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P1-H4

Resonant Current Transport in HTS Junctions via Localized States in the InterlayerI

I. A. Devyatov, D.V.Goncharov and M. Yu. Kupriyanov
Institute of nuclear Physics, Moscow State University, 119899, Moscow, Russia

Abstract— Motivated by an existing difficulties in the experimental observation of the zero bias anomaly (ZBA) in HTS NID and DID structures with thick dielectric barriers we have studied theoretically effect of scattering of quasiparticles on localized states (LS) within a barrier. Direct analytical calculations give that tunneling via LS does not contribute to ZBA independently on their space position in the barrier and on the relation between the metals Fermi energy and LS energy level. It is also shown that interference current component, which usually was not taken into account in the descriptions of NIS contacts, suppressed the amplitude of ZBA provided by the direct tunneling process. The resultant ZBA peak should be very narrow and its amplitude effectively suppressed by any scatters located in the barrier. Making use of Green's function technique we have calculated a Josephson current across DID structure with LS in interlayer.

I. Introduction

Properties of interfaces of HTS materials and their influence on transport characteristics of Josephson junctions and structures normal metal – barrier – *d*-wave superconductor (NID) are now a subject of intensive theoretical [1]-[2] and experimental [3],[4] investigations. There is a large difference in behavior of the superconducting order in the vicinity of small transparent interface of s- and *d*-wave superconductors. In the last case occurs formation of mid gap states [1] at an interface. These states essentially modify the conductance of NID structures and leads to appearance zero bias anomaly (ZBA) and finite voltage anomalies.

Experimentally ZBA has been observed in point contacts and Josephson junctions made on bicrystal substrates [3],[4]. Studies of the current transport in the ramp type DID and NID structures with the thick dielectric barrier have shown that normal current component takes variable hopping channels [5]-[6], no anomalies at zero bias voltage has been observed up to now. This stimulated us to attack this problem theoretically.

II. Junction Model

We will assume that the tunnel barrier potential $V(\mathrm{r})$ can be constitute as a sum

$$V(\mathbf{r}) = V_{rect} + V_{imp}, \tag{1}$$

in which the first item represents a two dimensional rectangular barrier of height V_0 and thickness $2d$

$$V_{rect}(x) = V_0 \Theta(|x| - d), \tag{2}$$

the second describes a LS located within the barrier at the point with coordinates $r_0 = (x_0, y_0)$

$$V_{imp}(\mathbf{r}) = \begin{cases} -\beta, |\mathbf{r} - \mathbf{r}_0| \le \rho \\ 0, |\mathbf{r} - \mathbf{r}_0| > \rho \end{cases}, \tag{3}$$

Here $\rho << k_0^{-1}$ is a radius of the "defect" and $\hbar k_0$ is the Fermi momentum.

We will also assume that LS density is rather small, so that their mutual influence is negligible, and the barrier thickness d is rather large: $\kappa_0 d >> 1$, $\hbar\kappa_0 = \sqrt{2m(V_0 - \mu)}$ is the quasiparticle momentum within the barrier, m is the mass of electron, μ is the Fermi energy.

III. NID Junction

Total current across a junction can be decomposed on three independent parts: direct, interference and resonant current components. Direct component describes tunneling across the junction without LS, the interference component describes the interference between direct tunneling and tunneling via LS and resonant component describes tunneling via LS.

