Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.

F. S. Bergeret 1, V.V. Pavlovskii 2, A. F. Volkov 1,2 and K. B. Efetov 1,3

(1) Theoretische Physik III,
Ruhr-Universität Bochum, D-44780 Bochum, Germany

(2) Institute of Radioengineering and Electronics of the Russian Academy of Sciences, 103907 Moscow, Russia

(3) L.D. Landau Institute for Theoretical Physics, 117940 Moscow, Russia

November 1, 2018

Abstract

We calculate the conductance of a superconductor/ferromagnet (S/F) mesoscopic structure in the dirty limit. First we assume that the ferromagnet exhibits a homogeneous magnetization and consider the case that the penetration of the condensate into the F wire is negligible and the case in which the proximity effect is taken into account. It is shown that if the exchange field is large enough, the conductance below the critical temperature $T_C$, is always smaller than the conductance in the normal state. At last, we calculate the conductance for a F/S structure with a local inhomogeneity of the magnetization in the ferromagnet. We demonstrate that a triplet component of the condensate is induced in the F wire. This leads to a increase of the conductance below $T_C$. 

1 Introduction

In the last decade transport properties of mesoscopic superconductor/normal metal (S/N) structures were intensively studied (see for example the review articles [1, 2] and references therein). It was established that in these nano-structures, i.e. in structures whose dimensions are less than the phase coherence length $L_\phi$ and the inelastic scattering length $l_\epsilon$, the conductance changes when decreasing the temperature below the critical temperature $T_c$ of the superconducting transition and this variation may be both positive ($\delta G > 0$) and negative ($\delta G < 0$) [3, 4, 5, 6, 7, 8, 54]. The increase or decrease of the conductance $G$ depends, in particular, on the interface resistances and is determined by a competition between two contributions caused by the proximity effect. One of them is due to the suppression of the density-of-state (DOS) and leads to a decrease of the conductance. The second one results in increasing $G$ and is similar to the Maki-Thompson term [9, 10].

These studies apparently were stimulated by the theoretical work [11] in which a weak-localization correction to the conductance of a S/N/S mesoscopic structure was calculated and an oscillatory dependence of this correction on the phase difference $\varphi$ between the superconductors S was predicted. These oscillations have been indeed observed [4, 5, 7, 12, 13, 14, 15] but their amplitude turned out to be two orders of magnitude larger than the predicted one. In order to overcome this discrepancy between theory and experiment one should take into account the proximity effect. It was shown that the latter leads to much larger amplitudes of the conductance oscillations than the weak-localization corrections [16, 17].

The proximity effect manifests itself also in other interesting peculiarities of the transport properties of S/N structures. One of them is an interesting dependence of the Josephson current $I_c$ on an additional dissipative current $I_{ad}$ through a N wire in a 4-terminal S/N/S structure similar to the one shown in the inset of Fig.4. According to theoretical predictions the current $I_c$ changes sign if the current $I_{ad}$ is large enough [18, 19, 20, 21, 22]. This behavior of the Josephson current ($\pi$-contact) was later experimentally confirmed [23]. Another interesting effect is a non-monotonic temperature $T$ (or voltage $V$) dependence of the correction to the conductance $\delta G$ [24, 4, 5, 6, 8]. When decreasing the
temperature, $\delta G$ increases, reaches a maximum and with further decrease of the temperature drops to zero. This behavior has been explained theoretically for the case of a short S/N contact ($E_{Th} >> \Delta$) and a "long" S/N structure ($E_{Th} << \Delta$) \[25\], here $E_{Th} = \hbar D/L^2$ is the Thouless energy, $D$ is the diffusion constant, $L$ is the length of the N wire (film). The reason for this behavior is the competition between two contributions to the conductance mentioned above.

It is interesting to note that low-energy states play an important role in transport properties of S/N structures. The reason is that the condensate penetrates into the N wire over a length $\xi_\varepsilon = \sqrt{\hbar D/\varepsilon}$ which may be much larger than the thermodynamic correlation length $\xi_T = \sqrt{\hbar D/2\pi T}$ provided the characteristic energy $\varepsilon \approx E_{Th} << T$. Therefore, in the limit of low (compared to $T$) Thouless energy $E_{Th}$, the phase coherence is maintained over distances $\xi_\varepsilon = \sqrt{\hbar D/E_{Th}}$ of the order of the length $L$ of the N wire. These long-range phase-coherent effects have been predicted in \[28, 52\] and observed on a 4-terminal S/N/S structure \[29, 30\]. In Ref.\[34\] the conductance oscillations related to the phase coherence were observed in a S/N/S mesoscopic structure. It was established for the first time that these oscillations survive in a temperature range where the Josephson coupling between the superconductors becomes negligible. The authors of Ref.\[29\] observed an increase in the Josephson critical current when an additional dissipative current was injected into the N region in a 4-terminal S/N/S structure. This current leads to long-range effects affecting the critical current.

Recently, similar investigations have been carried out on mesoscopic F/S structures in which ferromagnets (F) are used instead of normal (nonmagnetic) metals. It is well known that in the dirty limit, when the relaxation momentum time $\tau$ is very small ($\tau << h/h_{ex}$, where $h_{ex}$ is the exchange energy), the condensate penetrates into a ferromagnet over a length $\xi_F = \sqrt{\hbar D/h_{ex}}$. The latter is extremely short (5-50 Å) for strong ferromagnets like Fe or Ni. Therefore one might expect that the influence of the proximity effect on transport properties of such structures should be negligibly small. Experiments carried out recently on F/S structures showed however that the conductance variation $\delta G$ is quite visible (varying from about 1 to 10%) when the temperature decreases below $T_c$ \[30, 31, 32\]. It is worth mentioning that the conduction variation was both positive and negative. In
some experiments the variation $\delta G$ was related to a variation of the interface conductance (resistance) $\delta G_F$, whereas in others $\delta G$ to the conduction variation of the ferromagnetic wire $\delta G_F$. The theory also predicts both an increase $\delta G_F$ and decrease $\delta G_F$ of the conductance.

In Refs. $[33, 34, 35]$ a ballistic contact was analyzed. It was shown that at $h_{ex} = 0$ the contact conductance $G_{F/S}$ is twice as large as its conductance $G_{F/N}$ in the normal state (above $T_c$), which agrees with a theoretical prediction for a N/S ballistic contact, and drops to zero at $h_{ex} = E_F$, where $E_F$ is the Fermi energy. The conductance of a diffusive point contact $G_{F/S}$ has been calculated by Golubov $[34]$ who showed that $G_{F/S}$ is always smaller than the conductance $G_{F/N}$ in the normal state. In the case of a mixed conductivity mechanism (partly diffusive and partly ballistic) $G_{F/S}$ has been calculated in Ref. $[35]$ and it may be both larger or smaller than the conductance in the normal state $G_{F/N}$.

