

Triplet proximity effect in FSF trilayers

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We study the critical temperature T_c of FSF trilayers (F is a ferromagnet, S is a singlet superconductor), where the triplet superconducting component is generated at noncollinear magnetizations of the F layers. An exact numerical method is employed to calculate T_c as a function of the trilayer parameters, in particular, mutual orientation of magnetizations. Analytically, we consider limiting cases. Our results determine conditions which are necessary for existence of recently investigated odd triplet superconductivity in SF multilayers.

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A striking feature of the proximity effect between singlet superconductors and nonhomogeneous ferromagnets is the possibility of generating the triplet superconducting component [1, 2]. Recently, it was shown that the triplet component also arises in the case of several homogeneous but differently oriented ferromagnets [3]. Physically, the generating of the triplet component in SF systems [1–3] is similar to the case of magnetic superconductors [4].

In Ref. [3], the Josephson effect was studied having in mind that the superconductivity in the system is not suppressed by the ferromagnets. However, this issue requires separate study.

Although the SF (F – ferromagnet, S – singlet superconductor) proximity effect is rather well studied, the influence of the *mutual orientation* of F layers magnetizations (exchange fields) on T_c of layered SF structures has been mostly considered basing on the cases of parallel (P) and antiparallel (AP) alignment [5–10]. At the same time, those are the only cases when the triplet component is absent.

A FSF trilayer with homogeneous but noncollinear magnetizations of the F layers is the simplest example of a layered structure in which the triplet component is generated. The triplet component (correlations between quasiparticles with parallel spins) arises as a result of interplay between the Andreev reflections at the two SF interfaces. This mechanism is similar to the one described in Ref. [2], with the difference that instead of local magnetic inhomogeneity we deal with magnetic inhomogeneity of the structure as a whole.

The critical temperature of the noncollinear FSF system was studied in Ref. [11]. However, in that work the triplet component was not taken into account. Thus calculation of T_c in the noncollinear FSF trilayer is still an open question.

In this letter we study the critical temperature of a FSF trilayer at arbitrary angle between the in-plane magnetizations (see Fig.1), which makes it necessary to take the triplet component into account. We reduce the problem to the form, which allows to apply general nu-

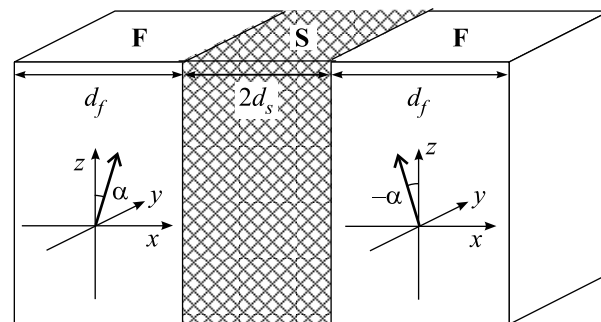


Fig.1. FSF trilayer. The system is the same as in Ref. [11]. The thickness of the S layer is $2d_s$, of each F layer — d_f . The center of the S layer corresponds to $x = 0$. The thick arrows in the F layers denote the exchange fields \mathbf{h} lying in the (y, z) plane. The angle between the in-plane exchange fields is 2α

merical methods developed in Refs. [12, 13]. This form also leads to some general conclusions about T_c and allows analytical progress in limiting cases.

1. General description. We consider the dirty limit, which is described by the Usadel equations. Near

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T_c , the Usadel equations are linearized and contain only the anomalous Green function \widehat{F} [1]:

$$\frac{D}{2} \frac{d^2 \widehat{F}}{dx^2} - |\omega_n| \widehat{F} + \Delta \widehat{\sigma}_3 - \frac{i}{2} \text{sgn} \omega_n (\widehat{F} \widehat{H}^* + \widehat{H} \widehat{F}) = 0, \quad (1)$$

$$\widehat{F} = \begin{pmatrix} f_{\uparrow\downarrow} & f_{\uparrow\uparrow} \\ f_{\downarrow\downarrow} & f_{\downarrow\uparrow} \end{pmatrix}.$$

