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Superconductivity and magnetism/Supraconductivité et magnétisme

Odd triplet superconductivity in superconductor–ferromagnet hybrid structures

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Abstract

We study the proximity effect in diffusive superconductor–ferromagnet structures. In particular we show that a triplet component of the condensate appears. This component is an odd function of the Matsubara frequency and an even function of the momentum. Therefore it is insensitive to the scattering by non-magnetic impurities. If the exchange field of the ferromagnet is homogeneous in space, only the triplet component with spin projection $S_z = 0$ appears. The latter penetrates into the ferromagnet over the short length $\xi_h = \sqrt{D/h}$ (D is the diffusion coefficient and h the exchange field) like the singlet component. However, if the exchange is not homogeneous the condensate has also the components with projections $S_z = \pm 1$. These are long-range components which penetrate the ferromagnet over the $\xi_T = \sqrt{D/T}$ (T is the temperature). We propose some experiments in order to detect the odd triplet component. ***To cite this article: F.S. Bergeret et al., C. R. Physique 7 (2006).***

Résumé

Superconductivité triplet impaire dans les structures hybrides supraconducteur–ferromagnétique. Nous étudions l'effet de proximité dans des structures supraconducteur–ferromagnétique diffusives. Nous montrons en particulier qu'une composante triplet du condensat apparaît. Cette composante est fonction impaire de la fréquence de Matsubara et fonction paire du moment. Elle est donc insensible à la diffusion par les impuretés non magnétiques. Si le champ d'échange du matériau ferromagnétque est spatialement uniforme, seule apparaît la composante triplet de spin projeté $S_z = 0$. Elle pénètre dans le ferromagnétique sur la courte distance $\xi_h = \sqrt{D/h}$ (où D est le coefficient de diffusion et h le champ d'échange), tout comme la composante singulet. Cependant, si l'échange n'est pas uniforme, le condensat possède aussi les composantes avec projections $S_z = \pm 1$. Ces dernières sont des composantes à longue portée qui pénètrent le ferromagnétique sur la distance $\xi_T = \sqrt{D/T}$ (où T est la température). Nous proposons plusieurs expériences destinées à détecter la composante triplet impaire. ***Pour citer cet article : F.S. Bergeret et al., C. R. Physique 7 (2006).***

Keywords: Superconductivity; Magnetism; Proximity effect

Mots-clés : Superconductivité ; Magnétism ; Effet de proximité

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1. Introduction

The study of the interplay between superconductivity and ferromagnetism in superconductor–ferromagnet (S/F) hybrid structures has attracted a great attention of researches (for review see [1–3] and references therein). As it is nowadays well known the proximity effect in such structures leads to the penetration of the superconducting condensate function into the F region over a distance of the order of $\xi_F \sim \sqrt{D/h}$, where h is the exchange field in the ferromagnet and D the diffusion coefficient (we consider the so called dirty limit when the mean free path l is the shortest length involved in the problem except for the Fermi wave length). Notice that for a strong ferromagnet with a large exchange field this length is very small (few Angstroms). The exponential decay of the condensate function into the ferromagnet is accompanied by oscillations in space. These oscillations lead, for example, to oscillations of the critical superconducting temperature T_c and the critical Josephson current I_c in S/F structures as a function of the thickness d_F. Such oscillatory behaviour, being predicted in Refs. [4,5], was observed in many works [6–11] on different S/F structures. Another manifestation of the oscillatory behavior of the condensate in the ferromagnet is the change of sign of the Josephson critical current through a $S/F/S$ junction (π-junction) [12]. This effect has been observed recently in many experiments [13–18].

The effects mentioned above were studied under the assumption that the exchange in the ferromagnet is homogeneous in space. In this case the condensate has a simple spin structure. However, it turns out that a non-homogeneous alignment of the exchange field leads to a complicated spin structure of the superconducting condensate [19,20]. As a result, not only the singlet component of the condensate exists but also a triplet one with all projections of the total spin of the Cooper pair ($S_z = 0, \pm 1$). In contrast to the singlet component, the spins of the electrons in the triplet state with $S = \pm 1$ are parallel to each other. Therefore, when generated, this triplet component is not destroyed by the exchange field and can penetrate the ferromagnet over long distances of the order of $\xi_N = \sqrt{D_F/2\pi T}$.