A. ZBA in direct tunneling channel

For zero bias conductance in the direct tunneling channel $G_d^{pot}(0)$ we have

$$G_d^{pot}(0) = \frac{2e^2}{\pi\hbar}\frac{k_0}{2\pi} L_y \int_{ZBA} d\theta \cos\theta \,. \tag{4}$$

Here ZBA denotes the integration area in which electron scatter into the regions with different sign of the order parameter, L_y is junction width. The half width of the

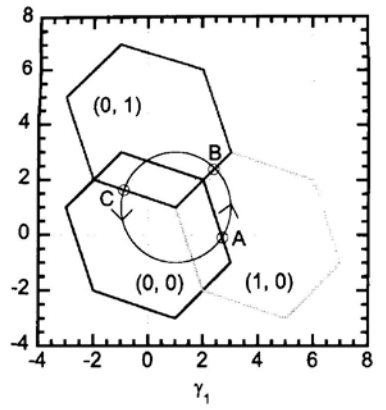

Fig. 2. Quantum diagrams for $(0,0)$, $(1,0)$, and $(0,1)$ for β_L = 0.25. Each hexagonal region is calculated by using the linearization method. Operation with two RF and one DC signals is represented as a circle. Transition points are indicated as A, B, and C.

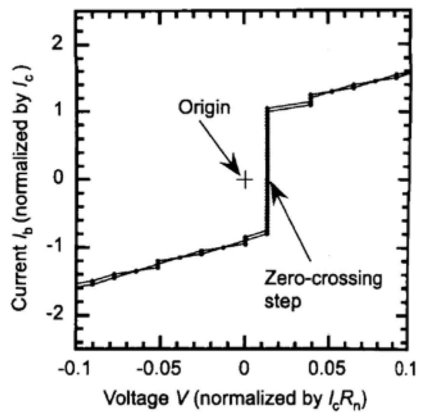

Fig. 3. Example of the calculated current-voltage characteristics. β_c of the junctions is set to 1. The normalized frequency of the RF signals is 0.013.

Fig. 2 shows the quantum diagram for $(0,0)$, $(1,0)$, and $(0,1)$, where (n_1, n_2) represents the quantum state of the SQUID. They are calculated by using the linearization method of $\sin^{-1} \simeq (\pi/2)x$ and the three conditions of $-1 \leq i_1/I_c \leq 1$, $-1 \leq (i_1 - i_2)/I_c \leq 1$, and $-1 \leq i_2/I_c \leq 1$. (Hence, each region becomes hexagonal.) β_L is set to 0.25 in Fig. 2. There is some overlap between each region. The overlap becomes larger for larger β_L.

III. Quantum Transitions and a Zero-Crossing Step

If we apply γ_1 of $(1 + 2\cos\omega t)$ and γ_2 of $(1 + 2\sin\omega t)$ to the SQUID, the operation point moves on the circle in the Fig. 2. Now let us start from the point in $(0,0)$. The operation point moves counterclockwise on the circle and goes across the point A, where the SQUID transits from $(0,0)$ to $(1,0)$. Next the operation point moves from $(1,0)$ to $(0,1)$ at the point B. And then, the operation point goes across the point C and returns to $(0,0)$. This sequence means that (i) one flux quantum enter the left loop through the left junction at the point A, (ii) it moves from the left to the right loop at the point B, (iii) and it exits from the SQUID through the right junction at the point C. These flux transitions occur without dc biasing, which results in a zero-crossing step on the current-voltage characteristics.

Fig. 3 shows an example of the calculated current-voltage characteristics. A big zero-crossing step appears at the voltage corresponding to the frequency of the RF signals.

IV. Conclusion

In this paper, we proposed a new operation of a 3J-SQUID, and numerically demonstrated a zero-crossing step on its current-voltage characteristics. The principle of generating a zero-crossing step is based on periodic flux transitions in the SQUID, and is different from the conventional method used in the Josephson voltage-standard systems [6]. The MRFDS can be used for a voltage-standard system with non-hysteretic high-T_c Josephson junctions. We can also apply the MRFDS to other applications such as a magnetometer and a logic device.

