Pis'ma v ZhETF, vol. 77, iss. 4, pp. 212 – 216

Influence of interface double electric layer on superconducting proximity effect in ferromagnetic metals

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Submitted 10 January 2003

Electric field within the double electric layer that occurs when two conductors with different work functions come into contact induces the interface spin-orbit coupling. In the case of contact of a conventional, s-wave superconductor with a ferromagnetic metal, the coupling is predicted to give rise to the appearance of a triplet superconductivity on both sides of the interface. Form of the triplet component of the condensation amplitude is determined and the ability of the triplet superconductivity to penetrate into the ferromagnet on the usual coherence length is shown.

PACS: 71.70.Ej, 73.40.Jn, 74.20.−z, 74.50.+r

The phenomenon of penetration of superconductivity into a normal metal placed in contact with a superconductor (S) is now reasonably well understood [1]. However, mechanisms through which superconductivity can penetrate into a ferromagnetic metal are still not clear-cut. In accord with theory developed up to now [2], the penetrating of the singlet Cooper pair into a ferromagnet (F) looks as follows. The Fermi surfaces in the F that correspond to the conduction electrons with spins directed along (up-states) and against (down-states) the exchange field do not coincide. Therefore, upon entering the F region, two electrons that the Cooper pair is composed of have to occupy different Fermi surfaces. The difference between the Fermi momenta of the surfaces makes it impossible for the total momentum of the pair be zero. As a result, the condensation amplitude, which in the case of a normal metal would exponentially decay on the length scale of the coherence length, $\xi_0 = v_f(2\pi T_c)^{-1}$ (here v_f is the average Fermi velocity in the normal metal), acquires an additional modulation factor which oscillates with a short period proportional to $\lambda_F = 2\pi(k_{f,d} - k_{f,u})^{-1} \ll \xi_0$ (here $k_{f.d}$ and $k_{f.u}$ are the Fermi wave vectors of down- and up-states of the F, respectively). So λ_F becomes to be the effective penetration depth of superconductivity into the ferromagnetic metal. Some of experimental data published recently [3] seem to conform the theory; there are also experimental results [4] that much less agree with the theory and indicate that superconductor may influence ferromagnet on a much long distance than λ_F. On the grounds of all that has been said, one can expect the condensation amplitude to cure the oscillations and hence to penetrate deep into the F region if two electrons of the Cooper pair gain the possibility at equal energy to have opposite momenta so that to reduce the total momentum to zero. For that the both electrons should occupy the same Fermi surface, i.e., should turn out in the same spin state. The very fact that it is easier for the triplet superconductivity than for the singlet one to coexist with ferromagnetism is a simple consequence of the weak sensitivity of some of the triplet states to an external magnetic field well known from studies of superfluid He^3. The problem is to reveal mechanisms that are able to transform near the contact the initially singlet superconductivity to the triplet one. The purpose of this paper is to point out a natural reason for such a transformation. The reason is the interface spin-orbit coupling that is always brought about by contact.

In fact, when two metals come into contact, electrons are known to transfer from metal with the lower work function to the metal with higher work function, so as to bring the two Fermi levels into coincidence. As a result, a double electric layer of a thickness of the order of the screening length r_D appears. When a conduction electron is in the layer, it is subject to the electric field that induces the spin-orbit (SO) coupling. Because r_D in metals is about the lattice constant that is much shorter than the expected decay length of the triplet superconductivity, this interface SO coupling can be put into the form $H_{SO} = \alpha(\mathbf{p} \times \mathbf{c}) \cdot \boldsymbol{\sigma}\delta(\mathbf{c} \cdot \mathbf{r})$, where $\mathbf{p}$, $\boldsymbol{\sigma}$, and $\mathbf{c}$ are, respectively, the electron momentum, the Pauli matrices, and the normal directed from the F to S region. The Plank's constant $\hbar$ is set to unity throughout. The interface SO coupling arises from the electric field in the double layer in the same way as the well known $\mathbf{p}$-linear Hamiltonian in the energy spectrum of bulk polar semiconductors [5] arises from the intra-crystalline electric field [6, 7]. The presence of H_{SO} violates the total spin

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conservation and should result in singlet-triplet mixing. Recently, some consequences of the mixing for the proximity effect in normal metals were pointed out [8]. Here we consider how the mixing facilitates the penetration of superconductivity in the ferromagnetic metal under the assumption that the exchange field is uniform and directed along the interface.