We will also show that the condensate function of the triplet state is an odd function of the Matsubara frequency and therefore this type of superconductivity is called *odd triplet superconductivity*. The singlet part is, as usual, an even function of ω but it changes sign when interchanging the spin indices. Both components are insensitive to the scattering by non-magnetic impurities and hence survive in the S/F structures even if the mean free path l is short.

The aim of this article is to review the properties and effects related to the odd triplet component (TC) of the condensate in S/F structures. In the first part we will study the structure of the condensate in the spin space and describe how to distinguish the triplet component from the usual BCS singlet one. Afterwards we will discuss the manifestation how the TC manifests itself in different physical properties of the system. Our analysis is based on the microscopic formalism for quasiclassical Green's functions (GF) [21,22]. Throughout this article the dirty limit is assumed.

2. The condensate function in a S/F structure

Consider a dirty S/F structure. Assuming that all energies are smaller than the Fermi energy and that all quantities vary over length scales larger than the Fermi wave length we can use the quasiclassical approach. In the dirty limit the GFs satisfy the Usadel equation [23], which in a 4×4 notation has the form [24]

$$D\partial(\check{g}\partial\check{g}/\partial r)/\partial r - \mathrm{i}\omega_n[\hat{\tau}_3, \check{g}] + [\check{\Delta}, \check{g}] + [\mathbf{h}\check{\mathbf{S}}, \check{g}] = 0 \tag{1}$$

where ω_n are the Matsubara frequencies, $\check{\Delta} = \mathrm{i}\Delta\hat{\tau}_2\hat{\sigma}_3$ is a matrix of the order parameter which is not zero only in the superconductor, $\check{\mathbf{S}}$ is the matrix vector $\check{\mathbf{S}} = (\hat{\sigma}_1, \hat{\sigma}_2, \hat{\tau}_3\hat{\sigma}_3)$, $\hat{\tau}_i$ and $\hat{\sigma}_i$, $i = 1, 2, 3$, are the Pauli matrices in particle-hole and spin space 'hat' 2×2 matrices. The GF is the sum of the normal component g and the condensate function f

$$\check{g} = \hat{g}\hat{\tau}_3 + \hat{f}\mathrm{i}\hat{\tau}_2$$

Solutions for the Usadel equation satisfy the normalization condition

$$\check{g}^2 = 1 \tag{2}$$

The Usadel equation is complemented by the boundary conditions presented in [25] on the basis of the Zaitsev's boundary conditions [26]. In the absence of spin-flip processes at the interface they take the form:

$$\check{g}_1\partial_x\check{g}_1 = \frac{1}{2\gamma_a}[\check{g}_1, \check{g}_2] \tag{3}$$

where $\gamma_1 = R_b\sigma_1$, σ_1 is the conductivity of the conductor 1 and R_b is the interface resistance per unit area, the x-coordinate is assumed to be normal to the plane of the interface.

First we consider the case of a homogeneous magnetization $\mathbf{h} = h(0, 0, 1)$. Assuming a weak proximity effect one can linearize Eq. (1) and obtain a simple equation for the condensate function $\hat{f}$

$$\left(\partial^2 \hat{f}/\partial x^2\right) - \kappa_\omega^2 \hat{f} + \mathrm{i}\kappa_h^2 \sigma_3 \hat{f} = 0 \tag{4}$$

where $\kappa_\omega^2 = 2|\omega|/D_F$ and $\kappa_h^2 = 2h \operatorname{sgn}\omega/D_F$. One can easily check that the condensate functions obeying Eq. (4) has the following form in spin space