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P3-D2

Thermal flows in nonthermolised Andreev bolometer

I. A. Devyatov and M. Yu. Kupriyanov
Institute of nuclear Physics, Moscow State University, 119899, Moscow, Russia

Abstract— The energy flows in Andreev bolometer irradiated by a high-frequency signal is calculated. It is shown that the account of the form of nonequilibrum function of electron *distribution exited by a high-frequency signal, radically changes thermoelectric properties of Andreev bolometer. S - N* boundaries of the bolometer absorber can not work as the effective «mirrors» for the energy reserved in it, like it was for a case of equilibrium electron distribution function and dc signal. Also, the exited electrons loss energy to phonon subsystem with extremely high rate.

I. Introduction

Development of a new type of detectors of electromagnetic radiation - bolometers on hot electrons with Andreev reflection in superconducting electrodes [1-4] stimulated an interest for studying thermoelectrical effects in structures containing boundaries between normal (N) and superconducting (S) materials. This bolometer consists of a thin film of a normal metal (absorber), placed between superconducting electrodes. In the case of clean S - N boundary the Andreev reflection from a superconductor provides exponentially small thermoconductivity for thermolised in absorber electrons. Therefore it is possible «to focus» energy of a signal carrying on superconducting wires in absorber, and the basic channel of energy loss of overheated electrons in absorber is the transfer it to phonons. The speed of electron-phonon energy loss also is small at low temperature. For measurement of a temperature of electronic gas in absorber two tunnel S-I-N junctions located in the middle part of N-film are used. The further development of the bolometer circuit assumes utilization additional S-I-N junctions on absorber for preliminary cooling of electron gas [2].

After publication of M. Nahum and J. Martines in 1993 pioneering paper with the idea of Andreev bolometer [1] till the present time the basic attention was focus on measurements at a dc current. These measurements reported the record value of sensitivity $NEP \approx 3*10^{-18}\ WHz^{-1/2}$ at temperature 100 mK [1], and stimulated activities in this direction. At the same time, the interval of frequencies, interesting for a radioastronomical applications is located around the frequency *1 THz* . At these frequencies a quantum character of signal absorption of electromagnetic radiation in absorber might be rather essential. Quantum absorption may lead to non Fermi-like electron energy distribution function, which in turn can essentially affect on bolometer parameters. In this paper we developed a theoretical model for a calculations of a thermal flows from absorber electrons subsystem to phonons and across S - N boundaries in a limiting case completely nonthermolised *energy distribution function of electrons in N-film.*

II. Energy flow across S-N boundary

The thermal flow through clean S-N boundary in a case of isotropic s-type superconductor can be calculated in the following way [5]:

$$j = \frac{k_0^2}{2\pi^2\hbar}\int d\varepsilon\, \varepsilon\left(f_N(\varepsilon) - f_S(\varepsilon)\right)\left(1 - |a(\varepsilon)|^2\right). \tag{1}$$

Here k_0 is wave vector on a Fermi surface, $f_N(\varepsilon), f_S(\varepsilon)$ are energy distribution functions of electrons in normal metal (absorber) and in superconductor, $|a(\varepsilon)|^2$ is a square of the module of Andreev reflection coefficient (see (2) of [5]). Thermolised electrons have Fermi energy distribution function $f_{N,S}(\varepsilon) = f_0(\varepsilon) = \left(1 - \exp((\varepsilon - \mu)/k_B T_{N,S})\right)^{-1}$, where $T_{N,S}$ is the electron temperatures of the absorber and the superconductor, μ is Fermi energy. Their substitution into (1) leads to Andreev result [6] for the thermoconductivity of a clean S- N boundary (for 3D films). This thermoconductivity is exponentially small for typical bolometer parameters [1] and one can neglect it compare to a speed of a leaving of energy to phonon subsystem $G_{e-ph} = 5\Sigma V T_N^4$ [7], where V is a volume of an absorber, Σ is the material parameter.

