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Triplet superconductivity and magnetoelectric effect near the *s*-wave-superconductor–normal-metal interface caused by local breaking of mirror symmetry

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An interface spin-orbit coupling induced by the absence of the ''left-right'' symmetry on the junction of two different conductors is argued to exist. The coupling is due to contact charge exchange between conductors with different work functions. The effect of the coupling on superconductivity near the interface of a conventional superconductor and a normal metal is studied in Gor'kov formulation by making use of the exact Green's functions of the interface scattering problem. The triplet component of the Cooper pair wave function and the spin polarization of the carriers in the state with supercurrent parallel to the interface are predicted to exist in a vicinity of the junction of the order of the coherence length.

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The discovery of optical activity in the 18th century has shown that studying effects that broken mirror symmetry may have on physical properties of matter is not only interesting as a fundamental problem but also can give rise to important practical applications. From this viewpoint, searching for uncommon features of superconductivity under broken space parity conditions is of the utmost relevance. Studies of the problem in three-dimensional (3D) polar superconductors done during past few decades has led to some positive results—novel magnetoelectric effects were predicted: (i) the supercurrent must be accompanied by the spin polarization of the carriers[1] and (ii) under the action of external magnetic field **B**, the magnitude of the critical supercurrent $\mathbf{J}_s(\mathbf{B})$ should be different for two opposite directions being dependent on the sign of the mixed product $\mathbf{c}\times\mathbf{B}\cdot\hat{\mathbf{J}}_s$,[2] where **c** is the polar axis. Several superconductors of polar symmetry are known at present (see references cited in Refs. 1 and 2). Unfortunately, they are not numerous and not yet widely adopted in technology.

Our purpose here is to point out a situation that is potentially important for emerging technology for producing and manipulating electronic-spin polarization[3] and where any conventional superconductor, such as Pb and Sn, can exhibit some properties characteristic of polar superconductors. The situation takes place near a contact of the superconductor with any normal metal. The uncommon properties will be shown to originate from a broken mirror symmetry at the contact.

Namely, if metals forming a contact are different, two normals to the interface that is supposed to be a plane are not equivalent. Owing to this symmetry fact, the electron Hamiltonian is allowed to include the interface spin-orbit (SO) term of the form

$$H_{so}=\alpha(\mathbf{p}\times\mathbf{c})\cdot\boldsymbol{\sigma}\delta(\mathbf{c}\cdot\mathbf{r}), \tag{1}$$

where **p**, $\boldsymbol{\sigma}$, and **c** are, respectively, the electron momentum, the Pauli matrices, and one of two nonequivalent normals. The Planck's constant $\hbar$ is set to unity throughout. The physics behind the term consists of the following. 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 appears. When a conduction electron is in this layer, it is subject to the electric field. It is this field that makes two normals nonequivalent and induces the SO coupling. Because the screening length in metals is short, the SO coupling energy can be put into the form of Eq. (1). An analogous Hamiltonian was introduced earlier in a phenomenological way in connection with the energy spectrum of bulk polar semiconductors.[4] Later on, a microscopic substantiation of the Hamiltonian was given for the case of bulk,[5] surface,[6] and heterostructure[7] electron states. A microscopic theory for the interface SO can possibly be developed by analogy with the works,[5–7] but it is not the subject of the present paper. Here, Eq. (1) will be considered phenomenologically.[8]

When one of the metals (say one on the interface's right) turns superconductor, the presence of H_{so} locally spoils the classification of the Cooper pair wave function in terms of the total spin and should result in singlet-triplet mixing. The problem of the triplet characteristics that a *s*-wave superconductor can show under no space parity conditions is, in fact, not a new one. It has been dealt with in a number of works in connection with the upper critical magnetic field,[9] the nonzero spin susceptibility[9,10] and the paramagnetic phase[10] of the condensate, the critical fluctuations,[11] a modification of the Ginzburg-Landau free-energy functional,[2] and the magnetoelectric effects mentioned above.[1,2] Recently, some of the results obtained in Ref. 10 were rediscovered and slightly extended.[12] However, the problem considered in this paper drastically differs from all problems investigated earlier. In the preceding works, all effects found were somewhat simple consequences of the fact that in those cases, the SO coupling constantly acts by lifting the spin degeneracy in the whole space available for the electron motion. Contrary to that, in the case of the contact, the SO coupling is absent in the whole 3D space except for the 2D plane of contact and the effect occurs due to a difference in the scattering efficiency of electrons with different helicity on the interface potential. However, in spite of the region of localization of the SO coupling is of zero measure, the effects discussed here will be shown to exist in a broad 3D domain of width of the

coherence length, $\xi_0 = v_F(2\pi T_c)^{-1}$, on both sides of the contact. From what has been just said, another important feature of the problem presented follows—it is essentially nonhomogeneous and, therefore, requires for its solution more involved theoretical means than those applied previously.

