Energy Spectroscopy of Andreev Levels between Two Superconductors

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We perform energy spectroscopy of Andreev reflection processes occurring at two superconducting electrodes connected in series via a ballistic two dimensional channel, by measuring the voltage dependence of that part of the conductance modulated by the macroscopic phase difference. The amplitude of the modulation oscillates as a function of energy and the phase exhibits an abrupt shift close to π at the energy for which the amplitude is minimum. A microscopic theoretical description ascribes the phenomenon to the precursor of a bound state formed between the two superconductors. [S0031-9007(97)04525-0]

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The supercurrent flowing through a variety of superconducting junctions is closely related to the existence of (Andreev) bound states [1] confined in the region between the superconducting electrodes by Andreev reflection [2] (AR) processes. Each bound state carries a finite amount of current and the total supercurrent flowing through a junction is the sum of the contributions from all the bound states. Consequently the energy spectrum, which is a function of the macroscopic phase difference ϕ between the superconductors (S), determines all the properties (i.e., critical current, its temperature dependence, current phase relation, etc.) of the junction.

The existence of bound states between a superconducting pair potential and a potential barrier has been predicted theoretically and verified experimentally long ago by means of tunneling spectroscopy [3]. However, in spite of their importance, the existence of discrete Andreev levels created by superconducting confinement on either side of a normal conductor has never been directly experimentally investigated.

As an alternative to tunneling spectroscopy, whose application to this problem is not straightforward, Andreev levels can be studied experimentally by means of a superconducting analog of resonant tunneling [4]. The central idea is to inject electrons in a phase coherent ballistic conductor containing two superconducting barriers connected *in series*. Andreev (quasi)bound states, formed between these barriers, manifest themselves via resonant *dips* of the AR probability as a function of electron energy and superconducting phase difference.

In this paper we address the above problem experimentally by investigating phase coherent transport through a ballistic normal channel to which two superconducting electrodes are connected in series (Fig. 1) via high transparency interfaces. As we will demonstrate, in spite of the weak confinement provided by one of the superconductors and of the large number of quantum modes ($\approx$40) in the channel, the precursor of an Andreev level can be revealed using the conductance oscillations modulated by the superconducting phase as a very sensitive experimental tool.

Figure 1 is a schematic top view of the samples. All the samples have been realized using as a normal conductor a channel etched in the two dimensional electron gas (2DEG) present in the InAs layer of a InAs/AlSb heterostructure (electron density $N = 1.75\,10^{16}$ m^{-2}; elastic mean free path $l_e = 1.9$ μm). The dark shaded region represents the ballistic channel, where the InAs is still covered by the top layer, connected to contacts A and B. Superconducting Nb electrodes make electrical contact to the InAs layer at the two sides ($S1$) and at the end of the channel ($S2$) (in Fig. 1 the regions defined by the black rectangles). These electrodes are part of a superconducting ring to which two more contacts for electrical measurements (not shown) are attached in C. Coupling superconductor $S1$ only to the side of the 2DEG channel is presently the best technological solution to study ballistic transport in a sample in which

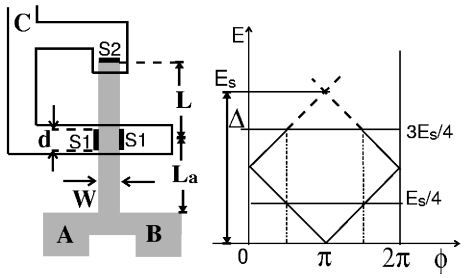

FIG. 1. Left: Schematic top view of the sample showing the 2DEG channel and the Nb ring. Superconductor $S1$ makes electrical contact with InAs only in the regions defined by the black rectangles. Right: Phase dependence of the Andreev energy levels for a 1D quantum channel between two superconductors with perfect interfaces. The picture is idealized in that the linear energy phase relation holds only for $E \ll \Delta$; E_s is the Andreev level spacing. At $E > \Delta$ the dashed line indicates the absence of true bound states.