The situation is essentially changed if we use the nonthermolised energy distribution function:

$$f_N(\varepsilon) = f_0(\varepsilon) - n_0\left(f_0(\varepsilon) f_0(-\varepsilon - h\nu) - f_0(\varepsilon - h\nu) f_0(-\varepsilon)\right), \tag{2}$$

$$n_0 = P\tau_{ee}/Vh\nu n_1,\ n_1 = n\int f_0(\varepsilon) f_0(-\varepsilon - h\nu)d\varepsilon \Big/ \int f_0(\varepsilon)d\varepsilon .$$

This function describe the electron distribution in normal metal irradiated by electromagnetic radiation with frequency ν and power P [8], n is the electron concentration. This function has characteristic step – like energy dependence at low temperatures. Previously it has been widely used in different theories [9]-[11], devoted to study both of an absorption of short laser pulses in metals [9]-[10] and the processes in transition-edge bolometers [11].

In this case from (1)-(2) the expression for a energy flow across S - N boundary has the form:

$$j = \frac{k_0^2}{2\pi^2 \hbar} 2n_0 \int_{\Delta}^{\hbar\omega} d\varepsilon\, \varepsilon \left(1 - |a(\varepsilon)|^2\right) \tag{3}$$

In contrast to a thermolised case, the energy flow j, described by (3), is not exponentially small. Moreover, the value of a energy flow increase linearly with signal frequency. For a frequency of a signal *1 THz* and concentration $n_0 = 5 * 10^{-7}$, the value of a thermal flow across S – N boundaries of the absorber is equal to $6 * 10^{-12}$ *W*, that is 5 order of magnitude larger than thermal flow across boundaries in thermolised case. The concentration $n_0 = 5 * 10^{-7}$ is correspond to photon flux power $2 * 10^{-14}$ *W* [1] and *e-e* relaxation time $\tau_{ee} = 10^{-7}\, C$ [12].

III. Energy flow to phonon subsystem

The another channel of electron energy lost in the absorber is the transfer of electron energy to phonons. In thermolised case an electron transfer energy to phonons at a rate $P_{ef} = \Sigma V(T_e^5 - T_f^5)$ [12]. The situation is significantly changed when we have non-equilibrium distribution (2). In this case Fermi's golden rule for energy transfer between electron-phonon subsystems gives [12]:

$$P = P_0 - P_1,$$

$$P_0 = \sum_{k,k'} f_N(E_k) f_N(-E_{k'}) \varepsilon_q \frac{2\pi}{\hbar} |M|^2 \delta(E_k - E_{k'} - \varepsilon_q),$$

$$P_1 = \sum_{k,k'} f_N(E_k) f_N(-E_{k'}) \varepsilon_q \frac{2\pi}{\hbar} |M|^2 n_q \{\delta(E_k - E_{k'} - \varepsilon_q) - \delta(E_k - E_{k'} + \varepsilon_q)\} \tag{4}$$

Here n_q is phonon occupation number, $|M|^2 = M_0^2 q$ is the matrix element of the electron – phonon interaction [12]. After substitution the nonequilibrium distribution (2) into (4) we come to the following relation for electron loss energy rate:

$$P_0 = \frac{\Sigma V}{\Gamma(5)\zeta(5)} \frac{n_0}{10} (h\nu)^5. \tag{5}$$

Here $\Gamma(n)$ is the gamma function, $\zeta(n)$ is the Rieman zeta function. For the bolometer parameters, described in [1], it follows that the value of spontaneous emission is approximately 10^4 larger than in thermolised case.

Thus, the carried out analysis has shown that the account of the form of nonequilibrum electron energy distribution function generated as a result of absorbtion of high-frequency signal, radically changes thermoelectric properties of Andreev bolometer. The S - N boundaries of the absorber can not work as the effective «mirrors» for the energy reserved in absorber, like it was for a case of equilibrium electron distribution function and dc signal. Also, the exited electrons loss energy to phonon subsystem with extremely high rate. The obtained results directly imposes restrictions on reduction of absorber length less $L_t = \sqrt{D\tau_{ee}} \approx 10^{-6}\, m$, where $D \approx 10^{-3}\, m^2/c$ is diffusion coefficient.

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