In order to obtain the resistance of the system shown in Fig.4 one should add the interface resistances $R_b$ and the resistance of the F wire $R_F$. The conductance of the F wire $G_F = R_F^{-1}$ under the assumption that $R_b$ is sufficiently small has been calculated in Refs. $[36, 37]$. As in Ref. $[35]$, the proximity effect was neglected. It turned out that the conductance of the F wire in the normal state was $G_{F_n} = G_\uparrow + G_\downarrow$ and was equal to $G_{F_s} = 4G_\uparrow G_\downarrow/(G_\uparrow + G_\downarrow)$ in the superconducting state (these formulas are valid for a F wire shorter than the spin relaxation length, see $[37]$). Thus, the conductance of the F wire decreases below $T_c$. The mechanism responsible for this behavior is in this case purely kinetic, since the form of the distribution function depends only on the boundary conditions at the interfaces: if the interface transparencies are perfect the distribution function is continuous in the normal state, while in the superconducting state the boundary condition is changed providing the spinless current through the F/S interface. These two different boundary conditions for the distribution function lead to the different values of the conductance above and below $T_c$. In the present paper we analyze the conductivity of a F/S structure in the dirty limit when the condition $\tau << h/h_{ex}$ is satisfied.

In the next section we study the kinetic mechanism of the conductance variation in a F/S meso-
scopic structure neglecting the proximity effect. The diffusion coefficients and the interface transparencies are assumed to be different for each spin direction. It will be shown that in this case the conductance variation is always negative, i.e. the conductance decreases with decreasing temperature below $T_c$.

In the third section we take into account the proximity effect (the condensate penetration into the ferromagnetic wire) and present the results for the conductance on the basis of exact calculations assuming equal diffusion constants and interface transparencies for spin up and down. We will see that for a strong exchange field $E_{ex} >> T_c$, the resistance variation $\delta R$ is positive at any temperature and interface resistance.

Finally, in section 4 we calculate the conductance variation assuming that the magnetization near the F/S interface has a local inhomogeneity. It will be shown that not only a singlet component arises in the F wire in this case but also a triplet one which penetrates into the F wire over a long distance of the order of $\xi_m = \sqrt{D/\epsilon_m}$, where $\epsilon_m = \min\{T, E_{Th}\}$. The penetration of the triplet component leads to a positive $\delta G$ which may be comparable with the conductance variation observed in N/S structures. The singlet component exists only in a very short region near the F/S interface and leads to a negligible contribution to the conductance.

2 Kinetic mechanism of the conductance variation

In this section, we calculate the conductance of the system shown in Fig.1. We neglect influence of the proximity effect on the conductance assuming that the penetration length of the condensate is very small in comparison with the length of the F wire. The penetration length of the condensate is equal to $\xi_F = \sqrt{\hbar D/\hbar_{ex}}$ in the dirty limit and to the mean free path $l = v_F \tau$ in the "clean" limit (to be more exact, in the limit of a strong ferromagnet, when the condition $\tau >> \hbar/\hbar_{ex}$ is satisfied).

The conduction variation is related in this case to different forms of the distribution function in the normal and superconducting states. The form of the distribution function is determined by boundary conditions at the interfaces F/N ($T > T_c$) or F/S ($T < T_c$). In the normal state the spin current
is finite, whereas in the superconducting one it is equal to zero. In the superconducting state at low temperatures, charge is transferred via Andreev reflections, i.e. an incident electron is reflected from the F/S boundary and moving back as a hole with the opposite spin direction (in the S region the charge is carried by singlet Cooper pairs with zero total spin). In order to find the charge or spin current we need to calculate the distribution function $N_\alpha = n_\alpha(\epsilon) - p_\alpha(\epsilon)$ obeying the diffusion equation

$$D_\alpha \partial^2_{xx} N_\alpha = I_{in}(N_\alpha). \tag{1}$$

Here $n_\alpha(\epsilon)$ is the distribution function of electrons for a given spin direction $\alpha$, $p_\alpha(\epsilon) = 1 - n_\alpha(-\epsilon)$ is the distribution function of holes, $D_\alpha$ is the diffusion coefficient which depends on the spin index $\alpha$, $I_{in}(N_\alpha)$ is the inelastic collision integral. The inelastic collision integral is of the order of $N_\alpha/\tau_{in}$, where $\tau_{in}$ is the inelastic scattering time. We ignore spin relaxation processes in the F wire assuming that the spin relaxation length exceeds the length $L$. It will be seen that the conductance in the superconducting state is always larger than in the normal one. The simplest formulas are obtained in the case of a perfect F/S interface. Therefore, we assume that there is no barrier at the F/S interface and the F'/F interface transparency is arbitrary (F’ is the ferromagnetic or normal reservoir). In the next section we will show that the non-zero reflection coefficient at the F/S interface leads also to an increase of the resistance when temperature is lowered below $T_c$.

The boundary conditions at the F'/F interface have the same form in the normal and superconducting states (see Appendix)

$$D_\alpha \partial_x N_\alpha = (\gamma R_b)_\alpha^{-1}(N_\alpha(0) - F_V); \quad x = 0 \tag{2}$$

Here $N_\alpha(0) = N_\alpha(x = 0)$, $R_b\alpha$ is the barrier resistance per unit area for a given spin direction $\alpha$, $\gamma = e^2\nu$, $\nu$ is the density-of-states (DOS). The function $F_V = [\tanh((\epsilon+eV)/2T) - \tanh((\epsilon-eV)/2T)]/2$ is the distribution function in the F' reservoir. At the F/S interface the boundary conditions are different in the normal and superconducting case. Above $T_c$ they are
Below $T_c$ they have the form (see Appendix)

$$N_\alpha = -N_{\bar{\alpha}}; \ x = L$$

(4)

Physically, Eqs.(3) and (4) mean that the electric potential at the perfect F/N(S) interface is chosen to be zero. The potential $V(x)$ is expressed in terms of $N_\alpha$ as $eV(x) = (1/4) \int d\epsilon [N_\uparrow + N_\downarrow]$. The condition (4) implies that the spin current is zero. This fact is a result of the Andreev reflections. It is worth mentioning that, in the model under consideration, the conductance related to the Andreev reflections differs from zero only if the amplitude of the condensate function at the F/S interface is not equal to zero. Therefore, strictly speaking, we neglect the proximity effect everywhere in the F wire except the nearest neighborhood of the F/S interface. The electrical current is given by the formula

$$I_Q = (1/4e) \int d\epsilon [(\gamma D)\alpha \partial_x N_\alpha + (\gamma D)_{\bar{\alpha}} \partial_x N_{\bar{\alpha}}]$$

(6)

The spin current $I_M$ is given by the same formula with sign “−” instead of “+” in the brackets and a factor of $g\mu/e$ in front of the integral ($\mu$ is the Bohr magneton). The problem is reduced to finding a solution for Eq.(3) with the boundary conditions (3-5). The inelastic collision integral has a complicated form and finding a solution is in general an unrealistic task. Therefore, we consider limiting cases only.