Here D is the diffusion constant (D_s and D_f for the S and F layers), $\omega_n = \pi T(2n + 1)$ are the Matsubara frequencies, and $\widehat{\sigma}_3$ is the third Pauli matrix. The function \widehat{F} is a matrix in the spin space. The $f_{\uparrow\uparrow}$ and $f_{\downarrow\downarrow}$ components describe the triplet superconducting correlations. In the P and AP cases it is sufficient to consider only the scalar equation for the singlet component $f_{\uparrow\downarrow}$.

Equation (1) is written in the general case when both pair potential and exchange field are present. In our system, in the F layers the pair potential is absent, $\Delta = 0$, while

$$\widehat{H} = h(\widehat{\sigma}_2 \sin \alpha + \widehat{\sigma}_3 \cos \alpha) \quad (2)$$

at the exchange field $\mathbf{h} = h(0, \sin \alpha, \cos \alpha)$. h is the exchange energy, and α describes the direction of the in-plane magnetization.

In the S layer, the exchange energy is zero, while the pair potential obeys the self-consistency equation

$$\Delta \ln \frac{T_{cs}}{T} = \pi T \sum_{\omega_n} \left(\frac{\Delta}{|\omega_n|} - f_{\uparrow\downarrow} \right), \quad (3)$$

where T_{cs} is the critical temperature of the S material. In the case of a single S layer, Δ can be chosen real.

The boundary conditions at the outer surfaces of the trilayer are

$$d\widehat{F}_f/dx = 0, \quad (4)$$

while at the SF interfaces

$$\xi_s(d\widehat{F}_s/dx) = \gamma \xi_f(d\widehat{F}_f/dx), \quad \gamma = \rho_s \xi_s / \rho_f \xi_f, \quad (5)$$

$$\pm \xi_f \gamma_b(d\widehat{F}_f/dx) = \widehat{F}_s - \widehat{F}_f, \quad \gamma_b = R_b \mathcal{A} / \rho_f \xi_f. \quad (6)$$

Here $\xi_{s(f)} = \sqrt{D_{s(f)}/2\pi T_{cs}}$ and $\rho_{s(f)}$ are the coherence lengths and the normal state resistivities of the S and F metals, R_b is the total resistance of the SF boundary, and \mathcal{A} is its area. The \pm sign in the l.h.s. of Eq. (6) refers to the left and right SF interface, respectively. The above boundary conditions were derived for SN interfaces [14] (N is a normal metal); their use in the SF case is justified by the small parameter $h/E_F \ll 1$ (E_F is the Fermi energy).

We expand the Green function \widehat{F} in the basis of the Pauli matrices $\widehat{\sigma}_i$, $i = 1, 2, 3$, and the unity matrix $\widehat{\sigma}_0$. It can be shown that the solution has the form

$$\widehat{F} = f_0 \widehat{\sigma}_0 + f_1 \widehat{\sigma}_1 + f_3 \widehat{\sigma}_3. \quad (7)$$

The f_0 component is imaginary, while f_1 and f_3 are real. The relations $f_0(-\omega_n) = -f_0(\omega_n)$, $f_1(-\omega_n) = -f_1(\omega_n)$, $f_3(-\omega_n) = f_3(\omega_n)$ make it sufficient to consider only positive Matsubara frequencies.

The f_1 components describes a special type of triplet condensate [1, 3], odd in frequency [$f_1(-\omega_n) = -f_1(\omega_n)$] and even in momentum, which is similar to the one proposed by Berezinskii [15]. It is independence on the momentum direction that allows the triplet condensate to survive in the diffusive limit.