$$\hat{f} = f_0\hat{\sigma}_0 + f_3\hat{\sigma}_3 \tag{5}$$

The amplitudes of the singlet and triplet components are related to the correlation functions $\langle\psi_\alpha\psi_\beta\rangle$ as follows [27]

$$f_3(t) \sim \left\langle\psi_\uparrow(t)\psi_\downarrow(0)\right\rangle - \left\langle\psi_\downarrow(t)\psi_\uparrow(0)\right\rangle$$
$$f_0(t) \sim \left\langle\psi_\uparrow(t)\psi_\downarrow(0)\right\rangle + \left\langle\psi_\downarrow(t)\psi_\uparrow(0)\right\rangle \tag{6}$$

Thus, the amplitude f_3 corresponds to the singlet component, while f_0 corresponds to the amplitude of the triplet component with zero projection of the magnetic moment of Cooper pairs on the z axis ($S_z = 0$). In the case of a non-ferromagnetic normal metal the condensate function has a singlet structure only, i.e., it is proportional to $\hat{\sigma}_3$. The exchange field breaks the spin symmetry and leads to the appearance of the triplet term proportional to $\hat{\sigma}_0$. This is an important point, since once one has understood the structure of the condensate (Eq. (5)), one can guess that a nonhomogeneous exchange field might lead to the appearance of the other two projections ($S_z = \pm 1$) of the triplet component. This will be discussed in the next section.

3. Triplet component of the condensate in $F/S/F$ structure

We consider the $F/S/F$ structure of Fig. 1. The system consists of one S layer and two F layers with magnetizations inclined at the angle $\pm\alpha$ with respect to the z-axis (in the yz plane). From the above analysis we know that each F layer generates the triplet component with projection $S_z = 0$, on the direction of the exchange field. If the magnetic moments of the layers are collinear (parallel $\alpha = 0$ or antiparallel $\alpha = \pi/2$), the total projection S_z remains equal to zero. However, if the moments of the ferromagnetic layers are not collinear the superposition of both triplet components may lead to a triplet component of the condensate with all possible projections, i.e., $S_z = 0, \pm 1$. Let us analyze this case.

The linearized Usadel equation in the F layer is now given by

$$\left(\partial^2 \check{f}_F/\partial x^2\right) - \kappa_\omega^2 \check{f}_F + \mathrm{i}\kappa_h^2 [\hat{\sigma}_3, \check{f}_F]_+ \cos\alpha \pm \hat{\tau}_3[\hat{\sigma}_2, \check{f}_F] \sin\alpha = 0 \tag{7}$$

The sign '+' ('−') corresponds to the right (left) F film.

The linearization may be justified in the two limiting cases: (a) T is close to the critical temperature of the structures T_c^* (the latter can be different from the critical temperature of the bulk superconductor T_c); and (b) the resistance of the S/F interface R_b is not small. In the latter case the condensate function in the S layer is weakly disturbed by the F film and the function

$$\check{f}_S = \check{f}_{BCS} + \delta\check{f}_S \tag{8}$$

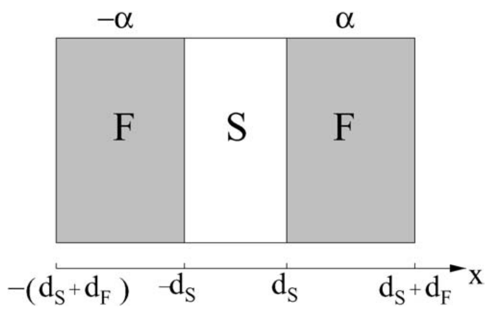

Fig. 1. Trilayer geometry. The magnetization of the left (right) F layer makes an angle α ($-\alpha$) with the z-axis.

where $\hat{f}_{BCS} = \mathrm{i}\hat{\tau}_2\hat{\sigma}_3\Delta/\mathrm{i}E_\omega$ is the bulk solution, $E_\omega = \sqrt{\omega^2+\Delta^2}$ and $\delta\check{g}_S$ a small correction. Thus, the linearized Usadel equation for $\delta\check{f}_S$ has the simple form

$$\partial^2_{xx}\delta\check{f}_S - \kappa_S^2\delta\check{f}_S = 0 \tag{9}$$

Here $\kappa_S^2 = 2E_\omega/D_S$.

