}^{(\gamma)} \rangle + \frac{2\langle \langle T_{a}^{(\gamma)} \rangle \langle T_{a}^{(\gamma)} \rangle - \langle T_{a}^{(\gamma)} \rangle \langle T_{a}^{(\gamma)} \rangle \rangle}{\langle T_{a}^{(\gamma)} \rangle + \langle T_{a}^{(\gamma)} \rangle + \langle T_{a}^{(\gamma)} \rangle + \langle T_{a}^{(\gamma)} \rangle} \right\}. \hspace{1cm} (16)$$

Here $\langle \rangle \equiv \frac{\int_{0}^{\infty} \left( -\frac{\partial n_{F}(E)}{\partial E} \right) f(E)dE}{n_{F}(E)}$ is the Fermi distribution, and the excitation energy $E > 0$ is measured from the Fermi level;

$T_{0}^{(\gamma)}(E)$ is the transition probability for the electron with energy $E$ incident from the left (right) reservoir to pass to the right (left) one as an electron; $T_{a}^{(\gamma)}(E)$ is its probability to reach the opposite reservoir as a hole (due to Andreev scattering somewhere in between), and $R_{a}^{(\gamma)}(E)$ is the probability for the electron incident from the left (right) to be reflected into the same reservoir as a hole.

In the weak coupling limit $\epsilon \to 0$ these transition and reflection probabilities are given by the following formulae;
\[ T_0^> \approx \frac{e^2}{|a_+|^2|a_-|^2} \left( |a_+|^2 + |a_-|^2 \right) \] (17)

\[ 2\text{Re} \left( \exp(i(\chi \bar{\chi} - \chi \bar{\chi}) a_- a_+^* \gamma_+ \gamma_- R) \right) ; \]
\[ T_a^> \approx \frac{e^2}{|a_+|^2|a_-|^2} \left( |a_+|^2 + |a_-|^2 \right) \] (18)

\[ 2\text{Re} \left( \exp(i(\chi \bar{\chi} - \chi \bar{\chi}) a_- a_+^* \gamma_+ \gamma_- R(\gamma_+ \gamma_-)^*) \right) ; \]
\[ T_a^< \approx \frac{e^2}{|a_+|^2|a_-|^2} \left( |a_+|^2 + |a_-|^2 \right) \] (19)

\[ 2\text{Re} \left( \exp(i(\chi \bar{\chi} + \chi \bar{\chi} - \chi h - \chi h) a_+ a_-^* \gamma_+ \gamma_- R(\gamma_+ \gamma_-)^*) \right) ; \]
\[ R_a^> \approx \frac{e^2}{|a_+|^2|a_-|^2} \left( |a_+|^2 + |a_-|^2 \right) \] (20)

\[ 2\text{Re} \left( \exp(i(\chi \bar{\chi} + \chi \bar{\chi} - \chi h - \chi h) a_+ a_-^* \gamma_+ \gamma_- R(\gamma_+ \gamma_-)^*) \right) ; \]
\[ R_a^< \approx \frac{e^2}{|a_+|^2|a_-|^2} \left( |a_+|^2 + |a_-|^2 \right) \] (21)

\[ 2\text{Re} \left( \exp(i(\chi \bar{\chi} + \chi \bar{\chi} - \chi h h) a_+ a_-^* \gamma_+ \gamma_- R(\gamma_+ \gamma_-)^*) \right) . \]

Here \( \chi_{e, h}(\bar{\chi}_{e, h}) \) is the phase gained by an electron (hole) on the interval CD (resp. DC).

The resonant denominators

\[ a_+ = 1 - t^4 \gamma_+ \gamma_- R \exp(i(\chi \bar{\chi} + \chi h)) , \] (23)
\[ a_- = 1 - t^4 \gamma_- \gamma_+ R \exp(i(\chi \bar{\chi} + \chi h)) \] (24)

define the Andreev levels in the system.