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AR's occur in series. A high transparency Nb contact all across the 2DEG channel would damage the transport properties underneath the contact itself [5].

The relevant sample dimensions, as obtained from scanning electron microscope micrographs, are (see Fig. 1) $L \simeq 1.1\ \mu$m, $W \simeq 390$ nm, $d \simeq 250$ nm, and $L_a \simeq 1.0\ \mu$m. The samples are realized by means of three step electron beam lithography in combination with chemical wet etching, Nb electron-beam evaporation and lift-off techniques. Details of the technology and of the material properties have been reported elsewhere [6,7]. Here we emphasize only that (1) transport between the two superconducting contacts is ballistic ($L < l_e$) and (2) the angular distribution of the holes reflected by the superconductors is essentially isotropic [7].

The experiments consist in measuring the (differential) resistance of the channel as a function of the dc voltage [8] and of the magnetic field. Current is sent from contact A to C and, at the same time, the voltage across B and C is measured. The differential resistance of one of our samples measured at 1.6 K as a function of the magnetic field B is shown in Fig. 2, for different values of the applied dc voltage. The magnetic field used in the experiments is so small ($B < 10$ G) that the magnetic flux piercing the 2D channel between the superconductors is negligible compared to a flux quantum. Hence the magnetic field has the only effect of tuning the macroscopic phase difference ϕ between the two superconducting electrodes, via the flux Φ induced through the Nb ring ($\phi = 2\pi\Phi/\Phi_0$).

The change of the macroscopic phase difference induces oscillations in the resistance due to quantum interference of electrons and holes Andreev reflected by the two superconductors. Upon increasing the electron energy (i.e., the dc voltage) we observe that two related phenomena occur. The oscillation amplitude first decreases with increasing E from 0 [trace (a) in Fig. 2] to $\simeq$0.4 meV [trace (b)]. However, a further increase of the electron energy results in an increase of the amplitude, which reaches a maximum at $E \approx 0.7$ meV [trace (c)], before eventually decaying and vanishing at energies above the superconducting energy gap Δ [trace (d); $\Delta \simeq 1.2$ meV is determined from the onset of the rapid increase of AR probability observed in the measurements of $dV/dI(V_{\rm dc})$]. This behavior has been observed in all three samples investigated. The detailed energy dependence of the oscillation magnitude measured in a second sample is shown in Fig. 3(a).

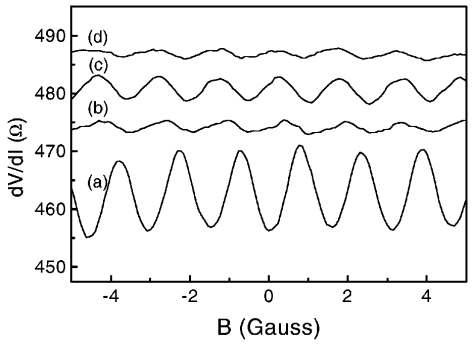

FIG. 2. Resistance oscillations at 0 mV (a), 0.4 mV (b), 0.7 mV (c), 1.3 mV (d) dc voltage (curves offset for clarity).

The data of Fig. 2 also clearly show a second phenomenon observed in all the samples, namely, that the minimum in the amplitude of the oscillations is accompanied by a shift in their phase. The phase shift occurs rather abruptly at $E \simeq 0.4$ meV and its magnitude is close but not equal to π ($\simeq 0.8\pi$). This characteristic behavior is illustrated in Fig. 3(b). At a given dc voltage, the corresponding phase has been inferred by fitting the differential resistance oscillations with a cosine function. The amplitude and phase of the conductance oscillations at energy $E < \Delta$ are described by $\delta R_{h,e}(E, \phi)$, the phase dependent part of the total AR probability [9]. In order to highlight the relevant microscopic processes giving a relevant contribution to $\delta R_{h,e}(E, \phi)$, we introduce a semiclassical model (λ_F, 21 nm, is much smaller than the sample dimensions) which provides us with a parametrization of the energy dependence of the oscillation amplitude and phase [10].