Equation (1) yields three coupled scalar equations (we consider $\omega_n > 0$):

$$\begin{aligned} \frac{D}{2} \frac{d^2 f_0}{dx^2} - \omega_n f_0 - i h f_3 \cos \alpha &= 0, \\ \frac{D}{2} \frac{d^2 f_1}{dx^2} - \omega_n f_1 + h f_3 \sin \alpha &= 0, \end{aligned} \quad (8)$$

$$\frac{D}{2} \frac{d^2 f_3}{dx^2} - \omega_n f_3 - i h f_0 \cos \alpha - h f_1 \sin \alpha + \Delta = 0.$$

Analyzing symmetries implied by Eqs. (8) and geometry of the system, we conclude that $f_0(x) = f_0(-x)$, $f_1(x) = -f_1(-x)$, $f_3(x) = f_3(-x)$. Thus we can consider only one half of the system, say $x < 0$, while the boundary conditions at $x = 0$ are

$$df_0/dx = 0, \quad f_1 = 0, \quad df_3/dx = 0. \quad (9)$$

Below we shall use the following wave vectors:

$$\begin{aligned} k_f &= \sqrt{2\omega_n/D_f}, & k_h &= \sqrt{h/D_f}, \\ \tilde{k}_h &= \sqrt{k_f^2 + 2ik_h^2}, & k_s &= \sqrt{2\omega_n/D_s}. \end{aligned} \quad (10)$$

The solution in the left F layer, satisfying the boundary condition (4), has the form

$$\begin{aligned} \widehat{F}_f &= C_1 (i\widehat{\sigma}_0 \sin \alpha + \widehat{\sigma}_1 \cos \alpha) \cosh [k_f (x + d_s + d_f)] + \\ &+ C_2 (\widehat{\sigma}_0 \cos \alpha + i\widehat{\sigma}_1 \sin \alpha + \widehat{\sigma}_3) \cosh [\tilde{k}_h (x + d_s + d_f)] + \\ &+ C_3 (\widehat{\sigma}_0 \cos \alpha + i\widehat{\sigma}_1 \sin \alpha - \widehat{\sigma}_3) \cosh [\tilde{k}_h^* (x + d_s + d_f)]. \end{aligned} \quad (11)$$

The matrix boundary condition (6) yields three scalar equations, which allow to express the coefficients C_1 , C_2 , C_3 in terms of the components f_0 , f_1 , f_3 of the Green function on the S side of the FS interface:

$$\begin{aligned} C_1 &= (-if_0 \sin \alpha + f_1 \cos \alpha) / (1 + \gamma_b A_f), \\ C_2 &= (f_0 \cos \alpha - if_1 \sin \alpha + f_3) / 2(1 + \gamma_b A_h), \\ C_3 &= (f_0 \cos \alpha - if_1 \sin \alpha - f_3) / 2(1 + \gamma_b A_h^*), \end{aligned} \quad (12)$$

where we have introduced the following notations:

$$A_f = k_f \xi_f \tanh(k_f d_f), \quad A_h = \tilde{k}_h \xi_f \tanh(\tilde{k}_h d_f), \\ V_f = \gamma A_f / (1 + \gamma_b A_f), \quad V_h = \gamma A_h / (1 + \gamma_b A_h). \quad (13)$$

Then the boundary condition (5) yields three scalar equations which entangle f_0 , f_1 , and f_3 . Thus the Green function of the F layer is eliminated, and we obtain equations for the S layer only. Moreover, we can proceed further, because in the S layer the unknown function $\Delta(x)$ only enters the equation for the f_3 component [see Eqs. (8)]. At the same time, taking boundary conditions (9) into account, we can write $f_0 = B_0 \cosh(k_s x)$, $f_1 = B_1 \sinh(k_s x)$. Excluding B_0 and B_1 , we arrive at the effective boundary condition for f_3 :

$$\xi_s (df_3/dx) = W f_3, \quad (14)$$

where

$$W = \text{Re } V_h + \frac{(\text{Im } V_h)^2}{k_s \xi_s A(\alpha) + \text{Re } V_h}, \quad (15)$$

and the angular dependence is determined by

$$A = \frac{k_s \xi_s \tanh(k_s d_s) + V_f [\sin^2 \alpha + \tanh^2(k_s d_s) \cos^2 \alpha]}{k_s \xi_s [\cos^2 \alpha + \tanh^2(k_s d_s) \sin^2 \alpha] + V_f \tanh(k_s d_s)}. \quad (16)$$