2.1 Mesoscopic limit ($L << l_{in}$)

In this case of a long inelastic relaxation length $l_{in} = \sqrt{D_{in}}$, we can neglect the inelastic collision integral and seek a solution in the form
\[ N_\alpha(x) = N_\alpha(0) - J_\alpha x/L. \] (7)

where \( J_\alpha = N_\alpha(L) - N_\alpha(0) \) is a constant that determines the current and, hence, the conductance. The constants \( N_\alpha(0) \) and \( J_\alpha \) can be determined from the boundary conditions (3-5) and are equal to

\[ J_\alpha = F_V \left( \frac{G_b}{G_L + G_b} \right)_\alpha, \quad N_\alpha(0) = J_\alpha L. \] (8)

in the normal case, and

\[ J_\alpha = F_V \left( \frac{2G_{L\alpha}}{G_{L\alpha}G_{L\alpha}(R_{b\alpha} + R_{b\alpha}) + G_{L\alpha} + G_{L\alpha}} \right), \quad N_\alpha(0) = F_V \left( \frac{G_{L\alpha}G_{L\alpha}(R_{b\alpha} - R_{b\alpha})}{G_{L\alpha}G_{L\alpha}(R_{b\alpha} + R_{b\alpha}) + (G_{L\alpha} + G_{L\alpha})} \right). \] (9)

in the superconducting case, where \( R_{L\alpha} = \rho_\alpha L \) is the resistance of the F wire in the normal state. Knowing the distribution function we can compute the conductances and the spin current \( I_M \). In the normal state, we obtain

\[ G_n = \left( \frac{G_bG_L}{G_b + G_L} \right)_\uparrow + \left( \frac{G_bG_L}{G_b + G_L} \right)_\downarrow, \quad I_M = (g\mu/e) V \left[ \left( \frac{G_bG_L}{G_b + G_L} \right)_\uparrow - \left( \frac{G_bG_L}{G_b + G_L} \right)_\downarrow \right]. \] (10)

whereas the result for the superconducting state reads

\[ G_s = \frac{4G_{L\uparrow}G_{L\downarrow}G_{b\uparrow}G_{b\downarrow}}{G_{L\uparrow}G_{L\downarrow}(G_{b\uparrow} + G_{b\downarrow}) + (G_{L\uparrow} + G_{L\downarrow})G_{b\uparrow}G_{b\downarrow}}, \quad I_M = 0. \] (11)

As follows from Eqs. (10-11), the spin current in the normal state is absent only if the conductances in each spin band \( G_{b\alpha} \) and \( G_{L\alpha} \) are the same. In the superconducting state this current is always zero. The conductances \( G_{b\alpha} \) and \( G_{L\alpha} \) enter the formulas for \( G_n \) and \( G_s \) (10-11) symmetrically. In the case of a small \( F'/F \) interface resistance, we obtain for \( G_n \) and \( G_s \) the formulas presented in [37] and in the introduction. The conductance in the normal state \( G_n \) is always larger than or equal to (if conductances in each spin channel are the same) the conductance in superconducting state \( G_s \).

Indeed, the difference between \( G_n \) and \( G_s \) can be written in the form
\[ G_n - G_s = \left[ G_{L\uparrow}G_{L\downarrow}(G_{b\uparrow} - G_{b\downarrow}) + G_{b\uparrow}G_{b\downarrow}(G_{L\uparrow} - G_{L\downarrow}) \right]^2 / D_{F/N}D_{F/S}, \]  

where \( D_n \) and \( D_s \) are the denominators in the expressions. Therefore, in the model under consideration (dirty limit, no condensate penetration, zero F/S interface resistance) the conductance decreases with decreasing temperature below \( T_c \).

Note that the formula for the conductance Eq.(11) is valid at low temperatures \( (T << \Delta) \) when the contribution of the states with \( |\epsilon| > \Delta \) to the conductance can be neglected. At arbitrary temperatures the conductance can be easily found with the help of Eq.(6) and has the form

\[ G_s(T) = G_s(0) \tanh(\Delta/2T) + G_n(1 - \tanh(\Delta/2T)) \]

where \( G_s(0) \) is the conductance of the F/S structure at zero temperature determined by Eq.(11). We took into account that at \( |\epsilon| > \Delta \) the boundary conditions for the distribution function are given by Eq.(3).

### 2.2 Semi-mesoscopic limit \( (L >> l_{in}) \)

Let us consider another limiting case when the length \( L \) of the F-wire exceeds the energy relaxation length \( l_{in} \). In this case the inelastic scattering term in Eq.(1) dominates. It turns to zero provided the distribution function has the equilibrium form

\[ N_\alpha(x) = \left[ \tanh((\epsilon + eV_\alpha(x))/2T) - \tanh((\epsilon - eV_\alpha(x))/2T) \right]/2. \]

where the potential \( V_\alpha \) depends on the spin index \( \alpha \) and is linear as a function of the coordinate \( x \):

\[ V_\alpha(x) = V_\alpha(0) - E_\alpha x. \]

Integrating Eq. (2) over the energies we obtain a relation between the current in the \( \alpha \) spin band and \( V_\alpha(0) \)

\[ I_\alpha = G_{b\alpha}(V - V_\alpha(0)) \]

Let us consider first the normal case. In the F wire we have
\[ I_\alpha = G_{Lo} V_\alpha(0)/L. \]  

From the last two equations we find \( V_\alpha(0) \)

\[ V_\alpha(0) = V \frac{G_{ba}}{G_{ba} + G_{Lo}}. \]  

The electric field and the total current are equal to the sum of \( E_\alpha \) and \( I_\alpha \) over spin indices respectively. We easily obtain an expression for the conductance \( G_\alpha \), which is identical to Eq. (10).

In the superconducting case the current \( I_\alpha \) is given by

\[ I_\alpha = G_{Lo}(V_\alpha(0) - V_\alpha(L))/L. \]  

The potentials \( V_\uparrow(L) \) and \( V_\downarrow(L) \) are related to each other by Eq. (4): \( V_\uparrow(L) = -V_\downarrow(L) \). Although the form of the distribution function differs from Eqs. (7-8), the conductance \( G_s \) below \( T_c \) remains unchanged with respect to the previous case (see Eq. (11)). Thus, regardless of the relationship between the lengths \( L \) and \( l_{in} \), the conductance of the structure under consideration is given by Eqs. (10-11), i.e. it decreases with decreasing temperature below \( T_c \).