It turns out that the form of the condensate function in the F layer is given by

$$\check{f}_F = \hat{f}_1\mathrm{i}\hat{\tau}_1 + \hat{f}_2\mathrm{i}\hat{\tau}_2 \tag{10}$$

The functions $\hat{f}_i$ are matrices in the spin-space. They can be represented in the form

$$\hat{f}_2(x) = f_0(x)\hat{\sigma}_0 + f_3(x)\hat{\sigma}_3 \tag{11}$$

$$\hat{f}_1(x) = f_1(x)\hat{\sigma}_1 \tag{12}$$

We see that in this case the condensate has an additional term, $\hat{f}_1$ which is related to the correlation functions $\langle\psi_\uparrow\psi_\uparrow\rangle$, $\langle\psi_\downarrow\psi_\downarrow\rangle$, and therefore corresponds to the triplet component with projection $S_z = \pm 1$. Notice that in the case of a homogeneous magnetization this component is absent (cf. Eq. (5)).

Solutions for Eqs. (7)–(9) can be written as a sum of exponential functions $\exp(\pm\kappa x)$, where the κ's are the eigenvalues of these equations. In the S layer the equations for $\delta f_3, f_{0,1}$ are decoupled and there is only one eigenvalue $\kappa = \kappa_S$. In the F-layers the equations are coupled and there are three different eigenvalues [28]

$$\kappa_{1,2} \equiv \kappa_\pm \simeq \sqrt{h/D_F}(1\pm\mathrm{i}) \tag{13}$$

$$\kappa_3 \equiv \kappa_\omega = \sqrt{2|\omega|/D_F} \tag{14}$$

It is assumed that $h \gg T$. We see from these equations that two completely different lengths $\xi_h \sim \sqrt{D_F/h}$ and $\xi_T \sim \sqrt{D_F/T}$ determine the decay of the condensate in the F layers. Usually $h > T_C$ and therefore at all temperatures $T < T_c^*$ the length ξ_T exceeds ξ_h. As expected, the triplet component with spin projection $S_z = \pm 1$ is long-ranged. Its decaying length is the same as the decay of the standard singlet condensate in a normal metal.

The solution in the right F layers for the condensate function is (for the detailed calculation see [29])

$$f_1(x) = b_1\frac{\cosh\kappa_\omega(x-d_S-d_F)}{\cosh(\kappa_\omega d_F)} + \operatorname{sgn}\omega\sin\alpha\left[-b_{3+}\mathrm{e}^{\kappa_+(x-d_s)} + b_{3-}\mathrm{e}^{-\kappa_-(x-d_s)}\right] \tag{15}$$

$$f_0(x) = -\tan\alpha\, b_1\frac{\cosh\kappa_\omega(x-d_S-d_F)}{\cosh(\kappa_\omega d_F)} + \operatorname{sgn}\omega\cos\alpha\left[-b_{3+}\mathrm{e}^{-\kappa_+(x-d_s)} + b_{3-}\mathrm{e}^{-\kappa_-(x-d_s)}\right] \tag{16}$$

$$f_3(x) = b_{3+}\mathrm{e}^{-\kappa_+(x-d_s)} + b_{3-}\mathrm{e}^{-\kappa_-(x-d_s)} \tag{17}$$

where

$$\tilde{b}_{3\pm} = b_{3\pm}(g_s + \gamma_b\xi_h\kappa_\pm) = f_{BCS}\frac{\tilde{\kappa}_S\tanh\Theta_S M_\mp}{M_+T_- + M_-T_+} \tag{18}$$

