Tunable supercurrent in superconductor/normal metal/superconductor Josephson junctions

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When two superconductors are connected by a weak link a supercurrent flows determined by the difference in the macroscopic quantum phases of the superconductors. Originally, this phenomenon was discovered by Josephson [1] for the case of a weak link formed by a thin tunnel barrier. The supercurrent I is related to the phase difference φ through the Josephson current–phase relation, $I = I_c \sin \varphi$, with I_c , the critical current, depending on the properties of the weak link. A similar relation holds for weak links consisting of a normal metal, a semiconductor or a constriction [2]. In all cases, the phase difference $\varphi = 0$ when no supercurrent flows through the junction, and φ increases monotonically with increasing supercurrent until the critical current is reached. Using nanolithography techniques we have succeeded in making and studying a Josephson junction with a normal metal weak link, in which we have direct access to the microscopic current-carrying states inside the link. We find that the fundamental Josephson relation can be changed from $I = I_c \sin \varphi$ to $I = I_c \sin(\varphi + \pi)$, i.e. to a π -junction, by suitably controlling the energy distribution of the current-carrying states in the normal metal. This fundamental change in the way these Josephson junctions behave has potential implications for their use in superconducting electronics as well as (quantum) logic circuits based on superconductors.

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The microscopic mechanism responsible for the supercurrent in a Josephson junction is the transport of correlated electrons. In a superconductor/normal metal/superconductor (SNS) junction, conduction electrons mediate current transport from superconductor 1 (S1) to superconductor 2 (S2) either by ballistic or diffusive transport through the N region. In a ballistic junction, in which the elastic mean free path is larger than the length of the normal region, Andreev bound states are formed [3–5]. The dispersion relation of these states is such that each subsequent state carries a supercurrent in the positive or negative direction at a given value of the macroscopic phase difference between the superconductors; the states are degenerate if the phase is zero. The net supercurrent that flows between the two superconductors, and its direction at a given phase, depends therefore not only on the actual phase difference φ , but also on the occupation of the Andreev bound

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Fig. 1. Current-carrying density of states of a diffusive SNS junction at $\varphi = \pi/2$, evaluated in the limit of the Thouless energy for the junction $E_T \ll \Delta$ (energy gap) of the superconductor.

states. The prediction is that the electron energy-distribution function in the normal region will change the supercurrent, even reversing its direction for specific distribution functions [6–8].

Transport in metals is usually diffusive, the electron trajectories are not well defined, and Andreev bound states are no longer the natural concept to describe the supercurrent. However, electron correlations induced by the superconducting electrodes are still present, with the energy scale determined by the Thouless energy $E_{Th} = \hbar D/L^2$, where D is the diffusion coefficient. The energy spectrum of the superconducting correlations is expressed in a so-called supercurrent-carrying density of states, $\text{Im}J(E, \varphi)$, which can be calculated directly using the quasiclassical Green's function theory of superconductivity [9–12]. The energy dependence of this function is given in Fig. 1 for $\varphi = \pi/2$. The supercurrent-carrying density of states is an odd function of energy and shows a phase-dependent minigap at low positive energies with respect to the Fermi energy, above which it has a positive maximum, after which it changes sign and approaches zero at high energies. The positive and negative parts of the supercurrent-carrying density of states represent, at a given phase, energy-dependent contributions to the supercurrent in the positive and negative direction. The size and direction of the total supercurrent at a given phase difference depends therefore on the occupied fraction of these states, analogous to the occupation of the discrete Andreev bound states in a ballistic system. It is given by:

$$I_c = \frac{d}{2R_d} \int_{-\infty}^{\infty} [1 - 2f(E)] \text{Im}J(E) dE \quad (1)$$