We will assume for simplicity that the electron spectrum on both sides of the contact is isotropic. Then the one-particle Hamiltonian of the system has the form

$$H_0(\mathbf{p}) = \mathbf{p}\frac{1}{2m(z)}\mathbf{p} + \theta(-z)\left[\Gamma + \frac{b}{2}(\mathbf{h}\cdot\sigma + 1)\right] + \\ + \left[\beta + \alpha(\mathbf{p}\times\mathbf{c})\cdot\sigma\right]\delta(z), \qquad (1)$$

where z is the coordinate along $\mathbf{c}$, the electron mass $m(z)$ equals m_F in the F and m_S in the S, (for definiteness, we assume $m_S \geq m_F$), $\theta(z)$ is the unit step function, Γ is the difference between the lower conduction band edge of the F and of the S, $\mathbf{h}$ is the unit vector along the exchange field, b is the exchange energy, and $\beta\delta(z)$ models the spin-independent interface potential. The straightforward way to look into the behavior of superconductivity is to investigate the self-consistent equation for the gap-matrix $\Delta_{\alpha\beta}(\mathbf{r},\mathbf{r}')$ which near T_c has the form [9]

$$\Delta_{\alpha\beta}(\mathbf{r_1},\mathbf{r_2}) = -T_c\sum_{\epsilon}\int dr_3 dr_4 \times \\ \times V^{\alpha\beta|\gamma\delta}(\mathbf{r_1},\mathbf{r_2}\mid\mathbf{r_3},\mathbf{r_4})F_{\delta\gamma}(\mathbf{r_4},\mathbf{r_3}\mid i\epsilon), \qquad (2)$$

$$F_{\kappa\rho}(\mathbf{r_1},\mathbf{r_2}\mid i\epsilon) = \int dr_3 dr_4 G_{\kappa\gamma}(\mathbf{r_1},\mathbf{r_3}\mid i\epsilon)\times \\ \times\Delta_{\gamma\delta}(\mathbf{r_3},\mathbf{r_4})(-1)G^t_{\delta\rho}(\mathbf{r_4},\mathbf{r_2}\mid i\epsilon), \qquad (3)$$

where $V^{\alpha\beta|\gamma\delta}(\mathbf{r_1},\mathbf{r_2}\mid\mathbf{r_3},\mathbf{r_4})$ is the inter-particle interaction, the superscript t denotes transposition of the spinor indices (and also space coordinates when refers to Green's functions), G is the Green's function of independent particles and F is the anomalous Green's function. But before being able to enter the subject one must ascertain the form of the G function. This can be done in the following way. (i) Since the interface is assumed translational invariant in the $x-y$ plane, the electron momentum $\mathbf{k}_\parallel$ along the plane is conserved and the one-particle scattering problem is effectively one dimensional. The eigenfunctions of the 1D Schrödinger equation can be easily found and are not presented here. (ii) It is convenient to expand the G function on the spinor basis that includes the projecton operators on the up- and down-states $G_{\alpha\beta} = \sum_{i=1}^{4} M^{(i)}_{\alpha\beta}G_{(i)}$, where $M^i = \left\{\Pi^{(u)},\Pi^{(d)}\,A^{(\uparrow\downarrow)}.A^{(\downarrow\uparrow)}\right\}$, $\Pi^{(u,d)} = \frac{1}{2}(1\pm\mathbf{h}\cdot\sigma)$, $A^{(\uparrow\downarrow)} = \frac{1}{2}h_+(-\mathbf{c}\cdot\boldsymbol{\sigma} - i\mathbf{h}\times\mathbf{c}\cdot\boldsymbol{\sigma})$, $A^{(\downarrow\uparrow)} = \left(A^{(\uparrow\downarrow)}\right)^+$, $h_\pm = h_x \pm ih_y$. (iii) The chemical potential μ is supposed to be grater than $\Gamma + b$ so that both the up- and down-states in the F are partially filled. The spectral decomposition of the G function with definite longitudinal momentum $\mathbf{k}_\parallel$ can be constructed in line with the general theory of 1D scattering [10] with the help of the full set of Jost functions - the orthogonal and normalized scattering states of particle corresponding to waves ingoing from the left and from the right. One should take into account that a particle with energy in the interval $(0,\Gamma)$ may be incident only from the right and one with energy in the interval $(\Gamma,\Gamma+b)$ may be incident from the left only being in the down-state . (iv) The integration on the spectral parameters (momenta $\mathbf{k}_\parallel$ and $k_\perp$) can be performed using the stationary phase approximation which holds on a scale large compared to the Fermi wave-vector of any of the occupied bands. (v) Finally, one should sum up over all spin states. In the sector $z<0$, $z'>0$ which is first of all needed for the following analysis the result turns out to be