As we see, the difference between various transmission and reflection coefficients is only in the quickly oscillating terms (as function of electron (hole) momentum). In the case of \( N_\perp >> 1 \) channels their relative contribution thus tends to zero, while it is quite insensitive to the temperature averaging. They correspond in our case to the Spivak-Khmel’nikskii’s oscillations \[12\). Disregarding this difference we get from (10) a simplified expression:

\[ G = \frac{2e^2}{\hbar} \left\{ \langle T_0^> \rangle + \langle R_a^> \rangle \right\} . \] (25)

The behaviour of the transition and reflection probabilities close to resonance is then governed by the ”slow” terms:

\[ T_{0,a}^{>,<}(E) \approx R_{a}^{>,<}(E) \approx e^2 \left( \frac{1}{|a_+|^2} + \frac{1}{|a_-|^2} \right) \] (26)

\[ \approx e^2 \sum_{n, \pm} \left\{ \left( \frac{2L}{\hbar \nu_F} \right)^2 \left( \frac{E - E_n^\pm}{\hbar \nu_F} \right)^2 + e^2 \left( \frac{\hbar \nu_F}{L} \right)^2 \right\}^{-1} . \]

Here the sum is taken over the resonant values of energy \( E_n \) are determined from (23-24):

\[ E_n^\pm = \frac{\pi \hbar \nu_F n}{L} + \frac{\hbar \nu_F}{2L} (\pi + \Delta \phi) . \] (27)

This is the result obtained by Kulik \[9\) for low-lying Andreev levels in a long clean SNS junction, but now due to leakage to the normal reservoirs, these levels acquire finite width \( e\hbar \nu_F / L \).

Substituting (23-24) in (23) and keeping only the terms within the accuracy of \( e^2 \), we obtain a formula

\[ G = \frac{2e^2}{\hbar} \frac{2e^2}{\cos \phi_J + 1 + 2e^2} . \] (28)

Here \( \phi_J = \phi_B - \phi_A + 2\Phi / \Phi_0 \) is the gauge invariant Josephson phase difference between the ends of a link BDCA (see Fig.1) in the presence of magnetic flux \( \Phi \) penetrating the loop BDCA; \( \Phi_0 = \hbar c / e \) is the magnetic flux quantum. [We had to keep the term \( 2e^2 \) in the denominator of (28), since it is important close to resonance, when \( \cos \phi_J + 1 \approx 0 \).]
As we can see, the normal conductance through the system has strongly resonant character (Fig.3). While we should expect the conductance of the order of \((2e^2/h)\epsilon^2\) (since the probability for an electron/hole to enter or leave the system at each of two junctions is of order \(\epsilon\)), the actual quantity has sharp peaks, where the conductance reaches its maximum possible value, \(2e^2/h\).

We can calculate the Josephson current in our system as well, e.g., using the Bardeen-Johnson approach \([14]\), which allows us to express it through the excitation spectrum at low energies. Specifically, the Josephson current in a one-mode case can be written as follows:

\[
j_j = \frac{2ev_F}{\pi L} v_s - \frac{2ev_F}{\pi L p_F} \sum_{E_n > 0} p(E_n) n_F(E_n - p(E_n) v_s).
\]

Here \(v_s = \hbar \phi_J/2mL\) is the supercurrent velocity. The first term in \((29)\) describes the (unperturbed) superflow due to superconducting phase gradient \(\phi_J/L\), and the second term is the contribution from the excited states (quasiparticle current). The interplay of these contributions gives rise to the characteristic "sawtooth" shape of the \(j_j(\phi_J)\) dependence in a long clean SNS junction \([14]\). It is important that at low temperatures only the lowest Andreev levels \(E_n\) in \((29)\) are occupied and contribute to the current, and therefore we can use in this expression the low-energy approximation for the levels \(E_n\).

If compare the expression \((30)\) to \((28)\) (Fig.3b), we see that the conductance at \(\epsilon \rightarrow 0\) actually behaves (within the accuracy of \(\epsilon^2\)) as the derivative of Josephson current with respect to the phase:

\[
G = \epsilon \left\{ \frac{\epsilon^2}{\pi h} - \frac{\epsilon L}{hv_F} \frac{d j_j^{(\epsilon)}}{d\phi_J} \right\}
\]

(cf. Eq.(2) of this paper). This result stresses the fact that while the maximum conductance is achieved when the Andreev levels in the system are tuned to the Fermi level \((17-21)\), the Josephson current is carried by the very Andreev levels (see, e.g., \([3]\)).