where R_d is the normal state resistance of the junction, d is the distance between the superconductors, $f(E)$ is the electron energy distributions and I_c is the maximum supercurrent through the junction. The phase dependence of the net supercurrent is still given by the Josephson relation. At zero temperatures all states above the Fermi energy in the normal metal are empty, and those below the Fermi energy are filled. States above and below the Fermi energy can therefore be used for superconducting correlations. This can be seen by realizing that an occupation of 1 or 0 gives a nonzero value of the kernel of the integral in eqn (1). Also, taking into account the odd symmetry of both the electron distribution and $\text{Im}J(E)$, it is clear that a nonzero supercurrent can flow through the junction. However, occupation of states at an energy ε above the Fermi energy will induce a (partial) suppression of the superconducting correlations at that specific energy, resulting in a zero supercurrent if $f(E) = 0.5$. The same argument holds for depletion of states below the Fermi energy. Excess holes and excess electrons have, therefore, the same effect on the total supercurrent.

Fig. 2. The devices used to create a higher effective electron temperature in the normal channel of an SNS junction. Two superconducting electrodes (dark horizontal parts in the center of the figures) are connected via a gold film (bright part) which is also connected to external reservoirs, which allows us to inject 'normal' current through the junction. The small vertical gold bar is $1\ \mu\text{m}$ long.

Thus, a possible way to influence the total supercurrent in a SNS junction is to change the electron energy-distribution function $f(E)$. This is done by using a sample geometry in which the normal region of the SNS junction is also coupled to two large electron reservoirs (see Figs 2 and 5). By applying a voltage over the reservoirs, the electron distribution in the N region of the junction can be controlled. The exact distribution function depends on the relaxation processes inside the wire and on the applied voltage.

The easiest way to influence the energy distribution of the electrons is by raising the effective electron temperature [13]. This was done using the device shown in Fig. 2. These samples have been realized on a thermally-oxidized Si substrate by means of a two-step electron beam lithography process. In the first step, the Au pattern is defined by means of electron beam deposition and lift-off (Au thickness $\sim 40\ \text{nm}$, Au square resistance $\sim 0.5\ \Omega$). In the second step, the Nb electrodes are deposited on top of the Au using sputtering and again lift-off. The part of the Au in contact with the Nb electrodes is cleaned by Ar bombardment in a plasma, *in situ*, prior to the Nb deposition, ensuring a high-quality electrical contact between the two metals. Devices having two different values for superconducting inter-electrode separation (190 and 370 nm, respectively) have been used. The width of the Au leads used to inject the 'normal' current is 150 nm. These leads are connected to the Au side contacts and the total length from one side contact to the other is $0.7\ \mu\text{m}$. Because the electronic mean free path, $l_e \sim 40\ \text{nm}$, is significantly smaller than the geometrical dimensions, the motion of the electrons in the Au is diffusive.

Again with reference to the device of Fig. 2, a typical experiment consists of measuring the current-voltage (I - V) characteristic of the junction for different values of the current flowing from one to the other of the Au side contacts (hereafter the path followed by the control current will be referred to as 'the control line'!). The result of these measurements, performed at 1.7 K are shown in Fig. 3. The supercurrent manifests itself in the region of the I - V curves where the current flowing through the junction is $I \neq 0$, while the voltage across the superconducting electrodes is $V = 0$. It is apparent that the width of this region decreases monotonically with an increase of control current I_{control} and that, by injecting a sufficiently large control current, it is possible to completely suppress the critical current of the junction. This behavior has been observed in all the

Fig. 3. I - V curves of the device shown in Fig. 2 (left), for different values of the current through the control line ($I_{\text{control}} = 0.90, 180, 280, \text{ and } 44 \mu\text{A}$): it is evident that the critical current is entirely suppressed by the largest control current.

devices investigated: it directly demonstrates the controllability of the supercurrent by means of a ‘normal’ current flowing through the normal region of a weak link.