$$G_{(u)}(\mathbf{r},\mathbf{r}'|i\epsilon) \cong \\ \cong \left(i\pi D_u(\mathbf{p})N_u^{1/2}(\mathbf{p})\right)^{-1}\exp\{i\Phi_u(\mathbf{p})\}\Big|_{\mathbf{p}=\mathbf{p}_u}, \qquad (4)$$

$$G_{(\uparrow\downarrow)}(\mathbf{r},\mathbf{r}'|i\epsilon) \cong 2V_{ud}(\mathbf{p})\times \\ \times\left(i\pi D_u(\mathbf{p})D_d(\mathbf{p})N_u^{1/2}(\mathbf{p})\right)^{-1}\exp\{i\Phi_u(\mathbf{p})\}\Big|_{\mathbf{p}=\mathbf{p}_u}, \qquad (5)$$

where $\mathbf{r} = (\boldsymbol{\rho}, z)$, $\mathbf{r}' = (\boldsymbol{\rho}', z')$,

$$D_u(\mathbf{p}) = im_S^{-1}\left(k_{f.s}^2 - p^2\right)^{1/2} + \\ + im_F^{-1}\left(k_{f.u}^2 - p^2\right)^{1/2} - 2\left(\beta + \alpha\mathbf{p}\times\mathbf{c}\cdot\mathbf{h}\right), \qquad (6)$$

$$N_u(\mathbf{p}) = |\rho-\rho'|p^{-1}\left[z'k_{f.s}^2\left(k_{f.s}^2-p^2\right)^{-3/2} - \right. \\ \left. -zk_{f.u}^2\left(k_{f.u}^2-p^2\right)^{-3/2}\right], \qquad (7)$$

$$\Phi_u(\mathbf{p}) = z'\left(k_{f.s}^2-p^2\right)^{-1/2}\left(k_{f.s}^2 + i\epsilon m_S\right) - \\ - z\left(k_{f.u}^2-p^2\right)^{-1/2}\left(k_{f.u}^2 + i\epsilon m_F\right), \qquad (8)$$

$p = |\mathbf{p}|$, $k_{f.s}$ is the Fermi momentum of the S ($k_{f.u} \leq k_{f.d} \leq k_{f.s}$), $V_{ud} = \frac{\alpha}{2i}(p_+ + p_-h_+^2)$, $p_\pm = p_x \pm ip_y$, $\mathbf{p}_u = [(\boldsymbol{\varrho}-\boldsymbol{\varrho}')/|\rho-\rho'|]\,p_u$ is the 2D vector whose value p_u is defined as the root of the equation

$$|\boldsymbol{\varrho}-\boldsymbol{\varrho}'| = p\left[z'\left(k_{f.s}^2-p^2\right)^{-1/2} - z\left(k_{f.u}^2-p^2\right)^{-1/2}\right], \qquad (9)$$