$$\tilde{b}_1 = b_1(g_{BCS} + \gamma_b\xi_h\kappa_\omega\tanh\Theta_F) = -f_{BCS}\sin\alpha\frac{\tilde{\kappa}_S^2(\tilde{\kappa}_+ - \tilde{\kappa}_-)\operatorname{sgn}\omega}{\cosh^2\Theta_S(M_+T_- + M_-T_+)} \tag{19}$$

and $\Theta_S = \kappa_s d_S$, $\Theta_F = \kappa_\omega d_F$, $\tilde{\kappa}_\pm = \kappa_\pm/(g_{BCS} + \gamma_b\xi_h\kappa_\pm)$, $\tilde{\kappa} = \kappa_\omega/(g_{BCS} + \gamma_b\xi_h\kappa_\omega\tanh\Theta_F)$, $\tilde{\kappa}_S = \kappa_S/(g_{BCS}\gamma)$, $g_{BCS} = \omega/E_\omega$, and

$$M_\pm = T_\pm(\tilde{\kappa}_S\coth\Theta_S + \tilde{\kappa}\tanh\Theta_F) + \tan^2\alpha, C_\pm(\tilde{\kappa}_S\tanh\Theta_S + \tilde{\kappa}\tanh\Theta_F)$$

$$T_\pm = \tilde{\kappa}_S\tanh\Theta_S + \tilde{\kappa}_\pm, \quad C_\pm = \tilde{\kappa}_S\coth\Theta_S + \tilde{\kappa}_\pm$$

The solutions in the left F layer can be easily obtained recalling that the function $f_1(x)$ is odd and $f_{0,3}(x)$ are even functions of x.

Thus, we see that the singlet component f_3 is an even function of ω and decays sharply in the ferromagnet over the short distance ξ_h. In contrast, the amplitudes of the triplet component f_0 and f_1 are odd functions of ω and penetrate the ferromagnet over the longer distance $\xi_T = \sqrt{D_F/2\pi T}$. The long-range part of TC determined by the amplitude b_1 has the maximum at $\alpha = \pi/4$. This value of α corresponds to a perpendicular orientation of the magnetizations in

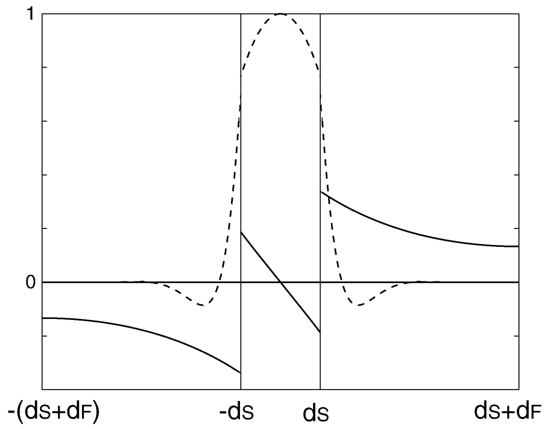

Fig. 2. The spatial dependence of Im(SC) (dashed line) and the long-range part of Re(TC) (solid line). We have chosen $\sigma_F/\sigma_S = 0.2$, $h/T_C = 50$, $\sigma_F R_b/\xi_F = 0.05$, $d_F\sqrt{T_C/D_S} = 2$, $d_S\sqrt{T_C/D_S} = 0.4$ and $\alpha = \pi/4$. The discontinuity of the TC at the S/F interface is because the short-range part is not shown in this figure. Taken from [29].

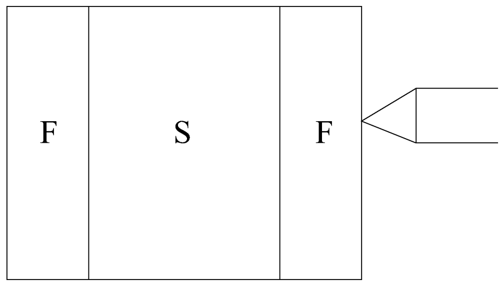

Fig. 3. Schematic: Measurement of the change of the density of states at the outer F interface by tunnelling spectroscopy. [14] performed such experiments on S/F structures.

the F layers. For a parallel ($\alpha = 0$) or antiparallel alignment of the magnetizations ($\alpha = \pi/2$) this amplitude decays to zero. In Fig. 2 we plot the spatial dependence of the SC and the long-range part of the TC. We see that both amplitudes are comparable at the S/F interface but the SC decays faster than the TC.