Referring to the explanation given above, it is quite easy to understand qualitatively the origin of the ‘interaction’ between the control current and the supercurrent. Because of the finite voltage difference present across the control line, the control current is carried by electrons which, on average, have a larger energy than those present at equilibrium in the weak link. This excess of high energy or ‘hot electrons’, modifies the electronic distribution in the weak link and tends to suppress the positive contributions in the supercurrent-carrying density of states. As a consequence, the magnitude of the supercurrent is suppressed.

A more precise analysis requires the knowledge of the shape of the nonequilibrium distribution induced by the control current. In this respect it is important to note that, in this experiment, the electrons injected in the control line from one of the side contacts have a rather low probability to scatter inelastically with phonons before reaching the opposite contact. In fact, the electron–phonon scattering time in Au at 4.2 K is ~ 1 ns. During this time an electron is able to diffuse over several microns, a distance significantly larger than the control line length.

In this transport regime, if we assume that electron–electron interaction is strong enough to bring the electrons into equilibrium among themselves, the electron population is described by a Fermi–Dirac distribution with an effective temperature,

$$T_{\text{eff}} = \sqrt{T^2 + (aV)^2} \quad (2)$$

where T is the physical temperature and a is a constant whose value, in the center of the control line, is equal to 3.2 K mV^{-1} [14]. We expect that for electrons at energies smaller than Δ (energy gap), this conclusion is not substantially influenced by the presence of the superconducting electrodes, since electrons cannot leak out into the electrodes and no heat can flow through an NS interface at that energy [15]. In Fig. 4, the dependence of the critical current on T_{eff} is compared with its dependence on the physical temperature T : the two curves essentially fall on top of each other if, however, $a = 6 \text{ K mV}^{-1}$.

We want to emphasize that it is the nonequilibrium of the electronic distribution, and not the fact that a ‘normal’ current actually flows through the junction, that is relevant in controlling the supercurrent flow. We have verified this result experimentally, using samples in which we can force a ‘normal’ current through Au sidelines parallel to the junction, instead of through the junction itself [13]. In that case no net current

Fig. 4. Critical current as a function of the physical temperature (filled squares) and of the effective electron temperature (open circles) [13].

flows through the Au between the Nb electrodes, however, the electronic distribution is modified by ‘hot electrons’ diffusing from the Au sideline, which are also allowed in this configuration to completely suppress the supercurrent flow.

On the other hand, in a mesoscopic wire the distribution of electrons over the energies is not necessarily thermal. As shown recently by Pothier *et al.* [14, 16], a normal wire attached to two large electron reservoirs at a voltage difference V may have a double-step structure. If the wire is sufficiently short, so that the diffusion time τ_D is smaller than both the electron–electron relaxation time τ_{e-e} , as well as the electron–phonon relaxation time τ_{e-p} , the electrons will conserve their energy over the length. This is in contrast to the preceding experiment, where $\tau_{e-e} < \tau_D < \tau_{e-p}$. The energy distribution of the electrons in the channel is now given by the superposition of the Fermi–Dirac distributions of the reservoirs. This results, at sufficiently low temperatures $kT \ll eV$, in a position-dependent double-step function, as shown in Fig. 5. The energy separation between the steps is eV . As predicted recently [9, 11], such an energy distribution will have a more profound effect on the supercurrent in an SNS junction than the experiment described above. For sufficiently large control voltages the direction of the supercurrent can even be reversed. The microscopic mechanism responsible for this is the fact that superconducting correlations are blocked over an energy interval $|2 eV|$ around the Fermi level, for within this region the electron occupation fraction of the energy levels in the N region is $1/2$, leading to a zero kernel in the integral of eqn (1). The form of the supercurrent-carrying density of states $\text{Im}(J(E))$ is such that all states carrying current in the positive direction will be blocked if a control voltage is applied of around 17 times the Thouless energy. This leaves only states with a minus sign available for the superconducting correlations, obviously resulting in a reverse of the supercurrent direction at a given phase. The stable zero current state, which corresponds to $\varphi = 0$ for a conventional Josephson junction, corresponds now to $\varphi = \pi$ at large control voltages, resulting in a current–phase relation given by $I = I_c \sin(\varphi + \pi)$, hence the name π -junction. This mechanism should not be confused with the π -junction behavior induced by magnetic impurities [17, 18] or resulting from the symmetry of the order parameter in ceramic superconductors [19]. Not only are the microscopic mechanisms different in these cases, but in our case the junction is also controllable, its state depending on the applied control voltage.