and $D_d(\mathbf{p})$ is obtained from $D_u(\mathbf{p})$ if one changes in the r.h.s. of Eq. (6) $k_{f.u} \to k_{f.d}$ and $\alpha \to -\alpha$. An inspection of Eq. (9) reveals that $p_u \le k_{f.u}$ at any $\mathbf{r}$ and $\mathbf{r}'$. The implicit dependence of the vector $\mathbf{p}_u$ (and hence $G_{(u)}$ and $G_{(\uparrow\downarrow)}$) on $\mathbf{r}$ and $\mathbf{r}'$ is a result of the Fermi wave vector mismatch, $k_{f.s} \neq k_{f.u}$. To obtain $G_{(u)}$ at $\epsilon < 0$ one should change in Eq. (4) $\epsilon \to -\epsilon, \mathbf{p}_u \to -\mathbf{p}_u$ and take complex conjugation. The function $G_{(\uparrow\downarrow)}$ at $\epsilon < 0$ is obtained in the same way except for the complex conjugating of V_{ud}. The second pair of the G-functions, $G_{(d)}$ and $G_{(\downarrow\uparrow)}$, has an analogous form. They are given by the r.h.s. of Eqs. (4)–(5), respectively, where all quantities denoted by a subscript u ($k_{f.u}$, D_u, N_u, p_u and Φ_u) are changed by the quantities denoted by a subscript d with the prescription $V_{ud} \to V_{du} = (V_{ud})^*$. There is an asymmetry between $G_{(\uparrow\downarrow)}$ and $G_{(\downarrow\uparrow)}$. The former contains the function $D_u(\mathbf{p}_d)$ whereas the latter contains $D_d(\mathbf{p}_u)$. The defining equation for p_d (which is nothing but Eq. (9) where $k_{f.u}$ is replaced by $k_{f.d}$) says only that $p_d \le k_{f.d}$. But it may appear that $p_d > k_{f.u}$ at some $\mathbf{r}$ and $\mathbf{r}'$. One can show that in such a case the equation $(k_{f.u}^2 - p_d^2)^{1/2}$, which enters $D_u(\mathbf{p}_d)$, should be changed by $i(p_d^2 - k_{f.u}^2)^{1/2}$. As opposed to that, the quantity $k_{f.d}^2 - p_u^2$, which enters $D_d(\mathbf{p}_u)$, is always positive. This asymmetry reflects the fact that at $k_{f.u} \le k_{f.d}$ not every electron incident from the S being in the down-state can pass to the up-states of the F even if it is allowed to pass into the down-state of the F.

The inter-particle interaction (in both metals) includes all spherical harmonics, however the conventional character of the S assumed means that the most strong attraction takes place for electron pairs in singlet s-wave state. In spite of the changes in the G-function induced by H_{SO}, upon solving Eqs. (2)–(3) with

$$V^{\alpha\beta|\gamma\delta}(\mathbf{r}_1, \mathbf{r}_2 \mid r_3, r_4) = V_s^{\alpha\beta|\gamma\delta}(\mathbf{r}_1, \mathbf{r}_2 \mid \mathbf{r}_3, \mathbf{r}_4) =$$

$$= \lambda_s(\mathbf{r}_1)\, g_{\alpha\beta} g^t_{\gamma\delta} \delta(\mathbf{r}_1 - \mathbf{r}_2)\delta(\mathbf{r}_3 - \mathbf{r}_4)\delta(\mathbf{r}_1 - \mathbf{r}_3),$$

where $g = i\sigma_y$, $\lambda_s(\mathbf{r}) = \theta(-z)\lambda_s(F) + \theta(z)\lambda_s(S)$, $\lambda_s(F)$ and $\lambda_s(S)$ are the coupling constants in the F and S, respectively, one obtains the singlet gap-matrix $\Delta^{(s)}_{\alpha\beta}(\mathbf{r}, \mathbf{r}') = \delta(\mathbf{r} - \mathbf{r}')g_{\alpha\beta}\Delta_{(s)}(z)$ with the usual behavior of $\Delta_{(s)}(z)$ described at the beginning of the paper up to small corrections of the order of $(\alpha m)^2 \ll 1$ – it is a smooth function at $z > 0$ that approaches a constant value $\Delta_{(s)}(\infty)$ in the bulk of the S and it is a rapidly oscillating function at $z < 0$. At the same time, $F_{\alpha\beta}(\mathbf{r}_1, \mathbf{r}_2)$ being given by the r.h.s of Eq. (3) with $\Delta_{\alpha\beta}(\mathbf{r}_1, \mathbf{r}_2) = \Delta^{(s)}_{\alpha\beta}(\mathbf{r}_1, \mathbf{r}_2)$ gains a different quality – it acquires the triplet component, i.e., the component which is even at $\alpha \rightleftarrows \beta$ and odd at $\mathbf{r}_1 \rightleftarrows \mathbf{r}_2$, and what is more – the component does not suffer from the rapid oscillations. To verify the fact, one needs an appropriate projector, i.e., such an operator that a nonzero result of its application to the F-function would definitely indicate the presence of the triplet part. As such a projector, one can choose the operator of convolution of the F-function with the inter-particle interaction in the triplet p-wave channel