### 3 Singlet component and proximity effect

In the preceding section we neglected the proximity effect and assumed that the F/S interface resistance was equal to zero. We have shown that the conductance decreases with decreasing temperature provided the conductivities of the ferromagnetic wire or the F'/F interface for spin up and down differ from each other. In the present section we calculate the conductance for an arbitrary F/S interface resistance taking into account the proximity effect (the penetration of the superconducting condensate into the F wire). In this case the problem becomes more complicated and, for simplicity, we assume that conductivities for both spin directions are the same. An equation for the distribution
function can be solved exactly and this allows one to present the normalized zero-bias conductance

\[ S = R_{F/N}(\partial I/\partial V) \]

in the form

\[ S = (1/2) \int d\epsilon \frac{(1 + r_L + r_0)\partial F_\nu/\partial V}{(1/M) + r_0/\nu(0) + r_L/(g_{QP} + g_A)}, \]

(19)

where \( r_{L,0} = R_{0,0}/R_L \) are the normalized resistances of the F/S and F'/F interfaces respectively (in the normal state), \( 2M(x) = 1 + |f^R(x)|^2 + |g^R(x)|^2, \nu(0) = \text{Re} g^R(0), \nu_S = \text{Re} g^R_S, g_{QP} = \nu(L)\nu_S \) is the quasiparticle normalized conductance at a given energy, \( g_A = \text{Im} f^R(L) \text{Im} f^R_S \) is the normalized subgap conductance (the conductance related to Andreev reflections), \( f^R_S = \Delta/\sqrt{\epsilon^2 - \Delta^2} \). The angle brackets mean the averaging over the length: \( \langle ... \rangle = \frac{1}{L} \int_0^L dx/\langle ... \rangle \). The physical meaning of the denominator is simple: the first term in the angle brackets is the normalized resistance of the F wire in the presence of the condensate, the second and third terms are the resistances of the F'/F and F/S interfaces below \( T_c \) respectively. One can see that the ”Andreev” conductance \( g_A \) is not zero only if the condensate function at the interface (from the ferromagnetic side) \( F^R(L) \) differs from zero. The factor \( (1 + r_L + r_0) \) in the numerator arises from the normalization of the conductance. In order to calculate the normalized conductance \( S \), we have to find the Green functions \( f^R(x) \) and \( g^R(x) \) in the ferromagnetic wire. These functions obey the Usadel equation which upon the substitution \( g^R(x) = \cosh u(x) \) and \( f^R(x) = \sinh u(x) \) acquires the well known form

\[ D\partial_{xx}^2 u + 2i(\hbar \epsilon + \epsilon) \sinh u = 0. \]

(20)

This equation is complemented by the boundary conditions

\[ r_0 L \partial_x u = \sinh u, \quad x = 0 \]

(21)

\[ r_L L \partial_x u = F^R_S \cosh u - G^R_S \sinh u, \quad x = L \]

(22)
The solution for Eq. (20) can be found analytically in some limiting cases [9, 39]. It has an especially simple form in the case of a weak proximity effect when $r_L \leq 1$ and Eq. (19) may be linearized (numerical calculations show that the linearized solution differs from the exact solution by less than 10% for $r_L \approx 1$). Here we present the results of a numerical solution of the Usadel equation Eq. (20) and of the calculation of the conductance $S$. In Fig. 2 the temperature dependence of the conductance variation $\delta S = S - 1$ is presented for a structure with good interface transparencies at various exchange energies $h$ (normalized to the Thouless energy $hD/L^2$).

It is seen from this figure that $\delta S$ is positive only if $h_{ex}$ is not too large compared to $\Delta$ or $T_c$. A small positive value $\delta S$ is observed near $T_c$ for $h_{ex}/\Delta = 5$ and 10. For $h_{ex}/\Delta \leq 20$ the conductance variation is negative and decreases with decreasing $T$. For example, in the case of F/Al structure the threshold exchange energy $h = 20\Delta$ above which the conductance deviation is negative for all temperatures, is equal to $h_{ex} = 44 K$, that is, this value is much less than the characteristic magnitudes of $h_{ex}$ for ferromagnets as Fe, Ni,Co etc.

A slight increase of $G_s$ above $G_n$ is related to the condensate penetration into the F wire. This increase occurs only near $T_c$ because in this temperature range the conductance $G_s$ is close to the conductance $G_{F/N}$ in the normal state due to a large contribution of the quasiparticle current. When the temperature decreases, the quasiparticle contribution also decreases and the weak (at large $h$) proximity effect can not overcome this decrease of the conductance. The contribution to the conductance due to proximity effect is suppressed even stronger if the F/S interface resistance is not small compared to $R_L$. In Fig. 3 we plot the temperature dependence of $\delta S$ for the case when the F/S interface resistance (in the normal state) is two times larger than the conductance of the F wire. One can see that the normalized conductance decreases with decreasing temperature and becomes very small at low $T$. The reason for this behavior of the conductance is quite clear. At high enough temperatures the main contribution to the conductivity is due to quasiparticles with energies $\epsilon > \Delta$. With lowering the temperature, this contribution decreases. The contribution caused by Andreev reflections (the term $g_A$ in Eq. (19)) at large $h$ and not small $r_L$ is very small as the condensate amplitude is small.
Indeed, in this case one can linearize the Usadel equation (20) and obtain for \( f^R \)

\[
f^R(x) = f^R_S/(r_L \sqrt{-2i\hbar}) \exp(\sqrt{-2i\hbar}(1 - x/L)) \tag{23}
\]

It follows from this equation that at \( h_{ex}/E_{Th} = h \gg (2r_L)^{-2} \) the amplitude of \( f^R(L) \) is small (we consider low energies \( \epsilon << \Delta \)). Therefore the "Andreev conductance" \( g_A \) is also small. The suppression of the Andreev reflections in the model considered follows directly from the Usadel equation and the widely used boundary conditions of Ref. [40]. It is worth mentioning that the Andreev reflections are responsible for a subgap conductance in N/I/S junctions, which was observed in Ref. [41]. The suppression of the Andreev conductance by the exchange field is analogous to the suppression of the subgap conductance by an external magnetic field[42, 43].

4 \hspace{1cm} \textbf{Triplet component and long-range proximity effect}

In this section, we consider again a ferromagnetic wire with a metallic contact with a superconductor. In the previous sections we have shown that the proximity effect in the presence of a strong exchange field \( h_{ex} \) could be neglected, since the superconducting condensate penetrates into the F wire only over a short distance (\( \sim \sqrt{\hbar D/h_{ex}} \) in the dirty limit). At the same time, the interface resistance \( R_{bs} \) increases when the normal reservoir becomes superconducting. In short, the total conductance decreases with decreasing temperature below the superconducting critical temperature \( T_c \).

However, in recent experiments on S/F structures a considerable increase of the conductance below \( T_C \) was observed [30, 32]. The measurements also demonstrate that the entire change of the conductance is due to an increase of the conductivity of the ferromagnetic wire as if the superconducting condensate penetrated into the ferromagnet. A question arises: how can the condensate function penetrate into the F wire over large distances? We understand that the Cooper pairs must be destroyed by the strong exchange field because the electrons of the pairs can no longer have opposite spins.