We first note that in a semiclassical description we can separate the trajectories of an electron propagating ballistically in the 2DEG channel towards the superconducting electrodes into two groups, depending on whether the electron hits electrode $S1$ or $S2$ first. Only after hitting one of the electrodes, due to the disorder present at the 2DEG/S interface, the electron is reflected into (either electron or hole) partial waves propagating in all possible directions that eventually interfere quantum mechanically. This makes it convenient to evaluate quantum interference at the position where an electron experiences the first collision with one of the superconductors [12]. It also follows that we can consider the interference of trajectories

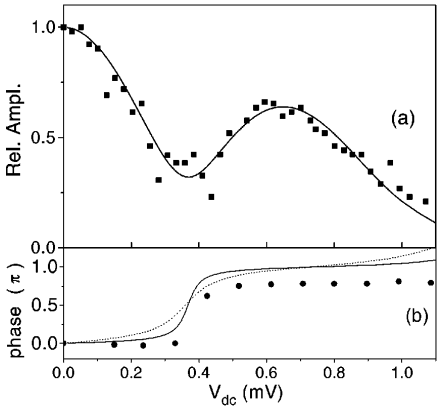

FIG. 3. (a) Energy dependence (normalized at the $E = 0$ value) of the conductance oscillation amplitude: the data (squares) are compared with theory (continuous line) in the case $c, d = 0$, $b/a \simeq 0.5$ (for the case $c, d \neq 0$ we obtain a fit of the same quality for $b/a \simeq 0.8$). (b) Energy dependence of the oscillation phase: experimental data (circles) are compared with theoretical prediction for $c, d = 0$ ($b/a \simeq 0.5$, dotted line) and $c, d \neq 0$ ($b/a \simeq 0.8$, continuous line).

in which an electron hits $S1$ first separately from those in which the electron hits $S2$ first [13].

Since superconductor $S1$ is coupled only to the sides of the 2DEG channel, the largest contribution to $\delta R_{h,e}$ is due to electron-hole trajectories hitting $S1$ the least number of times. These are the trajectories represented in Fig. 4, which describe the following processes: (1) An incoming electron is directly transformed into an outgoing partial hole wave; (2) normal reflection at $S1$ generates an electron partial wave that is Andreev reflected at $S2$ into a hole, which, after a second normal reflection at $S1$, produces an outgoing hole wave; (3) AR of the incoming electron at $S1$ generates a hole partial wave, which, after undergoing two more AR's at $S2$ and $S1$, produces an outgoing hole. Processes 4–6 are analogous to 1–3 apart from the fact that the incoming electron hits $S2$ first.

Note that in Fig. 4 we have drawn only those processes in which the electron-hole partial waves propagate from one superconducting electrode to the other and back along phase conjugated trajectories (i.e., trajectories along which electron and hole trace back each others paths), since we expect these trajectories to give the largest contribution to phase modulated quantum interference [7]. Note also that, in order to calculate the total quantum amplitude of processes 2 and 3 (and also 5 and 6), one has to sum the amplitudes of all possible phase conjugated trajectories.

We illustrate the estimate of the energy and phase dependence of the quantum amplitudes of the above processes by briefly discussing the evaluation of the amplitude A_2 of process 2. The amplitude A_2 can be factorized in terms of the amplitudes of the basic processes undergone by the electron and the hole during their motion along the trajectory represented in graph (2) of Fig. 4:

$$A_2 = \sum_i r_{e,e}^{(1)} P_i^{(e)} r_{h,e}^{(2)} P_i^{(h)} r_{h,h}^{(1)}, \qquad (1)$$