$$-T\sum_{\epsilon} \int dr_3 dr_4 V_p^{\alpha\beta|\gamma\delta}(\mathbf{r}_1, \mathbf{r}_2 \mid \mathbf{r}_3, \mathbf{r}_4) F_{\delta\gamma}(\mathbf{r}_4, \mathbf{r}_3 \mid i\epsilon), \tag{10}$$

Here

$$V_p^{\alpha\beta|\gamma\delta}(\mathbf{r}_1, \mathbf{r}_2 \mid \mathbf{r}_3, \mathbf{r}_4) = \lambda_p(\mathbf{r}_1)(\sigma^k g)_{\alpha\beta} \times$$

$$\times (g^t \sigma^k)_{\gamma\delta} \delta(\mathbf{r}_1 - \mathbf{r}_2)\delta(\mathbf{r}_3 - \mathbf{r}_4)\delta(\mathbf{r}_1 - \mathbf{r}_3)\nabla^i_{12}\nabla^i_{34}$$

is the $e - e$ interaction in the triplet p-wave channel [11] and $\nabla^i_{12} = (\partial/\partial r_2^i - \partial/\partial r_1^i)/2ik_f$, where k_f is a representative wave-vector of the order of $k_{f.s}$. Its precise value is not important because it always enters equations being multiplied by λ_p. If one considers the convolution only from mathematical viewpoint, λ_p can take any nonzero value. However, because the actual $e - e$ interaction in the metals surely has the nonzero triplet p-wave component, one can choose λ_p equal to its true value, $\lambda_p(\mathbf{r}) = \lambda_p(S)\theta(z) + \lambda_p(F)\theta(-z)$. Then the convolution will define the triplet component of the gap-matrix, $\Delta^{(t)}_{\alpha\beta}(\mathbf{r}_1, \mathbf{r}_2)$, just as the convolution with the singlet interaction defines the singlet component of the matrix.

Being represented as a function of the Cooper pair center of mass $\mathbf{r}$ and the momentum of the relative motion $\mathbf{p}$, the convolution takes the form $\Delta^{(t)}_{\alpha\beta}(\mathbf{r}, \mathbf{p}) = A_{ij}(\mathbf{r})\frac{p_j}{k_f}(\sigma_i g)_{\alpha\beta}$ (it will be seen below that only conserving components of $\mathbf{p}$ parallel to the interface enter this equation), where

$$A_{ki}(\mathbf{r}) = \lambda_p(\mathbf{r}) \int d^3 r' K^{ik}_{(st)}(\mathbf{r}, \mathbf{r}')\Delta_{(s)}(\mathbf{r}') \tag{11}$$

$$K^{ik}_{(st)}(\mathbf{r}_1, \mathbf{r}_2) = \lim_{\mathbf{r}'_1 \to \mathbf{r}_1} \nabla^i_{11'} T \times$$

$$\times \sum_{\epsilon_n} Tr\left[g^t \sigma^k G_{\epsilon_n}(\mathbf{r}_1, \mathbf{r}_2) g G^t_{-\epsilon_n}(\mathbf{r}_2, \mathbf{r}'_1)\right]. \tag{12}$$