To answer this question we notice that such a simple argument about the destruction of the condensate is based on the assumption that only the singlet pairing exists in the sample. This argument
cannot be correct if a triplet component of the superconducting condensate penetrates the ferromagnet. An arbitrary exchange field cannot destroy the triplet superconducting component since both the electrons of the Cooper pair are in the same spin band. In this section, we suggest a mechanism of formation of the triplet pairing, which is due to a local inhomogeneity of the magnetization in the vicinity of the S/F interface. We will show that the penetration length of the triplet component into the ferromagnet is of the order $\sqrt{\hbar D/\epsilon}$, where the energy $\epsilon$ is of the order of the temperature $T$ or the Thouless energy $E_{TH}$, and therefore the increase of the conductance due to the proximity effect may be comparable with that in a S/N structure. We consider the system shown in Fig. 4 and assume that there is a domain wall in the region $0 < x < w$ described by the angle $\alpha(x)$ between the magnetization $M$ and the $z$-axis. In this region the magnetization is given by

$$M = h_{ex} \left(0, \sin \alpha(x), \cos \alpha(x) \right).$$

(24)

For simplicity we assume that $\alpha(x)$ varies linearly according to $\alpha(x) = Qx$. We consider again the diffusive limit corresponding to a short mean free path and to the condition $h_{ex} \tau \ll 1$, which allows us to describe the system using the Usadel equation\[44\]. We assume that the resistance of the F/S interface is not too small and therefore the condensate amplitude $|\hat{f}|$ is small. We use the linearized Usadel equation for the retarded matrix (in spin space) Green function $\hat{f}^R$, which has the form (the index $R$ is dropped)

$$- iD\partial_x^2 \hat{f} + 2\epsilon \hat{f} - 2\Delta \sigma_3 + \left(\hat{f} \hat{V}^* + \hat{V} \hat{f} \right) = 0.$$  

(25)

Here the matrix $\hat{V}$ is defined as $\hat{V} = h_{ex} (\hat{\sigma}_3 \cos \alpha(x) + \hat{\sigma}_2 \sin \alpha(x))$. We also assume that the diffusion coefficient is the same for the both spin bands. This equation is supplemented by the boundary condition at the S/F interface, which after linearization takes the form \[10, 45\]

$$\partial_x \hat{f} \bigg|_{x=0} = \left(\rho/R_b\right) F_S ,$$

(26)

where $\rho$ is the resistivity of the ferromagnet, $R_b$ is the S/F interface resistance per unit area in the normal state, and $F_S = \hat{\sigma}_3 \Delta / \sqrt{\epsilon^2 - \Delta^2}$. We have assumed that there are no spin-flip processes at the
S/F interface.

The solution of Eq. (25) in the region $0 < x < w$ can be sought in the form

$$\hat{f} = \hat{U}(x) \hat{f}_n \hat{U}^*(x),$$

(27)

where $\hat{U}$ is an unitary transformation given by $\hat{U}(x) = \hat{\sigma}_0 \cos (Qx/2) + i \hat{\sigma}_1 \sin (Qx/2)$.

Substituting Eq. (27) into Eq. (25) we obtain the following equation for $\hat{f}_n$

$$-iD\partial_{xx}^2 \hat{f}_n + i(DQ^2/2) \{\hat{f}_n + \hat{\sigma}_1 \hat{f}_n \hat{\sigma}_1\} + 2 \epsilon \hat{f}_n + h_{ex} \{\hat{\sigma}_3, \hat{f}_n\} = 0.$$

(28)

where {...} is the anticommutator.

In the region $x > w$ the magnetization is homogeneous and $\hat{f}_n$ satisfies Eq. (28) with $Q = 0$. We see from Eq. (28) that the singlet and triplet component of the condensate function are mixed by the rotating exchange field $h_{ex}$. In the region $x > w$ these components decouple and they should be found by matching the solutions at $x = w$ (the function $\hat{f}$ is continuous in the entire F wire). The solution of Eq. (28) can be written in the form

$$\hat{f}_n = \hat{\sigma}_0 A(x) + \hat{\sigma}_3 B(x) + i\hat{\sigma}_1 C(x).$$

(29)

Here the function $C(x)$ is the amplitude of the triplet component, whereas $A(x)$ and $B(x)$ describe the singlet pairing. The structure of Eq. (28) allows to seek these amplitudes in the form

$$A(x) = \sum_{i=1}^{3} \left( A_i \exp(-\kappa_i x) + \bar{A}_i \exp(\kappa_i x) \right)$$

(30)

The functions $B(x)$ and $C(x)$ can be written in a similar way. The coefficients $A_i$, $B_i$ and $C_i$ obey the algebraic equations

$$(\kappa^2 - \kappa_i^2 - Q^2) C - 2(Q\kappa) A = 0$$

$$(\kappa^2 - \kappa_i^2) B - \kappa_i^2 A = 0$$

(31)

$$(\kappa^2 - \kappa_i^2 - Q^2) A - \kappa_i^2 B + 2(Q\kappa) C = 0,$$
where \( \kappa^2 = -2i\epsilon/D \) and \( \kappa_h^2 = -2h_{ex}/D \) (indices \( i \) were dropped). The eigenvalues \( \kappa_i \) are the values at which the determinant of Eqs. (31) turns to zero. In the case of a strong exchange field \( h_{ex} \), \( \kappa_h \) is large (\( \kappa_h \gg \kappa, Q \)) and the eigenvalues \( \kappa_i \) are given by

\[
\kappa_1, 2 \approx \frac{1 \pm i}{\xi_F}, \text{ for } 0 < x < L
\]

\[
\kappa_3 = \begin{cases} 
\sqrt{\kappa^2 + Q^2} & \text{for } 0 < x < w \\
\kappa_e & \text{for } w < x < L
\end{cases}
\]

The eigenvalues \( \kappa_{1,2} \) correspond to a sharp decay of the condensate in the ferromagnet, while \( \kappa_3 \) is associated with the slowly varying part. With the boundary condition, Eq. (26), we see that in a homogeneous case (\( Q = 0 \)) the amplitude of the triplet component \( C(x) \) is zero. If \( Q \neq 0 \), the function \( C(x) \) is coupled to the singlet components \( A(x) \) and \( B(x) \), and we consider this case.

If the width \( w \) is small, the triplet component changes only a little in the region \( (0, w) \) and spreads over a large distance of the order \( |\kappa^{-1}_e| \) in the region \( (0, L) \). In the case of a strong exchange field \( h_{ex} \), \( \xi_F \) is very short (\( \xi_F \ll w, \xi_T \)), the singlet component decays very fast over the length \( \xi_F \), and its slowly varying part \( B_3 \) is \( B_3 = 2 (Q\kappa_3/\kappa_h^2) C_3 \ll C_3 \). In order to obtain the expression for the triplet component \( C(x) \), we use Eq. (26) at \( x = 0 \) and assume that the solution vanishes at \( x = L \). Then, we find

\[
C^{(A)}(x) = \mp i \left\{ Q\epsilon(0) \sinh (\kappa_e(L - x)) [\kappa_e \cosh \Theta_3 + \kappa_3 \sinh \Theta_3 \cosh \Theta_3]^{-1} \right\}^{R(A)}, \quad (34)
\]

where \( w < x < L \), \( B^{R(A)}(0) = (\rho \xi_h/2 R_b) f^{R(A)}_S \) is the amplitude of the singlet component at the S/F interface, \( \Theta_e = \kappa_e L, \Theta_3 = \kappa_3 w, \) and \( \kappa_e^{R(A)} = \sqrt{-2i\epsilon/D} \).