The above results are directly generalized to the case of \(N_\perp > 1\) non-mixing transverse modes. Indeed, the resonant denominators in Eq.(28) now acquire dependence on the transverse energies, \(E_{\perp,\nu}\), and we have to sum over the transverse mode indices, \(\nu\) (or integrate over \(dE_{\perp}\) if \(N_\perp \gg 1\)). It is easy to see that in either case the result will be simply \(N_\perp\) times the 1D conductance, if only \(k_BT < \Delta E_{\perp}\). (In a realistic case of 1000 Å wide Ag wire \(\Delta E_{\perp} \simeq 2K\).)

On the other hand, the Josephson current in a clean SNS contact is directly proportional to its area, i.e. to \(N_\perp\), so that the relations \([11,12]\) hold (we should only instead of the longitudinal Fermi velocity \(v_F\) write its average over \(N_\perp\) transverse modes, \(v_F = 1/N_\perp \sum_{\nu=1}^{N_\perp} v_{F,\nu}\)).

It is noteworthy that the amplitude of the phase-dependent conductance oscillations can now significantly exceed \(\epsilon^2/h\), which is its characteristic value in dirty SNS systems \([14,15]\). This is the manifestation of the ballistic motion of the excitations in the system.

Now we would like to suggest a direct experimental application of the result Eqs.(31). Namely, it allows a direct measurement of the phase dependence of Josephson current in a metallic wire by measuring the equilibrium thermal noise in it (Fig.4). Indeed, by virtue of Nyquist theorem \([13]\) the spectral density of equilibrium voltage noise is

\[
\langle \delta V^2 \rangle_\omega = 4k_BT R,
\]

where \(R = 1/G\) is the normal resistance of the system. The resonant dips in the noise level will then follow the phase derivative of the Josephson current through the normal bridge (Fig.4).

Let us estimate the conditions under which such a measurement is possible. The width of the Andreev level in the system, of order \(\epsilon h v_F / L\), must be larger than temperature, i.e.,

\[
\frac{T}{T_c} < \frac{\xi_0}{L}
\]

(33)
where $\xi_0 \sim h v_F / T_c$ is the superconducting coherence length. This gives an estimate $T < 2K$ for $\epsilon = 0.1$, $T_c = 10K$ and micrometer-sized system. The thermal voltage noise intensity is then confined to an interval

$$\epsilon \cdot \frac{2k_B T_c \xi_0 h}{L e^2} \leq \langle \delta V^2 \rangle_\omega \leq \frac{1}{\epsilon} \cdot \frac{2k_B T_c \xi_0 h}{L e^2},$$

(34)

which for $T_c \sim 10K, L \sim \xi$ and $\epsilon = 0.1$ yields

$$7 \cdot 10^{-19} \frac{V^2}{Hz} \leq \langle \delta V^2 \rangle_\omega \leq 7 \cdot 10^{-17} \frac{V^2}{Hz},$$

(35)

which can be measured experimentally.

In conclusion, we have investigated the normal current in a quasi1D mesoscopic system, controlled by the phase difference in the superconductor, connected to it. The magnitude of the conductance oscillations can significantly exceed the conductance quantum, $2e^2/h$. We have shown also that the normal current contains a term which behaves as the phase derivative of the Josephson current in the system. This allows to measure directly the dependence of the Josephson current on the superconducting phase difference and magnetic field, by measuring the equilibrium voltage fluctuations in the normal part of the system. The size of the effect is estimated.

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FIG. 1. The schematic view of the ballistic normal metal system in contact with the superconductor and thermal reservoirs (L and R). The scattering occurs only at the NS boundary (A and B) and in the nodes (C and D).

FIG. 2. The effective scattering matrices in the nodes C, D.
FIG. 3. Josephson current and normal conductance of the system as functions of the superconducting phase difference between A and B, $\phi_J$.

FIG. 4. Suggested scheme of equilibrium noise measurement in a clean normal wire (a) and the noise dependence on $\phi_J$. 