By evaluating the kernel $K^{ik}_{(st)}$ by means of the G-function obtained one should remember that rapidly oscillating terms should be removed, i.e., one should consider the kernel averaged over distances large compared to k_f but small compared to $\xi_0 = v_f(2\pi T_c)$ (here and below v_f is a characteristic Fermi velocity of the system), and also that because by deriving Eq.(12)

we did not introduce the frequency cutoff of the order of the Debye frequency ω_D, the equation is valid at $|z_1 - z_2| > \delta \sim v_f/\omega_D$ [12]. By expanding the Green's functions on the spinor basis M^i, one can see that major contributions come from the pairs of the Green's functions $G_{(u)}(\mathbf{r}_1, \mathbf{r}_2|i\epsilon)$, $G_{(\uparrow\downarrow)}(\mathbf{r}_2, \mathbf{r}_1'| - i\epsilon)$ and $G_{(d)}(\mathbf{r}_1, \mathbf{r}_2|i\epsilon)$, $G_{(\downarrow\uparrow)}(\mathbf{r}_2, \mathbf{r}'|-i\epsilon)$. In these contributions, the largest parts of the phases Φ_u and Φ_d (see Eq. (8)) that do not depend on the imaginary frequencies, $i\epsilon_n = i\pi T(2n+1)$, exactly cancel each other, thereby avoiding the oscillating behavior of $A_{ji}(\mathbf{r})$. The first pair of the functions forms a part of $K^{ik}_{(st)}$ that gives rise to the triplet component with $\mathbf{S}\cdot\mathbf{h} = 1$ (where $\mathbf{S}$ is the total spin of the Cooper pair) and the second one results in formation of the component with $\mathbf{S}\cdot\mathbf{h} = -1$. The triplet component with $\mathbf{S}\cdot\mathbf{h} = 0$ can be shown to rapidly oscillate like $\Delta_{(s)}(z)$. So one gets $A_{ji}(\mathbf{r}) = h_i(h\times c)_j\Delta_{(t1)}(z) + ih_ic_j\Delta_{(t2)}(z)$. The explicit form of the kernel $K^{ik}_{(st)}$ is rather awkward at arbitrary parameters of the energy spectrum of the F. Therefore, we consider below the limiting case of equal masses $m_S = m_F \equiv m$ and weak exchange energy $\xi_0^{-1} \ll k_{f.d} - k_{f.u} \ll k_{f.d} + k_{f.u}$. The difference between $k_{f.s}$ and $\frac{1}{2}(k_{f.d} + k_{f.u})$ will be ignored as well. Then, with an accuracy up to corrections of the order of $(b/\mu)^2$, at $z < 0$ one gets

$$\Delta_{(t1,2)}(z) = -\alpha m\lambda_p(N)N(0)\xi_0^{-1}\times \\ \times\int\limits_0^\infty dz'\int\limits_0^\infty du\, L_{(1,2)}(z,z',u)\Delta_{(s)}(z'), \qquad (13)$$

$$L_{(1)}(z,z',u) = \frac{B}{2}u(1+u)^{-1}\left[1+B^2(1+u)\right]^{-2}\times \\ \times\sinh^{-1}\left(\frac{\zeta}{\xi_0}\sqrt{1+u}\right), \qquad (14)$$

$$L_{(2)}(z,z',u) = \frac{b}{2\mu}\frac{|z|}{|z|+z'}\times \\ \times\left[1+\frac{|z|+z'}{4\xi_0}\sqrt{1+u}\coth\left(\frac{|z|+z'}{\xi_0}\sqrt{1+u}\right)\right]\times \\ \times L_{(1)}(z,z',u). \qquad (15)$$

Here $N(0) = mk_f(2\pi^2)^{-1}$ is the electron density of states per unit energy interval at the Fermi level, and $B = m\beta/k_F$. It is seen that the functions $\Delta_{(t1,2)}(z)$ and hence the triplet part of the F-function concentrate in a domain of width ξ_0. A contribution to the r.h.s. of Eq. (13) from integration over negative z' is small because of the oscillating behavior of $\Delta_{(s)}(z)$ at $z < 0$. One can show that A_{ji} on the S side of the contact also concentrates in a domain of the same order and at any energy spectrum of the F has only $h_i(h\times c)_j$ and $(c\times h)_ih_j$ components but not ih_ic_j component. In the weak exchange energy limit, the components combine to $e_{jin}c_n$ yielding the same triplet order parameter as in the case of N/S contact [8] plus corrections of the order of $(b/\mu)^2$.