Finally, it is worth mentioning that in [19,20] the conductance of a ferromagnetic wire attached to a superconductor was calculated. It was assumed that the F wire had a domain wall located at the S/F interface. This inhomogeneity of the magnetization induces the long-range triplet component studied here. The latter leads to an increase of the conductance for temperatures below T_c as observed in the experiments of [31,30]

4. How to detect the triplet component of the condensate?

In order to verify the existence of the triplet component it is necessary to separate it from the singlet one. Some experimental realizations were presented in [28,29].

For example, one can consider the system shown in Fig. 3. It is clear that, for distances from the S/F interface larger than ξ_h, only the TC leads to a variation of the local DoS. Thus, if the thickness d_F is much larger than ξ_h one can detect directly the presence of the TC performing measurements of the DoS at the outer side of one of the F layers. Any deviation from the normal value would be only due to the TC. Similar measurements as those of [14] can be performed. In this experiments the DoS was determined with the help of a planar tunneling spectroscopy.

From the results of the previous section it is easy to determine the DoS at $x = d_S + d_F$. The expression for the normalized DoS is (we ignore the difference in the DoS for the up and down spin directions in the ferromagnet above the critical temperature T_c)

$$\delta\nu = \frac{1}{8}\operatorname{Tr}(\hat{\tau}_3\hat{\sigma}_0)\big(\check{g}^{\mathrm{R}} - \check{g}^{\mathrm{A}}\big) \tag{20}$$

Here the indices ‘A’ and ‘R’ denote advanced and retarded. As it was mentioned before, in the case $d_F \gg \xi_h$ only the TC (i.e., the functions $f_0(x)$ and $f_1(x)$) contributes to the DoS. For the DoS one obtains [29]

$$\delta\nu = \frac{1}{2}\operatorname{Re}\frac{(b_1^{\mathrm{R}})^2}{\cos^2\alpha\cosh^2\Theta_F^{\mathrm{R}}} \tag{21}$$

where $\Theta_F^{\mathrm{R}} = \sqrt{-2\mathrm{i}\varepsilon/D_F}\,d_F$, and b_1^{R} is the amplitude of the retarded Green’s function in Eqs. (15)–(16). It is obtained from b_1 by replacing ω by $-\mathrm{i}\varepsilon$. In Fig. 4 the dependence of $\delta\nu$ on ε for different α is shown.

Another system suitable for detecting the triplet component of the condensate is shown in Fig. 5. The idea is to measure the Josephson current between the S layers of the FSFSF structure. The thickness of the F layers d_F is

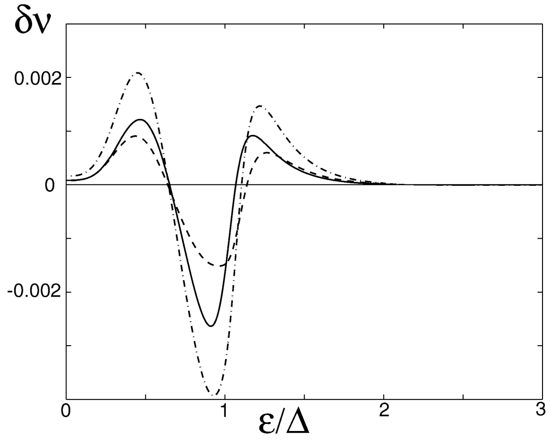

Fig. 4. The normalized DoS $\delta\nu$ as a function of the energy for $\alpha = 3\pi/8$ (solid line), $\alpha = \pi/8$ (dashed line) and $\alpha = \pi/4$ (point-dashed line). Note that for $\alpha = 0, \pi/2$ $\delta\nu = 0$. We have chosen $\sigma_F/\sigma_S = 0.05$, $h/\Delta = 25$, $\sigma_F R_b/\xi_h = 0.5$, $d_F/\xi_\Delta = 0.5$, and $d_S/\xi_\Delta = 0.4$. Here $\xi_\Delta = \sqrt{D_S/\Delta}$ and Δ is the BCS order parameter. Taken from [29].