Although it is not a limitation of the method used, we will assume for simplicity that the electron spectrum is isotropic and the effect of the Fermi wave-vector mismatch can be ignored. Then the one-particle Hamiltonian of the system has the form

$$H_0(\mathbf{p}) = \frac{p^2}{2m} + [\beta + \alpha(\mathbf{p}\times\mathbf{c})\cdot\boldsymbol{\sigma}]\delta(z), \tag{2}$$

where z is the coordinate along $\mathbf{c}$ and $\beta\delta(z)$ models the spin-independent interface potential. The interparticle interaction (in both metals) includes all spherical harmonics $(s,p,d,\ldots)$, however, the conventional character of the superconductor assumed means that the strongest attraction takes place for electron pairs in the singlet s-wave state. So the main part of the interaction Hamiltonian is

$$H_s = \frac{1}{2}\int d^3r_{1234}\psi^+_\alpha(1)\psi^+_\beta(2)V_s^{\alpha\beta|\gamma\delta}(\mathbf{r}_1,\mathbf{r}_2|\mathbf{r}_3,\mathbf{r}_4) \times\psi_\delta(4)\psi_\gamma(3), \tag{3}$$

where $\psi_\alpha(\mathbf{r})$ is the electron quantum-field operator, $V_s^{\alpha\beta|\gamma\delta}(\mathbf{r}_1,\mathbf{r}_2|\mathbf{r}_3,\mathbf{r}_4) = \lambda_s(\mathbf{r}_1)g_{\alpha\beta}g^t_{\gamma\delta}\,\delta(\mathbf{r}_1-\mathbf{r}_2)\,\delta(\mathbf{r}_3-r_4)\,\delta(\mathbf{r}_1-\mathbf{r}_3)$ is the interparticle interaction in the singlet s-wave channel, $g = i\sigma_y$, $\lambda_s(\mathbf{r}) = \theta(-z)\lambda_s(N) + \theta(z)\lambda_s(S)$, $\lambda_s(N)$ and $\lambda_s(S)$ are the coupling constants in the normal metal and the superconductor, respectively, $\theta(z)$ is the unit step function, and the superscript t denotes transposition of spinor indices (and also space coordinates when refers to Green's functions). Near T_c, the gap matrix $\Delta_{\alpha\beta}(\mathbf{r},\mathbf{r}')$ satisfies the equations[13]

$$\Delta_{\alpha\beta}(\mathbf{r}_1,\mathbf{r}_2) = -T\sum_\epsilon \int dr_3 dr_4 V_s^{\alpha\beta|\gamma\delta}(\mathbf{r}_1,\mathbf{r}_2|\mathbf{r}_3,\mathbf{r}_4) \times F_{\delta\gamma}(\mathbf{r}_4,\mathbf{r}_3|i\epsilon), \tag{4}$$

$$F_{\kappa\rho}(\mathbf{r}_1,\mathbf{r}_2|i\epsilon) = \int dr_3 dr_4 G_{\kappa\gamma}(\mathbf{r}_1,\mathbf{r}_3|i\epsilon)\Delta_{\gamma\delta}(\mathbf{r}_3,\mathbf{r}_4) \times(-1)G^t_{\delta\rho}(\mathbf{r}_4,\mathbf{r}_2|i\epsilon), \tag{5}$$

where G is the Green's function of independent particles and F is the anomalous Green's function. When $\alpha = 0$, i.e., the interface potential is spin independent, the spinor structure of the gap matrix and the F function are the same and the equations admit a solution of the form $\Delta^{(s)}_{\alpha\beta}(\mathbf{r},\mathbf{r}') = g_{\alpha\beta}\Delta_{(s)}(z)\,\delta(\mathbf{r}-\mathbf{r}')$, where the singlet gap-function $\Delta_{(s)}(z)$ describes the usual proximity effect.[14] Things drastically change if one takes into account H_{so}.

First, the one-particle G function ceases to be a diagonal matrix in spinor space. Necessary steps to obtain the function are the following.

(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.