Effectively, we obtain the following problem:

$$\Delta \ln \frac{T_{cs}}{T} = 2\pi T \sum_{\omega_n > 0} \left(\frac{\Delta}{\omega_n} - f_3 \right), \quad (17)$$

$$\frac{D_s}{2} \frac{d^2 f_3}{dx^2} - \omega_n f_3 + \Delta = 0, \quad (18)$$

$$\xi_s \frac{df_3(-d_s)}{dx} = W(\omega_n) f_3(-d_s), \quad \frac{df_3(0)}{dx} = 0 \quad (19)$$

— this is exactly the problem that was solved in Refs. [12, 13]. Inserting the new function W , we can use the methods developed in those works. At $\alpha = 0$, Eq. (15) reproduces W from Refs. [12, 13].

All information about the F layers is contained in a single function W , all information about the misorientation angle — in its part $A(\alpha)$. Knowledge of W is already sufficient to draw several general conclusions about the behavior of T_c . First, if the S layer is thick, i.e. $d_s \gg \xi_s$, then $\tanh(k_s d_s) \approx 1$ at characteristic frequencies, and T_c does not depend on α . Qualitatively, this happens because the effect of mutual orientation of the F layers is due to “interaction” between the two SF interfaces, which is efficient only in the case of thin S layer. Second, T_c does not depend on d_f if $d_f \gg \max(\xi_f, k_h^{-1})$. Qualitatively, this is due to the fact

that quasiparticles penetrate from S to F layer only on the scale $\max(\xi_f, k_h^{-1})$.

The triplet component is “nonmonotonic” as a function of α : it vanishes at $\alpha = 0$ and $\alpha = \pi/2$ (P and AP case, respectively), and arises only between the two boundary values. However, the $T_c(\alpha)$ dependence is always monotonic. It can be directly proven from the monotonic behavior of $A(\alpha)$, and, hence, W . This rigorously derived conclusion disproves the result obtained by the approximate single-mode method in Ref. [7], where it was claimed that T_c in the AP configuration can be smaller than in the P case.

Numerical results obtained by the methods developed in Refs. [12, 13], are shown in Figs. 2, 3. A question arises: why is there pronounced angular dependence in

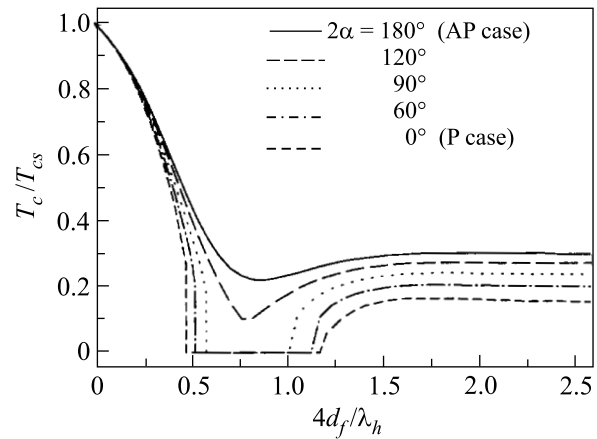


Fig. 2. Critical temperature T_c vs. thickness of the F layers d_f , which is normalized on the wavelength of the singlet component oscillations $\lambda_h = 2\pi/k_h$. Parameters $d_s/\xi_s = 1.2$, $h/\pi T_{cs} = 6.8$, $\gamma = 0.15$, $\gamma_b = 0.02$ correspond to Ref. [13]. The curves are calculated at different angles 2α between the in-plane exchange fields in the F layers

the case $d_s > \xi_s$, when the S layer is not thin? The answer is that the condition $d_s \ll \xi_s = \sqrt{D_s/2\pi T_{cs}}$ is a *sufficient* condition of thin S layer, whereas the *necessary* condition is weaker: $d_s \ll \xi = \sqrt{D_s/2\pi T_c}$, since the characteristic energy for a particular system is πT_c with its own value of T_c . The two conditions become essentially different if T_c is notably suppressed, and in this case T_c can exhibit pronounced angular dependence at $d_s \ll \xi$, while it is possible to have $d_s > \xi_s$.