It is clear from Eq.(34) that, at the interface, the triplet component is of the same order of magnitude as the singlet one. Indeed, for the case \( w \ll L \) we obtain from Eq. (34) \( |C(0)| \sim B(0)/\sinh \alpha_w \), where \( \alpha_w = Qw \) is the angle characterizing the rotation of the magnetization. Therefore, provided the angle \( \alpha_w \leq 1 \) and the S/F interface transparency is not too small, the singlet and triplet components are not small. They are of the same order in the vicinity of the S/F interface but, while the singlet component decays fast over a short distance (\( \sim \xi_F \)), the triplet one varies smoothly along the ferro-
magnet turning to zero at the F reservoir (see Fig.4). One can see also that the singlet component
oscillates, which is well known [46].

The penetration of the triplet component into the ferromagnet is similar to the penetration of the
superconducting condensate into a normal metal. The presence of the condensate function (triplet
component) in the ferromagnet can lead to long-range effects and therefore to a significant change of
the conductance of a ferromagnetic wire in a S/F structure (see inset in Fig.4) when the temperature
decreases below $T_c$. The normalized conductance variation $\delta S = (G_s - G_n)/G_n$ is given by the
expression [39, 47]:

$$\delta S = \frac{-1}{32 T} \text{Tr} \int d\epsilon F'_V \left\langle \left[ \hat{f}_R(x) - \hat{f}_A(x) \right]^2 \right\rangle.$$  

(35)

Here $G_n$ is the conductance in the normal state, $<..>$ denotes the average over the length of the
ferromagnetic wire between the F reservoirs, and $F'_V$ is given by the expression

$$F'_V = \frac{1}{2} \left[ \cosh^{-2}((\epsilon + eV)/2T) + \cosh^{-2}((\epsilon - eV)/2T) \right].$$  

(36)

Substituting Eqs. (29, 34) into Eq. (35) one can determine the temperature dependence $\delta S (T)$ shown
in Fig.4. We see that $\delta S$ increases with decreasing temperature and saturates at $T = 0$. In order to
explain the reentrant behavior of $\delta S(T)$ observed in Refs. [30, 32], one should take into account other
mechanisms as those analyzed in Refs. [37, 34, 48] and in section 2. However, this question is beyond
the scope of the present analysis.

It is also interesting to note that a triplet component of the condensate function with the same
symmetry (odd in frequency $\omega$ and even in momentum $p$) has been suggested long ago by Berezinskii
[49] as a possible phase in superfluid $^3He$ (this, so called “odd superconductivity”, was discussed
in a subsequent paper [50]). Being symmetric in space, this component is not affected by potential
impurities, in contrast to the case analyzed in Ref. [51], where the triplet component of the condensate
was odd in space. While this hypothetical condensate function is not realized in $^3He$ (in $^3He$ it is
odd in $p$ but not in frequency), this odd (in $\omega$) triplet component does exist in the system considered
here, although under special conditions described above.
As it was mentioned in the introduction, experimental data are still controversial. It has been established in an experiment [31] that the conductance of the ferromagnet does not change below $T_c$ and all changes in $\delta S$ are due to changes of the S/F interface resistance $R_b$. However, in other experiments $R_b$ was negligibly small [30]. The mechanism suggested in our work may explain the long-range effects observed in the experiments [30, 32]. At the same time, the result of the experiment [31] is not necessarily at odds with our findings. The inhomogeneity of the magnetic moment at the interface, which is the crucial ingredient of our theory, is not a phenomenon under control in these experiments. One can easily imagine that such inhomogeneity existed in the structures studied in Refs. [30, 32] but was absent in those of Ref. [31]. The magnetic inhomogeneity near the interface may have different origins. Anyway, a more careful study of the possibility of a rotating magnetic moment should be performed to clarify this question.

5 Conclusion

We have analyzed the conductance variation $\delta G_s \equiv G_s - G_n$ of the F/S mesoscopic structure in the dirty limit when the condition $h_{ex} < \hbar/\tau$ is satisfied. First, neglecting the proximity effect we have shown that below $T_c$ the conductance variation $\delta G_s$ is negative and its magnitude increases with increasing the difference between the conductances for spin up and down. The change of the conductance is related to changes in the boundary conditions at the F/S interface. Below $T_c$ one of the boundary conditions requires the spin current to be zero. The formulas for the conductance remain valid even if inelastic scattering (not spin flip) processes are taken into account.

We also studied how the proximity effect affects the conductance for an arbitrary transparency of the F/S interface. The account for the condensate penetration (the proximity effect) leads to an increase of the conductance near $T_c$. This is possible, however, only if the exchange energy $h_{ex}$ is not too large: the parameter $hr_L^2$ should not be large, where $h$ is the exchange energy in units of the Thouless energy $\epsilon_{Th} = \hbar D/L^2$ and $r_L$ is the ratio of the F/S interface resistance in the normal state to the resistance of the F wire. If the parameter $hr_L^2$ is large, the conductance $G_s$ decreases
with decreasing temperature and becomes very small at low $T$. This behavior is related to a strong suppression of the Andreev reflection processes that determine the conductance at low $T$ by the exchange field.

It is important to note that we neglected the inverse proximity effect, that is, suppression of superconductivity in $S$. In the case of large $h$ one has to take into account this effect and calculate a deviation of the Green’s functions in the superconductor from their bulk values. It turns out that in the case of a large exchange energy $h$, the energy gap in the superconductor may be suppressed and the DOS $\nu_S$ is not zero even at $|\epsilon|<\Delta$. In this case the quasiparticle conductance is not zero and the charge transfer is possible through the F/S interface at low energies $\epsilon$. The conversion of the quasiparticle current into the current of Cooper pairs is realized in the superconductor over the coherence length $\xi_S$ (at low energies).

At last, we have calculated the conductance of a F/S structure in which the magnetization is non-homogeneous near the F/S interface. It was shown that, in this case, a triplet component arises alongside with the singlet one. This triplet component in another systems (He$^3$, high $T_c$ superconductors etc) was studied in several papers [50, 49]; it is odd in Matsubara frequencies (the so called odd superconductivity) and even in momentum space (in the main approximation in the parameter $(h/h_{ex}\tau)$). The triplet component penetrates into the F wire over a large distance of the order $\sqrt{hD/T}$ and leads to a positive conductance variation. The singlet component penetrates over much shorter length of the order $\sqrt{hD/h_{ex}}$ and leads to a negligible contribution to the conductance. It would be interesting to realize experimentally the situation that we analyzed in the last section and to observe the triplet component not in an exotic system but in an ordinary F/S structure.