(ii) Because the helicity operator $\hat{\nu} = (\mathbf{k}_\parallel/|\mathbf{k}_\parallel|)\times\mathbf{c}\cdot\boldsymbol{\sigma}$ is conserved, the Green's function falls into two parts corresponding to positive and negative helicity, $\nu = \pm 1$.

(iii) The spectral decomposition of the Green's function with definite longitudinal momentum $\mathbf{k}_\parallel$ and helicity ν can be constructed according to the general theory of 1D scattering[15] with the help of the full set of Jost functions—the orthogonal and normalized scattering states of particle with energy $(\mathbf{k}_\parallel^2 + \mathbf{k}_\perp^2)/2m$ corresponding to waves ingoing from the left and from the right.

(iv) The integration on the spectral parameters (momenta $\mathbf{k}_\parallel$ and $k_\perp$) can be performed using the stationary phase approximation that holds on a scale large compared to k_F^{-1}. The result turns out to be

$$G(\mathbf{r},\mathbf{r}'|i\epsilon)_{\alpha\beta} \cong \sum_{\nu=+,-} \Pi^{(\nu)}_{\alpha\beta}(\mathbf{r}-\mathbf{r}'|\epsilon)G_{(\nu)}(\mathbf{r},\mathbf{r}'|i\epsilon), \tag{6}$$

$$\Pi^{(\nu)}(\mathbf{r}|\epsilon_n) = \frac{1}{2}\left(1 + \mathrm{sgn}(\nu)\mathrm{sgn}(\epsilon_n)\frac{\mathbf{r}\times\mathbf{c}}{|\mathbf{r}\times\mathbf{c}|}\cdot\boldsymbol{\sigma}\right), \tag{7}$$

where at $\epsilon_n > 0$,

$$G_{(\nu)}(\mathbf{r},\mathbf{r}'|i\epsilon) = G_0(\mathbf{r}-\mathbf{r}') + G_0(\mathbf{r}-\mathbf{r}'_{ref})R_{(\nu)}(\mathbf{r}-\mathbf{r}'_{ref}), \quad z>0,\ z'>0, \tag{8}$$

$$G_{(\nu)}(\mathbf{r},\mathbf{r}'|i\epsilon) = G_0(\mathbf{r}-\mathbf{r}')[1 - R_{(\nu)}(\mathbf{r}-\mathbf{r}')], \quad z>0,\ z'<0. \tag{9}$$

Here, $G_0(\mathbf{r}) = -(m/2\pi|\mathbf{r}|)\exp[(ik_F - \epsilon/v_F)|\mathbf{r}|]$ is the free 3D Green's function,

$$R_{(\nu)}(\mathbf{r}) = \frac{m[\beta + \alpha k_F \mathrm{sgn}(\nu)|\boldsymbol{\rho}|/|\mathbf{r}|]}{ik_F|z|/|\mathbf{r}| - m[\beta + \alpha k_F \mathrm{sgn}(\nu)|\boldsymbol{\rho}|/|\mathbf{r}|]}, \tag{10}$$

and $\mathbf{r}_{ref} = \mathbf{r} - 2\mathbf{c}(\mathbf{c}\cdot\mathbf{r})$ is the "reflected" vector with respect to the vector $\mathbf{r} = (\boldsymbol{\rho}, z)$. To obtain $G_{(\nu)}$ at $\epsilon < 0$, one should change in Eqs. (8) and (9) $\epsilon \to -\epsilon$ and take complex conjugation.

In spite of the changes in the spinor structure of G function, upon solving Eqs. (4) and (5), one obtains the same gap matrix, $\Delta^{(s)}_{\alpha\beta}$, up to small corrections of the order of $(\alpha m)^2 \ll 1$. At the same time, $F(\mathbf{r_1},\mathbf{r_2})_{\alpha\beta}$ being given by the right-hand side of Eq. (5) 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 $\alpha \to \leftarrow \beta$ and odd at $\mathbf{r_1} \to \leftarrow \mathbf{r_2}$. To verify this 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. The operator of convolution of the F function with the interparticle interaction in the triplet p-wave channel

$$-T\sum_\epsilon \int dr_3 dr_4 V_p^{\alpha\beta|\gamma\delta}(\mathbf{r_1},\mathbf{r_2}|\mathbf{r_3},\mathbf{r_4})F_{\delta\gamma}(\mathbf{r_4},\mathbf{r_3}|i\epsilon) \tag{11}$$