To be able to apply a nonthermal distribution function to the normal conducting weak link we used the device shown in Fig. 6 [20]. The fabrication process is similar to the device shown in Fig. 2. The main difference is that the gold channel (40 nm thick and 200 nm wide) is connected to two very thick (475 nm) electron reservoirs of millimeter size and two voltage probes at the end of the channel. The reservoirs are deposited in a last, extra deposition step. They are of millimeter size because they should also act as effective

Fig. 5. B, Schematic representation of the sample layout. The control voltage over the channel induces a position-dependent electron distribution, shown in A, for positions 1, 2 and 3 in B. The current through the Josephson junction is indicated by I_{SNS} [20].

cooling fins to prevent unwanted electron heating at high control voltages and hence, to maintain the desired step-like electron distribution [21]. The absence of those large reservoirs in the device in Fig. 2 is the reason that the step-like distribution function could not be obtained in this sample, even at the lowest temperatures, resulting in an absence of the described reversing of the supercurrent. The distance between the two niobium electrodes is 350 nm. The measurements are performed at 100 mK in order to realize a sharp Fermi–Dirac distribution function inside the reservoirs. RC filtering at room temperature and copper powder filtering at 100 mK is used to reduce external noise and hence, unintentional heating of the electrons. The gold of the junction has a diffusion coefficient of $180 \text{ cm}^2\text{s}^{-1}$, resulting in a Thouless energy $E_T \sim 140 \text{ } \mu\text{eV}$, so that at the measuring temperature $k_B T = 10 \text{ } \mu\text{eV} < E_T < \Delta = 1500 \text{ } \mu\text{eV}$ (k_B is Boltzman's constant). The length of the control channel between the reservoirs, L , is $1 \text{ } \mu\text{m}$, resulting in an estimated diffusion time of $\tau_D = 50 \text{ ps}$. This is much smaller than the estimated energy relaxation time for electrons in a thin dirty gold film [22]. The contacts between the gold and the niobium and the gold channel and the reservoirs are cleaned *in situ* using Argon sputtering to ensure a high interface transparency.

In the experiment, the I – V curves of the SNS junction are measured for different values of the control voltage over the control channel (V_{control}). From these curves we determine the critical current of the junction as a function of the control voltage. The results are given in Fig. 7. The critical current of the SNS junction shows a nonmonotonic behavior as a function of control voltage. At zero control voltage the $I_c R_n$ product, $200 \text{ } \mu\text{V}$, is in the order of the Thouless energy, in good agreement with the theory of diffusive SNS junctions [23]. When a voltage is applied the distribution function inside the gold wire changes, and therefore the occupation of the energy levels carrying the supercurrent changes. A decrease of the supercurrent is observed, reaching zero at $V_{\text{control}} = 520 \text{ } \mu\text{V} \equiv V_{\text{critical}}$, corresponding to approximately four times the Thouless energy. As can be verified in the inset, no sign of a supercurrent can be detected. At higher voltages the supercurrent reappears. In Fig. 7 we plot the critical supercurrent in this region with a minus sign,

Fig. 6. Device used to measure the π -behavior. The niobium is weakly coupled by means of a gold channel connected to two large thick reservoirs (top and bottom). Voltage probes are connected to the gold channel directly.

Fig. 7. Critical current of the SNS junction as a function of the control voltage. The current is plotted negative above $V_{\text{critical}} = 0.52$ meV because of the π -state of the junction. The inset shows selected I - V curves for the following control voltages: 0.38 mV (1), 0.44 mV (2), 0.52 mV (3), 0.64 mV (4), 0.84 mV (5), 1.70 mV (6) [20].