We would like to thank SFB 491 for financial support.
Here we show how the boundary conditions for the distribution functions $N_\alpha = n_\alpha - p_\alpha$ can be derived from matching conditions for the quasiclassical Green functions, where $n_\alpha$ is the distribution function of electrons and $p_\alpha = 1 - n_\alpha(-\epsilon)$ is the distribution functions of holes (having in mind a superconductor, it is better to speak about electron- and hole-like excitations). We use the boundary conditions for the quasiclassical matrix Green functions $g$ the matrix elements of which consist of the retarded (advanced) Green’s functions $\hat{g}^R$ and the Keldysh Green’s function $\hat{g}^K$. The last matrix is expressed in terms of the matrix distribution function $\hat{f}$: $\hat{g}^K = \hat{g}^R \hat{f} - \hat{f} \hat{g}^A$, where the matrix $\hat{f}$ can be represented in the form: $\hat{f} = f_1 \hat{1} + f_3 \hat{\sigma}_3$. The first function $f_1$ determines the order parameter in the superconductor and is equal $f_{1\alpha} = 1 - (n_\alpha + p_\alpha)$. The second function $f_3$ determines the electrical or spin current and equals $f_{3\alpha} = -(n_\alpha - p_\alpha)$. The matching condition at the F(N)/S interface for the $4\times4$ matrix Green function has the form

$$L\hat{g} \partial_x \hat{g} = \left(1/2r_L\right)[\hat{g}, \hat{g}_S].$$

(37)

where $r_L = R_b/\rho L$ is the F/S interface resistance (in the normal state) normalized to the resistance of the F wire $\rho L$. The physical meaning of Eq.(A1) is simple. If we take the Keldysh (the element (12)) of Eq.(A1), multiply it by $\hat{\sigma}_3$ and calculate the trace, we obtain the current at a given energy $\epsilon$. If we integrate over energies, we obtain on the left side the usual expression for the current

$$I_Q = (1/4)e\nu_n D \text{Tr} \hat{\sigma}_3 \int \text{d}\epsilon [\hat{g}^R \partial_x \hat{g}^K + \hat{g}^K \partial_x \hat{g}^A].$$

(38)

where $\nu_n$ is the DOS in the normal state. On the right hand side we obtain the current at a given energy which is well known in the tunnel Hamiltonian method. We take the Keldysh component of Eq. (A1), multiply it by $\hat{\sigma}_3$ and $\hat{1}$ and calculate the trace. We obtain thus the equations

$$M_3 \partial_x f_3 = (1/r_L) \left[ (\nu \nu_s + g_\sigma) f_3 - (g - f_{eq} + g_\sigma f_1) \right].$$

(39)
\[ M_1 \partial_x f_1 = \left(1/r_L \right) \left[ (\nu \nu_s + g_1) (f_1 - f_{eq}) - g_3 \right] \]  

(40)

where \( \nu = \text{Re} \, g^R \), \( \nu_S = \text{Re} \, G_S^R \) is the DOS in the ferromagnet and superconductor, respectively and

\[
M_3 = (1/2)[1 - g^R g^A - f^R f^A] \\
M_1 = (1/2)[1 - g^R g^A + f^R f^A] \\
g_1 = -(1/4)(f^R - f^A)(F^R_s - F^A_s) \\
g_\pm = +(1/4)\{(f^R \mp f^A)(\hat{F}^R_s \pm \hat{F}^A_s)\}_3 \\
g_A = -(1/4)(f^R + f^A)(F^R_s + F^A_s)
\]

The symbol \( \{\ldots\}_3 \) stands for \( \{\ldots\}_3 = (1/2)\text{Tr} \, \hat{\sigma}_3 \{\ldots\} \). In the case under consideration \( f^{R(A)} \propto i \hat{\sigma}_2 \) and \( F^{R(A)}_S \propto i \hat{\sigma}_2 \), so that \( g_\pm = 0 \). We are interested in the states with subgap energies (low temperatures) \( |\epsilon| < \Delta \). For these states one has \( F^R_s = F^A_s = \Delta/\sqrt{\Delta^2 - \epsilon^2} \) and \( G^R_s = G^A_s = \epsilon/\sqrt{\Delta^2 - \epsilon^2} \). Therefore we obtain that \( g_1 = 0 \), and \( \nu_S = 0 \), but \( g_A \neq 0 \). One can show that at any non-zero \( r_L \) the functions \( M_3 \) and \( M_1 \) are not zero and we find from Eqs.(A3-A4)

\[
\partial_x f_1 = 0 \implies \begin{cases} 
\partial_x (n_\uparrow + p_\downarrow) = 0 \\
\partial_x (n_\downarrow + p_\uparrow) = 0 
\end{cases}
\]

(41)

From Eq.(A5) we find the boundary condition

\[
\partial_x N_\uparrow - \partial_x N_\downarrow = 0
\]

(42)

This condition means the absence of the spin current. It may be easily generalized to the case of different diffusion coefficients

\[
D_\uparrow \partial_x N_\uparrow - D_\downarrow \partial_x N_\downarrow = 0
\]

(43)
At small $r_L$ (a good F/S contact) the condition (A3) yields

$$f_3 = 0 \implies \begin{cases} n_\uparrow - p_\uparrow = 0 \\ n_\downarrow - p_\downarrow = 0 \end{cases} \quad (44)$$

We took into account that in the sub-gap region the "Andreev" conductance $g_A$ is not zero. From Eq.(A8) we get

$$N_\uparrow = -N_\downarrow \quad (45)$$

In the normal state the distribution function is continuous across a good F/S interface. Hence we get $n_\alpha = p_\bar{\alpha} \implies N_\alpha = 0$. We use these boundary conditions in section 2.

References

[1] C. W. J. Beenakker, Rev. Mod. Phys. 69, 731 (1997).
[2] C. Lambert and R. Raimondi, J. Phys. Cond. Matt. 10, 901 (1998).
[3] V. T. Petrashov and V. N. Antonov, JETP Lett. 54, 242 (1991).
[4] A. Dimoulas, J. P. Heida, B. J. van Wees, and T. M. Klapwijk, Phys. Rev. Lett. 74, 602 (1995).
[5] P. Charlat, H. Courtois, P. Gandit, D. Mailly, A. F. Volkov, and B. Pannetier, Phys. Rev. Lett. 77, 4950 (1996).
[6] C. J. Chien and V. Chandrasekhar, Phys. Rev. B 60, 3655 (1999).
[7] C. J. Chien and V. Chandrasekhar, Phys. Rev. B 60, 15356 (1999).
[8] S. Shapira, E. H. Linfield, C. J. Lambert, R. Serviour, A. F. Volkov, and A. V. Zaitsev, Phys. Rev. Lett. 84, 159 (2000).
[9] A. A. Golubov, F. K. Wilhelm, and A. D. Zaikin, Phys. Rev. B 55, 1123 (1997).
[10] A. F. Volkov and V. V. Pavlovskii, in *Proceedings of the XXXI Rencontres de Moriond*, edited by T. Martin, G. Montambaux, and J. Tran Thanh Van (Frontiers, France, 1996).

[11] B. Z. Spivak and D. E. Khmel’nitskii, JETP Lett. **35**, 412 (1982).

[12] V. T. Petrashov, V. N. Antonov, P. Delsing, and T. Claeson, Phys. Rev. Lett. **74**, 5268 (1995);
    V. T. Petrashov, V. N. Antonov, P. Delsing, and T. Claeson, Phys. Rev. Lett. **70**, 347 (1993).

[13] H. Pothier, S. Gueron, D. Esteve, and M. M. Devoret, Phys. Rev. Lett. **73**, 2488 (1994).

[14] P. G. N. de Vegvar, T. A. Fulton, W. H. Mallison, and R. E. Miller, Phys. Rev. Lett. **73**, 1416 (1994).

[15] J. Nitta, T. Akazaki, and H. Takayanaki, Phys. Rev. B **49**, 3659 (1994).