Experimentally, the conditions for observing the angular dependence of T_c are more easily met when T_c is essentially (but not completely) suppressed. Accordingly, the effect of α on $T_c(d_f)$ dependence is most pronounced near the reentrant behavior. Experimental detection of such behavior was reported in Ref. [16].

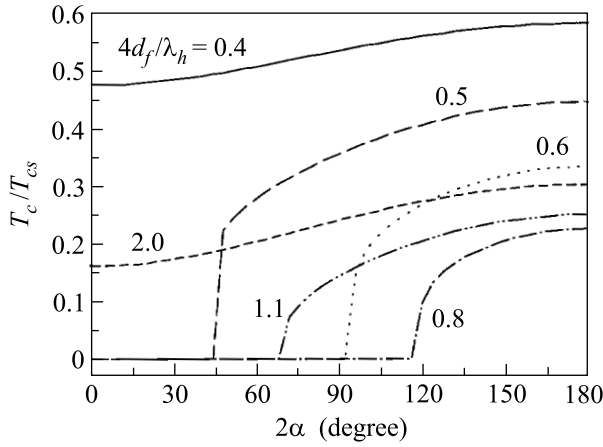


Fig.3. T_c vs. misorientation angle 2α . The curves correspond to different thicknesses of the F layers d_f . The parameters are the same as in Fig.2

2. Thin S layer. If $d_s \ll \xi_s$, then Δ is constant. The Usadel equation (18) can be solved, and the equation determining T_c takes the form

$$\ln \frac{T_{cs}}{T_c} = 2\pi T_c \sum_{\omega_n > 0} \left(\frac{1}{\omega_n} - \frac{1}{\omega_n + W\pi T_{cs} \xi_s / d_s} \right), \quad (20)$$

where W is given by Eq. (15) with simplified function $A(\alpha)$:

$$A = \frac{k_s^2 \xi_s d_s + V_f [\sin^2 \alpha + (k_s d_s)^2 \cos^2 \alpha]}{k_s \xi_s [\cos^2 \alpha + (k_s d_s)^2 \sin^2 \alpha] + V_f k_s d_s}. \quad (21)$$

For the P and AP alignments, under additional assumption of strong ferromagnetism ($h \gg \pi T_{cs}$), we obtain:

$$\ln \frac{T_{cs}}{T_c^P} = \text{Re} \psi \left(\frac{1}{2} + \frac{V_h \xi_s T_{cs}}{2 d_s T_c^P} \right) - \psi \left(\frac{1}{2} \right), \quad (22)$$

$$\ln \frac{T_{cs}}{T_c^{AP}} = \psi \left(\frac{1}{2} + \frac{W \xi_s T_{cs}}{2 d_s T_c^{AP}} \right) - \psi \left(\frac{1}{2} \right), \quad (23)$$

where ψ is the digamma function, V_h is determined by Eqs. (13) with $\tilde{k}_h = (1+i)k_h$, and in the region of parameters, where $T_c \neq 0$ [the corresponding conditions can be extracted from the results for the critical thickness — see Eqs. (25), (26) below], we may write

$$W = \text{Re} V_h + (\text{Im} V_h)^2 d_s / \xi_s. \quad (24)$$

Due to symmetry, the result for the P case (22) reproduces that for the SF bilayer with S layer of thickness d_s [13]. In the AP case, if the second terms in the r.h.s. of Eq. (24) can be neglected (e.g., at $k_h d_f \gg 1$ in the region of parameters where $T_c \neq 0$), then $W = \text{Re} V_h$ and we reproduce the result of Ref.