Fig. 8. A, V_{SNS} and the resistance of the control channel, B, as a function of the current for three different values of the control voltage. The curves correspond to curves 2, 3 and 4 of the inset of Fig. 7 (offset for clarity). The resistance variation changes sign, although the corresponding I - V curves look similar for $V_{\text{control}} < V_{\text{critical}}$ (2) or $V_{\text{control}} > V_{\text{critical}}$ (4) [20].

anticipating that the junction has entered the π -state. Above voltages of 1.7 meV no supercurrent can be detected anymore.

To determine the current-phase relation, $I = I_c \sin(\varphi)$ or $I = I_c \sin(\varphi + \pi)$, we need an independent way to determine the phase difference over the junction for a fixed direction of the applied supercurrent. For this we recall that the normal conductance of the control channel is also modulated by the phase difference φ , a phenomenon known as Andreev interferometry and extensively studied in recent years [24, 25]. The modulation of the conductance arises from the fact that Andreev reflected holes scatter from both superconductors with a different phase shift and, therefore, might interfere constructively ($\varphi = 0$) or destructively ($\varphi = \pi$). This leads to a phase-dependent increase in the diffusion constant in the N region. The resulting phase dependence of the resistance is approximately given by $\Delta R = -R_0(1 + \cos(\varphi))$, with R_0 a temperature and energy-dependent amplitude.

In Fig. 8 we show the resistance of the normal control channel (Fig. 8B) together with the observed I - V curves of the SNS junction (Fig. 8A). Clearly the dependence of ΔR on the applied current changes sign upon crossing the critical control voltage. The middle curve (2) shows no modulation at all, and also the corresponding I - V curve is linear. Suppose the current-phase relation is given by the conventional Josephson relation $I = I_c \sin(\varphi)$. Then the phase difference between the two superconductors changes from $-\pi/2$ through 0 to $\pi/2$, if I_{SNS} is varied from $-I_c$ to $+I_c$. Hence the resistance of the control channel will show a minimum at $I_{\text{SNS}} = 0$, corresponding to $\varphi = 0$. This is observed for all control voltages $V_{\text{control}} < V_{\text{critical}}$. The upper curve in Fig. 8B is an example of the behavior for $V_{\text{control}} > V_{\text{critical}}$. Now a maximum is observed in the resistance for $I_{\text{SNS}} = 0$, consistent with the assumption that the current-phase relation is now given by $I = I_c \sin(\varphi + \pi)$ and $\varphi = \pi$ when $I_{\text{SNS}} = 0$. The phase π between the two superconductors appears

whenever $V_{\text{control}} > V_{\text{critical}}$. This is taken as a direct proof that the junction switches from a normal state to a π -state as a function of the control voltage.

The qualitative dependence of the critical current on control voltage and the energy on which the phase jumps, are all in agreement with recent theoretical work on diffusive SNS junctions [9, 10]. However, the relative magnitude of the supercurrent in the π -state range of control voltages is smaller than one would expect in comparison to the critical supercurrent. The most important reason for this is probably the geometry of the sample. In our case the width over length ratio W/L is about 0.4, so that the distribution function is not constant over the entire junction length. This will reduce the magnitude of the supercurrent in the π -regime, as predicted theoretically [9].

The considerations just made indicate the concrete possibility of using the hot-electron tunable supercurrent principle to realize superconducting transistors which perform better than those realized in the past. Apart from possible general use, which cannot be seriously predicted at this stage, these transistors can be useful in research applications at dilution refrigerator temperature. In addition, more specific applications of our devices are possible. A few examples are very fast bolometer mixers [25] with a large intermediate frequency bandwidth and any superconducting quantum interference device (SQUID) application where one would like the two junctions in the SQUID to have variable (e.g. finely tunable or modulable at high frequency) supercurrents.

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