In principle, one ought from the very beginning to introduce both singlet and triplet components of the gap-matrix and the s- and p-wave inter-particle interactions. Then the spinor self-consistency equation reduces to a symmetrical set of two simultaneous equations for $\Delta_{(s)}(\mathbf{r})$ and $A_{ji}(\mathbf{r})$. However, the conventional character of superconductivity assumed in the bulk of the S means that the s-wave interaction is attractive and larger in magnitude. Therefore, in the bulk, where the influence of H_{SO} disappears, only the s-wave interaction is relevant in the sense that the convolution of the F-function with the p-wave interaction vanishes. Thus there is only $\Delta_{(s)}(\mathbf{r})$ in the bulk. The p-wave interaction becomes to be relevant near the interface. If λ_p is small, the set of equations transforms into an usual equation for $\Delta_{(s)}$ and Eq. (11). Thus, although both components of the gap-matrix, $\Delta^{(s)}_{\alpha\beta}$ and $\Delta^{(t)}_{\alpha\beta}$, are simultaneously initiated at the critical temperature T_c, they do not possess equal rights - the singlet order parameter is the source for the triplet one. It should be also stressed the important difference between the triplet part of the F-function, which is the basic conceptual element of the pairing theory, and that of the gap-matrix, which is only one of characteristics of the F-function. The latter is proportional to λ_p whereas the former at small λ_p is independent of λ_p.

Since $B \sim U_b/\mu$, where U_b is the height of the interface barrier, one may estimate $B \sim 1$. The small quantity in Eq. (13) is $\alpha m \sim \alpha_V/v_F$, where α_V is the 'bulk' SO constant within the double electron layer. A tentative estimate of the order of magnitude of α_V (and hence αm) done previously [8] shows that αm can be of the order of 10^{-2} under favorite conditions. Thus the F-function in the F has the rapidly oscillating singlet component and small but smooth triplet component [13]. (In the general case, there are no reasons for λ_p/λ_s to be anomalously small. So $\Delta_{(t1,2)}/\Delta_{(s)} \sim (\alpha m)(\lambda_p N(0))$ can well appear to be of the order of $10^{-2} \div 10^{-3}$. However, discussion of possible effects of $\Delta_{(t1,2)}$, in particular, its non-unitarity in the F, is out the scope of this paper.)

Finally, we remark that all presented above shows that the problem of the F/S contact admits rigorous consideration within Gor'kov-Nambu formulation of the pairing theory. The triplet – singlet mixing found is the consequence of inversion symmetry breaking due to the **p**-linear interface SO coupling. Still more interesting effects of simultaneous breaking of T and P symmetry can

be apparently brought out in non-equilibrium properties of the contact. So this system certainly merits further examination.

This work was partly supported by Grant No 01-02-16506 from RFBR.

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13. It should be mentioned that the triplet superconductivity on the F side of the F/S interface was also considered in the article by F.S.Bergeret, A.F.Volkov, and K.B.Efetov [Phys. Rev. Lett. **86**, 4096 (2001)] as a result of the magnetization rotation near the interface like that in the Bloch domain wall. In addition to the difference between models studied in the present and that paper, there is a fundamental difference between approaches applied. In that paper, the Usadel equation for the anomalous Green's matrix-function (AGF) as a function of the Cooper pair center of mass coordinates was solved. The anti-symmetric components of the AGF were, as usually, attributed to the singlet superconductivity whereas the symmetric ones were interpreted as the appearance of the triplet superconductivity. However, the total AGF, $F(\mathbf{r}_1, \mathbf{r}_2)_{\alpha\beta}$, (as well as its singlet and triplet parts separately) must change sign at $(\alpha, \mathbf{r}_1) \rightleftarrows (\beta, \mathbf{r}_2)$ due to the Pauli principle. It is unclear how one can assure this property for the symmetric component found within the Usadel approach, i.e., without a consideration of the relative coordinate dependence of the AGF.