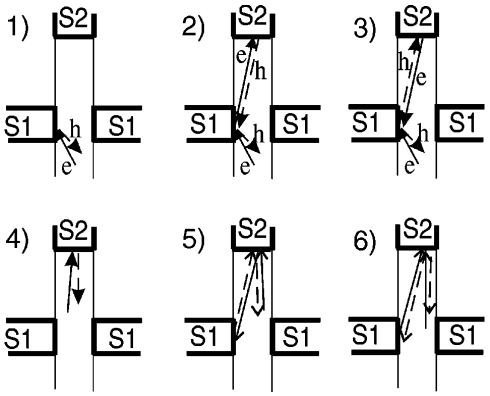

FIG. 4. The most relevant processes contributing to the conductance oscillations. Electrons (holes) are represented by continuous (dashed) lines; the arrows indicate the propagation direction (see description in the text).

where $r_{e,e}^{(1)}$, $r_{h,h}^{(1)}$, and $r_{h,e}^{(2)}$ are the amplitudes for normal reflection (respectively, electron into electron and hole into hole) at $S1$ and AR at $S2$; $P_i^{(e)} = e^{ik_e(E)x_i}$ ($P_i^{(h)} = e^{ik_h(E)x_i}$) is the dynamical phase that an electron (hole) acquires in propagating from $S1$ to $S2$ (from $S2$ to $S1$) along trajectory i, of length x_i. The $\sum_i$ is a sum over all phase conjugated trajectories connecting $S1$ to $S2$.

In order to proceed in the evaluation of Eq. (1) at energy $E < \Delta$ [14] we approximate the sum over the phase conjugated trajectories as an integral over the possible lengths ($\sum_i \approx \int_{L_{\min}}^{L_{\max}} dx_i$ with $L_{\min}$ and $L_{\max}$ minimum and maximum length of the relevant phase conjugated trajectories). In this approximation we obtain the E and ϕ dependence of A_2 (and analogously the amplitude of all the other processes shown in Fig. 4). By taking the squared modulus of the sum of the amplitudes of the processes represented in Fig. 4 we arrive at the following parametrization of $\delta R_{h,e}(E, \phi)$:

$$\delta R_{h,e}(E,\phi) = \Gamma(E)\{a\cos(\phi + k_F L_{\rm av} E/E_F) + b\cos(-\phi + k_F L_{\rm av} E/E_F) + c\cos[-\phi + k_F L_{\rm av} E/E_F - 2\arccos(E/\Delta)] + d\cos[\phi + k_F L_{\rm av} E/E_F - 2\arccos(E/\Delta)]\}. \qquad (2)$$

Here k_F and E_F are the Fermi wave vector and energy, respectively, and $L_{\rm av} = (L_{\min} + L_{\max})/2$. $\Gamma(E) = \sin(k_F \Delta L E/E_F)/(k_F \Delta L E/E_F)$, where $\Delta L = (L_{\max} - L_{\min})/2$ is the spreading in the length of the relevant phase conjugated trajectories, describes the decay of the oscillation amplitude at finite energy due to the breaking of phase conjugation. $-\arccos(E/\Delta)$ is the energy dependent part of the phase picked up in the AR process [14].

a, b, c, and d are energy and phase independent real constants that define the relative strength of the four terms in Eq. (2). With reference to Fig. 4, the terms proportional to a and c come from the interference of trajectories (1) with (2) and (1) with (3), respectively (electrons that hit $S1$ first), whereas the terms proportional to b and d are produced by the interference of trajectories (4) with (5) and of (4) with (6) (electrons that hit $S2$ first) [15].

After convoluting with the Fermi-Dirac distribution to account for the finite temperature [9] we obtain from Eq. (2) the energy dependence of the oscillation amplitude and phase. In order to decrease the number of parameters we initially neglect trajectories involving three AR's [(3) and (6) in Fig. 4] by setting $c = d = 0$ and we normalize the oscillation amplitude to the $E = 0$ value, so that the theoretical curve depends only on the ratio b/a. $L_{\rm av}$ and ΔL are also varied to optimize the comparison between model and experiments. The continuous line in Fig. 3(a) is a best fit of the oscillation amplitude energy dependence under the conditions just