[8]. However, the second term becomes essential in the Cooper limit, defined by conditions $d_s \ll \sqrt{D_s / 2\omega_D}$, $d_f \ll \min(\sqrt{D_f / 2\omega_D}, k_h^{-1})$, $\gamma_b = 0$, with ω_D the Debye energy of the S material. In this case $\text{Re} V_h = 0$ and Eqs. (23), (24) reproduce the result of Tagirov [5].

The critical thickness d_{sc} of the S layer, below which the superconductivity vanishes, immediately follows from Eqs. (22), (23) for the P and AP cases:

$$d_{sc}^P / \xi_s = 2e^C |V_h|, \quad d_{sc}^{AP} / \xi_s = 2e^C W \quad (25)$$

at

$$d_{sc} / \xi_s \ll 1. \quad (26)$$

Here $C \approx 0.577$ is Euler's constant. Condition (26) is necessary for applicability of Eqs. (25). If this condition is not satisfied, then Eqs. (25) only tell us that at $d_s / \xi_s \ll 1$ the superconductivity is certainly absent, i.e., $T_c = 0$. According to the monotonic growth of $T_c(\alpha)$, the function $d_{sc}(\alpha)$ decreases monotonically, hence $d_{sc}^P > d_{sc}^{AP}$. At $\gamma_b = 0$, $k_h d_f \gg 1$, Eqs. (25) reproduce the results of Ref. [11] for the P and AP cases.

The $T_c(\alpha)$ dependence can be most easily studied in the Cooper limit. In this case a simple analysis (see, e.g., Appendix A1 in Ref. [13]) can be done already on the level of the Usadel equations, and the system is described as a uniform layer with the effective exchange energy²⁾

$$h_{\text{eff}} = (\tau_f / \tau_s) h \cos \alpha, \quad (27)$$

where $\tau_{s(f)} = 2d_{s(f)} R_b \mathcal{A} / \rho_{s(f)} D_{s(f)}$. The accuracy of this result is limited to the first order over h , which becomes insufficient in the vicinity of $\alpha = \pi/2$. At $\alpha = \pi/2$, the first-order effect of h vanishes, while a more accurate analysis (Ref. [5] and Eqs. (23), (24)) reveals the second-order effect of h on T_c .

Let us now consider the same limit as in Ref. [11]:

$$d_s \ll \xi_s, \quad k_h d_f \gg 1, \quad h \gg \pi T_{cs}, \quad \gamma_b = 0, \quad (28)$$

$$\gamma k_h \xi_f d_s / \xi_s \ll 1. \quad (29)$$

The condition to have superconductivity at least at some orientations has the form $d_{sc}^{AP} < d_s \ll \xi_s$, and in the case under discussion, Eqs. (25), (26) yield:

$$2e^C \gamma k_h \xi_f < d_s / \xi_s \ll 1, \quad (30)$$

hence condition (29) becomes redundant.

²⁾Since ω_n was neglected in comparison with h in the Usadel equation, the result of the Cooper limit is valid only at $\tau_s \gg \tau_f$.

Starting from Eqs. (20), (15), (21), we finally obtain the following equation for T_c :

$$\ln \frac{T_{cs}}{T_c} = Q\psi\left(\frac{1}{2} + \frac{\Omega_1}{2\pi T_c}\right) + R\psi\left(\frac{1}{2} + \frac{\Omega_2}{2\pi T_c}\right) - \psi\left(\frac{1}{2}\right), \quad (31)$$

where

$$Q = \frac{1}{2} + \frac{\sin^2 \alpha}{2\sqrt{\sin^4 \alpha - 4 \cos^2 \alpha}}, \quad R = 1 - Q, \\ \Omega_{1,2} = \frac{d_0}{d_s} \pi T_{cs} \left(1 + \cos^2 \alpha \pm \sqrt{\sin^4 \alpha - 4 \cos^2 \alpha}\right), \\ d_0 = \gamma k_h \xi_f \xi_s / 2. \quad (32)$$

In the P and AP cases, where the triplet component is absent, Eqs. (31), (32) reproduce the results of Refs. [6, 11]. At the same time, at a noncollinear alignment the results are clearly different.