[16] F. W. J. Hekking and Y. V. Nazarov, Phys. Rev. Lett. **71**, 1625 (1993); F. W. J. Hekking and
    Y. V. Nazarov, Phys. Rev. B **49**, 6847 (1994).

[17] A. V. Zaitsev, Phys. Lett. A **194**, 315 (1994).

[18] B. J. Van Wees, K.-M. Lenssen, and C. J. P. M. Harmans, Phys. Rev. B **44**, 470 (1991).

[19] A. F. Volkov, Phys. Rev. Lett. **74**, 4730 (1995).

[20] S. K. Yip, Phys. Rev. B **58**, 5803 (1998).

[21] F. K. Wilhelm, G. Schön, and A. D. Zaikin, Phys. Rev. Lett. **81**, 1682 (1998).

[22] G. Wendin and V. S. Sumeiko, Phys. Rev. B **53**, R6006 (1996).

[23] J. J. Baselmans, A. F. Morpurgo, T. M. Klapwijk, and B. J. van Wees, Nature **43**, 43 (1999).

[24] V. N. Gubankov and N. M. Margolin, JETP Lett. **29**, 673 (1979).

[25] S. N. Artemenko, A. Volkov, and A. V. Zaitsev, Sol. St. Comm. **30**, 771 (1979).

[26] Y. V. Nazarov and T. H. Stoof, Phys. Rev. Lett. **76**, 823 (1996).
[27] A. F. Volkov, N. Allsopp, and C. J. Lambert, J. Phys. Condens Matter 8, 45 (1996).

[28] A. F. Volkov and V. V. Pavlovskii, JETP Lett. 64, 670 (1996).

[29] J. Kutchinsky, R. Taboryski, C. B. Sorensen, J. B. Jonsen, and P. E. Lindelof, Phys. Rev. Lett. 83, 4856 (1999).

[30] V. T. Petrasov, I. A. Sosnin, I. Cox, A. Parsons, and C. Troadec, Phys. Rev. Lett. 83, 3281 (1999).

[31] J. Aumentado and V. Chandrasekhar, Phys. Rev. B 64, 054505 (2001).

[32] M. Giroud, H. Courtois, K. Hasselbach, D. Mailly, and B. Pannetier, Phys. Rev. B 58, 11872 (1998).

[33] M. J. M. de Jong and C. W. J. Beenakker, Phys. Rev. Lett. 74, 1657 (1994).

[34] A. A. Golubov, Physica C 326-327, 46 (1999).

[35] W. Belzig, A. Brataas, Y. V. Nazarov, and G. E. W. Bauer, Phys. Rev. B 62, 9726 (2000).

[36] F. J. Jedema, B. J. van Wees, B. H. Hoving, A. T. Filip, and T. M. Klapwijk, Phys. Rev. B 60, 16549 (1999).

[37] V. I. Falko, A. F. Volkov, and C. J. Lambert, Phys. Rev. B 60, 15394 (1999); V. I. Falko, A. F. Volkov, and C. J. Lambert, JETP Lett. 69, 532 (1999); F. Taddei, J. H. Jefferson, and C. J. Lambert, Phys. Rev. Lett. 82, 4938 (1999).

[38] A. I. Larkin and Y. N. Ovchinnikov, in Nonequilibrium Superconductivity, edited by D. N. Langenberg and A. I. Larkin (Elsevier, Amsterdam, 1984).

[39] A. F. Volkov, A. V. Zaitsev, and T. M. Klapwijk, Physica C 210, 21 (1993).

[40] M. Y. Kuprianov and V. F. Lukichev, Sov. Phys. JETP 64, 139 (1988).
[41] A. Kastalsky, A. W. Kleinsasser, L. H. Greene, R. Bhat, F. P. Milliken, and J. P. Harbison, Phys. Rev. Lett. 67, 3026 (1991).

[42] B. J. van Wees, P. de Vries, P. Magnee, and T. M. Klapwijk, Phys. Rev. Lett. 69, 510 (1992).

[43] A. F. Volkov, JETP Lett. 55, 746 (1992).

[44] K. L. Usadel, Phys. Rev. Lett. 25, 507 (1970).

[45] A. V. Zaitsev, JETP 59, 863 (1984).

[46] A. I. Buzdin and M. Y. Kupriyanov, JETP Lett. 53, 321 (1991).

[47] A. F. Volkov and H. Takayanagi, Phys. Rev. B 56, 11184 (1997).

[48] F. S. Bergeret, A. F. Volkov, and K. B. Efetov, Phys. Rev. Lett. 86, 4096 (2001).

[49] V. L. Berezinskii, JETP Lett. 20, 287 (1975).

[50] A. Balatsky and E. Abrahams, Phys. Rev. B 45, 13125 (1992).

[51] A. I. Larkin, JETP Lett. 2, 130 (1965).

[52] F. Zhou, B. Spivak, and A. Zyuzin, Phys. Rev. B 52, 4467 (1995).

[53] H. Courtois, Ph. Gandit, D. Mailly, and B. Pannetier, Phys. Rev. Lett. 76, 130 (1996).

[54] F. Wilhelm, A. D. Zaikin, and H. Courtois, Phys. Rev. Lett. 80, 4289 (1998).
Figure Captions

Fig.1: S/F system.

Fig.2: The dependence of the conductance variation $\delta S(T) = S(T) - 1$, for different values of the exchange field $h = h_{ex}/E_{Th}$. Here $\Delta/E_{Th} = 10$, $r_0 = 1.10^{-3}$ and $r_L = 0.1$.

Fig.3: The dependence of the conductance variation $\delta S(T) = S(T) - 1$, for different values of the exchange field $h = h_{ex}/E_{Th}$. Here $\Delta/E_{Th} = 10$, $r_0 = 1.10^{-3}$ and $r_L = 2$. The solid, dashed and dotted lines correspond to $h = 50$, $h = 100$ and $h = 500$ respectively.

Fig.4: Schematic view of the structures under consideration.

Fig.5: Spatial dependence of the singlet (dashed line) and the triplet (solid line) components of $|\hat{f}|$ in the F wire for different values of $\alpha_w$. Here $w = L/5$, $\epsilon = E_T$ and $h/E_T = 400$. $E_T = D/L^2$ is the Thouless energy.

Fig.6: The $\delta S(T)$ dependence. Here $\gamma = \rho\xi_h/R_b$. $\Delta/E_T \gg 1$ and $w/L = 0.05$. 
Figure 1: F. S. Bergeret, A.F. Volkov and K.B. Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”
Figure 2: F. S. Bergeret, A.F. Volkov and K.B. Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”
Figure 3: F. S. Bergeret, A.F. Volkov and K.B.Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”
Figure 4: F. S. Bergeret, A.F. Volkov and K.B. Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”
$\alpha_w = \pi/5$

$\alpha_w = \pi$

Figure 5: F. S. Bergeret, A.F. Volkov and K.B. Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”
Figure 6: F. S. Bergeret, A.F. Volkov and K.B.Efetov. “Transport and triplet superconducting condensate in mesoscopic ferromagnet-superconductor structures.”