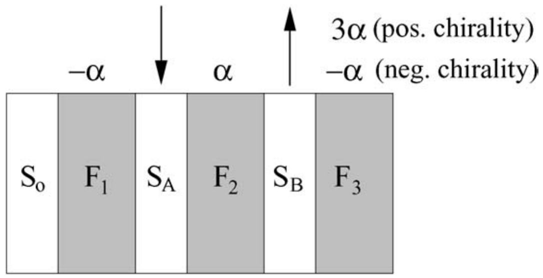

Fig. 5. The multilayered structure considered. The arrows show the bias current. In the case of positive (negative) chirality the magnetization vector $\mathbf{M}$ of the layer F_3 makes an angle 3α $(-\alpha)$ with the z-axis, i.e., in the case of positive chirality the vector M rotates in one direction if we go over from one F layer to another whereas it oscillates in space in the case of negative chirality. Taken from [29].

much larger than ξ_h. According to the previous sections it is clear that in this case the Josephson coupling between the S layers is due to the long range part of the TC. Therefore the supercurrent in the transverse direction is caused by the triplet component of the condensate, while the in-plane superconductivity is caused mainly by the ordinary singlet component. Thus, the macroscopic superconductivity due to the Josephson coupling between the layers is an interesting combination of the singlet superconductivity within the layers and the odd triplet superconductivity in the transversal direction. The critical Josephson current was calculated in [29,28] and is given by the expression

$$\mathrm{e}R_F I_c = \pm 2\pi T \sum_{\omega} \kappa_\omega d_F b_1^2(\alpha)\big(1 + \tan^2\alpha\big)\mathrm{e}^{-d_F\kappa_\omega}, \tag{22}$$

where $b_1(\alpha)$ is given in Eq. (19) and the sign '+' ('−') corresponds to the positive (negative) chirality. In the case of negative chirality the critical current is negative (π-contact). Thus, experimentally one can switch between the 0 and π-contacts by changing the angles of the mutual magnetization of the layers.

In Fig. 6 we plot the dependence of I_c on the angle α. If the orientation of $\mathbf{M}$ is parallel ($\alpha = 0$) or antiparallel ($\alpha = \pi/2$), the amplitude of the triplet component is zero and therefore there is no coupling between the neighboring S layers, i.e., $I_c = 0$. For any other angle between the magnetizations the amplitude of the triplet component is finite. This leads to a non-zero critical current. At $\alpha = \pi/4$ (perpendicular orientation of $\mathbf{M}$) I_c reaches its maximum value.

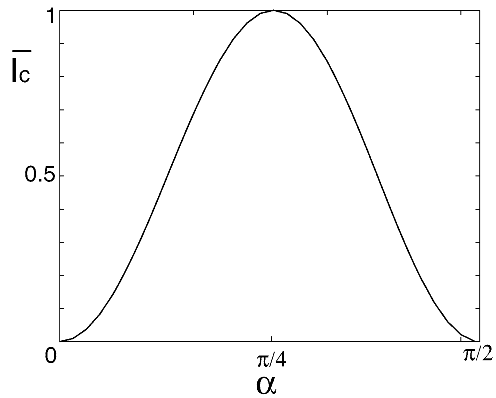

Fig. 6. Dependence of the critical current (normalized with respect to the maximum value) on the angle α. We have chosen the same values as in Fig. 4. Taken from [29].

Finally, we would like to mention that the influence of the triplet component on the variation of the critical temperature of a $F/S/F$ structure studied in [32].