discussed, with $b/a \simeq 0.5$. It is obtained for $L_{\rm av} \simeq 1.3\ \mu$m and $\Delta L \simeq 0.5\ \mu$m. These values, which cannot be varied significantly without decreasing the quality of the fit between theory and experiment, are only slightly larger than those inferred from the samples geometry. This indicates that trajectories connecting $S1$ to $S2$ and bouncing once or twice off the channel sides give a relevant contribution. As for the E dependence of the phase, Eq. (2) with $c, d = 0$ reproduces the energy at which the jump occurs, but not its shape and magnitude [Fig. 3(b)]. For the shape the situation is improved if one takes into account nonzero values for c and d. In this case one has the best fit of both amplitude and phase for $b/a \approx 0.8$ and c/a, d/a in the range 0 to ≈ 0.3 (for $L_{\rm av}$ and ΔL same values as above). However the magnitude of the jump is still not exactly reproduced, for which a more elaborate theoretical description is needed. Note also that, in spite of the high transparency of the Nb/2DEG interfaces, processes involving three AR's are not the most relevant.

Finally, our parametrization of $\delta R_{h,e}(E, \phi)$ predicts that the oscillation amplitude measured at 1.6 K is approximately 3 times larger than at 4.2 K (at $V_{\rm dc} = 0$) not far from the experimental value, 2.6.

We now discuss the significance of the experimental results in relation to the problem of Andreev levels. Well defined bound states can be formed between $S1$ and $S2$, only if the confinement due to $S1$ is strong enough. In that case the levels induce resonances [16] in the AR probability as a function of E and ϕ. However, if the confinement due to $S1$ is not strong enough, multiple reflections responsible for the full formation of the bound states have a low probability to occur and the resonances are broadened into an oscillatory dependence of the AR probability on E and ϕ, which is accounted for by the lowest order processes only. These are the processes responsible for the behavior of our samples: that is why Eq. (2), valid in the regime of low coupling to $S1$, reproduces the experimental observations.

In conclusion we like to demonstrate that, in spite of the number of modes (≈ 40) present in the channel, we can interpret qualitatively our results in terms of the idealized 1D level diagram of Fig. 1 [the energy scale in the diagram is determined by E_s, the Andreev level energy spacing, which for the idealized (linear) $E(\phi)$ relation, is $E_s = 2\pi E_F/k_F L \simeq 1.5$ mV]. As we sweep the phase ϕ at fixed electron energy E, an energy level $E(\phi)$ decreases the conductance most effectively when $E(\phi) = E$ (recall that Andreev levels produce conductance dips). For every value of E two main contributions to the conductance oscillations are therefore present [$E(\phi) = E$ has two solutions] whose relative phase is 0 at $E = 0$ (for $E = 0$ the solutions coincide at $\phi = \pi$), and increases with E. At $E_s/4 \simeq 0.38$ mV the two contributions are π out of phase, and the oscillation magnitude shows the observed minimum [Fig. 3(a)]. At $E_s/2 \simeq 0.75$ meV the relative phase is 2π, i.e., the two contributions are again in phase: this produces the local maximum in the amplitude observed at $\simeq 0.7$ mV. One expects to see a second minimum at $3E_s/4 \simeq 1.1$ mV but the oscillation amplitude keeps on decreasing at higher energy, because $E > \Delta$ and the particles are no longer efficiently reflected by the gap.

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[12] The situation is similar to the semiclassical treatment of weak localization: in that case it is convenient to evaluate quantum interference at impurity sites. [See, e.g., G. Bergmann, Phys. Rep. **107**, 1 (1984)].

[13] For the same reason we consider separately trajectories hitting $S1$ on the opposite side of the 2DEG channel.

[14] As stated in the text Eq. (2) holds if $E < \Delta$ and in the presence of highly transparent S/2DEG interfaces. In that case we can neglect the energy dependence of the modulus of r in (1) and approximate the energy dependent part of the phase acquired upon AR with $-\arccos(E/\Delta)$.

[15] Interference of trajectories (2) with (3) and of (5) with (6) give a contribution of higher order in the coupling to $S1$.

[16] In the present experimental configuration a finite probability of normal scattering at the Nb/2DEG interfaces is required to observe the resonances.