can be this projector. Here, $V_p^{\alpha\beta|\gamma\delta}(\mathbf{r_1},\mathbf{r_2}|\mathbf{r_3},\mathbf{r_4}) = \lambda_p(\mathbf{r_1})(\sigma^k g)_{\alpha\beta}(g^t\sigma^k)_{\gamma\delta}\delta(\mathbf{r_1}-\mathbf{r_2})\delta(\mathbf{r_3}-\mathbf{r_4})\delta(\mathbf{r_1}-\mathbf{r_3})\nabla^i_{12}\nabla^i_{34}$ is the *e-e* interaction in the triplet *p*-wave channel[16] and $\nabla^i_{12}=(\partial/\partial r^i_2-\partial/\partial r^i_1)/2ik_F$. If one considers the convolution only from mathematical viewpoint, λ_p can take any nonzero value. The convolution has the physical meaning as well. 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, $\lambda_p(\mathbf{r})=\lambda_p(S)\theta(z)+\lambda_p(N)\theta(-z)$. Then the convolution will define the triplet component of the gap-matrix, $\Delta^{(t)}_{\alpha\beta}(\mathbf{r_1},\mathbf{r_2})$, like the convolution with the singlet interaction [see Eq. (4)] 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 $\Delta^{(t)}_{\alpha\beta}(\mathbf{r},\mathbf{p})=A_{ij}(\mathbf{r})(p_j/k_F)(\sigma_i g)_{\alpha\beta}$, where

$$A_{ki}(\mathbf{r})=\lambda_p(\mathbf{r})\int d^3r' K^{ik}_{(st)}(\mathbf{r},\mathbf{r}')\Delta_{(s)}(\mathbf{r}') \tag{12}$$

$$K^{ik}_{(st)}(\mathbf{r_1},\mathbf{r_2})=\lim_{\mathbf{r}'_1\to\mathbf{r_1}}\nabla^i_{11'}T\sum_{\epsilon_n}Tr[g^t\sigma^k G_{\epsilon_n}(\mathbf{r_1},\mathbf{r_2}) \times g G^t_{-\epsilon_n}(\mathbf{r_2},\mathbf{r}'_1)] \tag{13}$$

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 ξ_0) and also that because by deriving Eq. (13), we did not introduce the frequency cutoff of the order of the Debye frequency ω_D, the equation is valid at $|z_1-z_2|>\delta\sim v_F/\omega_D$.[17] In this way, one gets $A_{ji}(\mathbf{r})=e_{jim}c_m\Delta_{(t)}(z)$, where

$$\Delta_{(t)}(z)=-\alpha m\lambda_p(N)N(0)\xi_0^{-1}\int_{-\infty}^{\infty}dz'\,\mathrm{sgn}(z')Q_1(|z|+|z'|)\Delta_{(s)}(z'), \tag{14}$$

at $z<0$ and $\Delta_{(t)}(z)=-\Delta_{(t)}(-z)\lambda_p(S)/\lambda_p(N)$ at $z>0$. Here

$$Q_1(\zeta)=\frac{B}{2}\int_0^\infty du\,u(1+u)^{-1}[1+B^2(1+u)]^{-2}\times\sinh^{-1}\left(\frac{\zeta}{\xi_0}\sqrt{1+u}\right), \tag{15}$$

$N(0)=mk_F(2\pi^2)^{-1}$ 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 function $Q(\zeta)$ as well as $\Delta_{(t)}(z)$, and hence the triplet part of the *F* function concentrate in a domain of width ξ_0 on both sides of the interface.

As was noticed above, the conventional nature of superconductivity assumed in the bulk of the superconductor means that the interparticle interaction in the singlet *s*-wave channel is negative and larger in magnitude than in any other channel. Therefore, the critical temperature T_c is the temperature, below which the Eqs. (4) and (5) begin to reveal the nontrivial solution $\Delta^{(s)}_{\alpha\beta}$ of the singlet type. Then in the bulk, where the influence of H_{so} disappears, interactions in other channels become irrelevant in the sense that the convolution of the *F* function with any of them except for the singlet *s*-wave channel vanishes. Thus there is only the singlet component of the gap matrix. However, two partial interactions, the singlet and triplet ones, become to be relevant near the interface that gives rise to two parts of the gap matrix, the singlet and triplet ones, respectively. The main difference between these components is that the singlet component, as seen from Eq. (12), is the source for the triplet one.