The critical thickness is found from Eqs. (31), (32):

$$d_{sc}(\alpha)/d_0 = 4\sqrt{2}e^C \cos \alpha \times \left(\frac{1 + \cos^2 \alpha + \sqrt{\sin^4 \alpha - 4 \cos^2 \alpha}}{1 + \cos^2 \alpha - \sqrt{\sin^4 \alpha - 4 \cos^2 \alpha}}\right)^{\frac{\sin^2 \alpha}{2\sqrt{\sin^4 \alpha - 4 \cos^2 \alpha}}}. \quad (33)$$

Although the square root in this expression can become imaginary, the whole expression remains real (z^i is real if $|z| = 1$). Figure 4 illustrates the result (33).

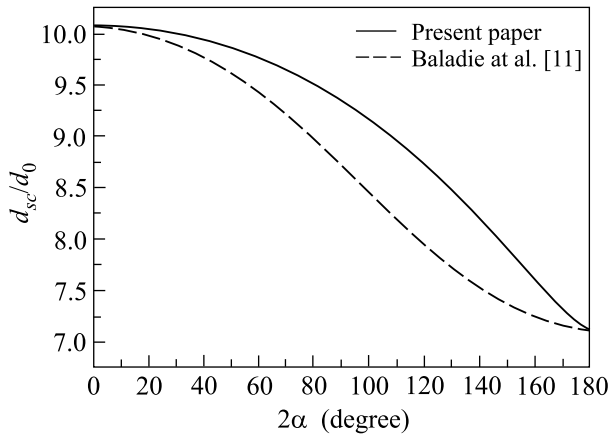


Fig.4. Critical thickness of the S layer d_{sc} vs. misorientation angle 2α . Dashed line is the result of Ref. [11], obtained without account of the triplet component

Now we turn to analyze the conditions of applicability for the results reported in Ref. [3]. A noncollinear FSF trilayer is a unit cell of the multilayered structure studied in that work. The main result of Ref. [3], the Josephson current due to the long-range triplet component, requires that the S layer is thin $d_s \ll \xi_s$, while the

F layers are thick for the singlet component and moderate for the triplet one: $k_h^{-1} \ll \xi_f < d_f$ [3]. In this case the condition that superconductivity is not completely suppressed at least in the vicinity of the AP alignment (Eqs.(25),(26)) takes the form

$$4e^C \gamma k_h \xi_f \frac{1 + 2\gamma_b k_h \xi_f}{(1 + 2\gamma_b k_h \xi_f)^2 + 1} < \frac{d_s}{\xi_s} \ll 1. \quad (34)$$

At $\gamma_b = 0$ (as it was assumed in Ref. [3]), this yields $2e^C \gamma k_h \xi_f < d_s/\xi_s \ll 1$, which is a rather strong condition for γ , since $k_h \xi_f \gg 1$. Finite interface transparency relaxes this condition: already at $\gamma_b \sim 1$, Eq. (34) yields $2e^C \gamma/\gamma_b < d_s/\xi_s \ll 1$.

The condition that superconductivity exists at all orientations has the form similar to Eq. (34) but with the corresponding expression for d_{sc}^P instead of d_{sc}^{AP} in the l.h.s. This only leads to a minor difference, since the two critical thicknesses are of the same order: $d_{sc}^P = \sqrt{2}d_{sc}^{AP}$ at $\gamma_b = 0$, while $d_{sc}^P = d_{sc}^{AP}$ at $\gamma_b > 1$.

In conclusion, we have studied T_c of a FSF trilayer as a function of its parameters, in particular, the angle between magnetizations of the F layers. The angular dependence becomes pronounced when the S layer is thin, and can lead to switching between superconducting and non-superconducting states as the angle is varied. Our results directly apply to multilayered SF structures, where a FSF trilayer is a unit cell. We have formulated the conditions which are necessary for existence of recently investigated odd triplet superconductivity in SF multilayers [3].

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