5. Discussion and summary

We have shown that in S/F hybrid structures the appearance of a triplet component of the condensate is inevitable. If the system has only one quantization spin axis, i.e., an homogeneous exchange field, then only the triplet component with spin projection $S_z = 0$ appears. If the exchange field is not homogeneous then the other projections, $S_z = \pm 1$ also are generated. Obviously these are not affected by the exchange field and therefore the condensate may penetrate over large distance of the order of $\xi_T = \sqrt{D/T}$.

The triplet component of the condensate is an odd function of the Matsubara frequency and an even function of the momentum. The latter property make it insensitive to the scattering by non-magnetic impurities.

Historically, such odd triplet superconductivity was suggested by Berenziskii as a possible phase in superfluid ^{3}He [33]. However, this hypothetical condensate function is not realized. In ^{3}He the condensate function is odd in momenta, but not in frequency. The possibility to realize this type of superconductivity in bulk superconductors (mainly in high T_c materials)was discussed in [34–37].

Note an important fact. The quasiclassical Green's function in the diffusive case can be expanded in spherical harmonics. In the present approach only the first two terms of this expansion are taken into account, thus

$$\check{g} = \check{g}_0 + \check{g}_1 \cos\vartheta \tag{23}$$

where ϑ is the angle between the momentum p and the x-axis, and $\check{g}_1 = -l\check{g}_0 \partial \check{g}_0/\partial x$. The first term is symmetric respect to the momentum while the second term is antisymmetric. The antisymmetric part of $\check{g}$ determines the electric current in the system. Higher order terms in the expansion of $\check{g}$ are, in the diffusive limit, small and can be neglected. In the case of a weak proximity effect the antisymmetric part of the condensate function in the F region can be written as

$$\check{f}_1 \cos\vartheta \approx -l\hat{\tau}_3 \otimes \hat{\sigma}_0 \operatorname{sgn}\omega \partial \check{f}_0/\partial x \cos\vartheta \tag{24}$$

This expression follows from the fact that $\mathrm{e}\check{g}_0 \approx -l\hat{\tau}_3 \otimes \hat{\sigma}_0 \operatorname{sgn}\omega$ (corrections to $\check{g}_0$ are proportional to $\check{f}_0^2$). Eq. (24) holds for both the singlet and triplet component. Consider the triplet component. As we noted above, the symmetric part $\check{f}_0$ is an odd function of ω and therefore the antisymmetric part is an even function of ω so that the total condensate function $\check{f} = \check{f}_0 + \check{f}_1 \cos\vartheta$ is neither odd or even function of ω. However, in the diffusive limit it is valid to speak about odd superconductivity since the symmetric part is much larger than antisymmetric part of $\check{f}$. If the parameter $h\tau$ is not small, i.e., the system is not diffusive, the symmetric and antisymmetric parts are comparable, and one cannot say about odd superconductivity. All this distinguishes superconductivity in S/F structures from the odd superconductivity suggested by Berenziskii who assumed that the order parameter $\Delta(\omega)$ is an odd function of ω. In our discussion it is assumed that the order parameter Δ is an ω-independent quantity and it is determined by the singlet component of the condensate function $\check{f}_0$.

As to experimental studies, one has to note that convincing experimental evidences in favour of the odd long-range triplet component are still lacking. In [38–41] the conductance of a S/F mesoscopic structure was measured. It was found that the influence of the superconductor on the conductance spreads over a length exceeding essentially the exchange length ξ_h. This long-range effect may be related to the long-range triplet component [40,41].

In another experiment [42] a multilayered $S/F/S\ldots$ structure was studied. It consists of the high T_c material $YBa_2Cu_3O_7$ and the ferromagnetic half metal $La_{0.7}Ca_{0.3}MnO_3$. It was established that the superconducting coupling between neighboring S layers exists even if the thickness of the F layer exceeds the exchange length ξ_h. The authors assumed that a possible mechanism for this long-range effect may be again related with the triplet component. It may arise at the spin-active S/F interface [43]. However further experiments are needed in order to get more direct evidences which would confirm the existence of the odd long-range triplet component.

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