The presence of the triplet part allows one by analogy with bulk polar superconductors[1] to anticipate the magnetoelectric effect (MEE), i.e., the appearance of the magnetization near the interface in the state with supercurrent parallel to the interface. The mean value of the spin density is given by the equation $\langle\psi^{(+)}_\alpha(\mathbf{r})\sigma^i_{\alpha\beta}\psi_\beta(\mathbf{r})\rangle$, where the angular brackets mean the thermal average. According to general rules of quantum statistics, the equation can be expressed through the total electron propagator, which near T_c can be expanded in powers of the gap matrix. So one gets

$$\langle\sigma^i(\mathbf{r})\rangle=T\sum_{\epsilon}\int d^3r_1d^3r_2Tr\{\sigma^iG_\epsilon(\mathbf{r},\mathbf{r_1})\Delta^{(s)}(\mathbf{r_1}) \times(-1)G^t_{-\epsilon}(\mathbf{r_1},\mathbf{r_2})\Delta^{+(s)}(\mathbf{r_2})G_\epsilon(\mathbf{r_2},\mathbf{r})\}. \tag{16}$$

One should substitute here the *G* function from Eqs. (6)–(10) and $\Delta^{(s)}_{\alpha\beta}(\mathbf{r})=g_{\alpha\beta}\Delta_{(s)}(z)\exp(i\mathbf{Q}\cdot\mathbf{r})$, $\mathbf{Q}\perp\mathbf{c}$, where the real function $\Delta_{(s)}(z)$ is the solution of the usual zero-current proximity problem with spin-independent interface potential and vector $\mathbf{Q}$ is defined by the supercurrent density $\mathbf{J}_s$ far from the interface through the relation $\mathbf{J}_s=(2e/m)\mathbf{Q}N_s$, where $N_s=n_3 7\zeta(3)\Delta^2_{(s)}(\infty)/8\pi^2T_c^2$ is the superfluid electron density, $n_3=k_F^3/3\pi^2$ is the electron density, and $\zeta(3)$ is the Riemann ζ function. An inspection of the integrand in Eq. (16) reveals that at given $\mathbf{r}$ and $\mathbf{r_1}$, the dominant contribution to the integral is given by the domain of integration with respect to $\mathbf{r_2}$ near the straight line drawn through the points $\mathbf{r}$ and $\mathbf{r_1}$. The result takes the form

$$\langle\boldsymbol{\sigma}(\mathbf{r})\rangle=\frac{\mathbf{c}\times\mathbf{J}_s}{2ev_F}(\alpha m)M(z\xi_0^{-1}), \tag{17}$$

where

$$M(w)=\int_w^\infty dv\int_{-\infty}^\infty du\,\mathrm{sgn}(u)Q_2(|u|+|v|) \times\Delta_{(s)}(u\xi_0)\Delta_{(s)}(v\xi_0)/\Delta^2_{(s)}(\infty), \tag{18}$$

at $w>0$, and

$$M(w)=-\int_{-\infty}^w dv\int_{-\infty}^\infty du\,\mathrm{sgn}(u)Q_2(|u|+|v|) \times\Delta_{(s)}(u\xi_0)\Delta_{(s)}(v\xi_0)/\Delta^2_{(s)}(\infty), \tag{19}$$

at $w<0$. Here

$$Q_2(t) = 0.8Bt \int_0^\infty dv v^3 [1+(1+v^2)B^2]^{-2} \sinh^{-1}(t\sqrt{1+v^2}). \quad (20)$$

One can show that $M(w) \sim 1$ at $w=0$ and $M(w) \sim w^{-1}\exp(-w)$ at $w \gg 1$. Equation (17) shows that the MEE is entirely due to the central symmetry breaking because velocity and spin have opposite parity under space inversion.

Since $B \sim V_b/\epsilon_F$, where V_b is the height of the interface barrier, one may estimate $B \sim 1$. The small quantity in Eqs. (17) and (14) 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) can be obtained in the indirect way. For the case of polar crystal, $\alpha_v \mathbf{p}\times\mathbf{c}\cdot\boldsymbol{\sigma} = \Sigma_{nm}\langle c|m^{-1}(-i\mathbf{p}\cdot\boldsymbol{\nabla})|n\rangle (E_n - E_c)^{-1}\langle n|\zeta \mathbf{l}\times\mathbf{s}|m\rangle (E_m - E_c)^{-1}\langle|eE\mathbf{r}\cdot\mathbf{c}|c\rangle$, where $|c\rangle$ refers to the conduction band, $|m\rangle$ and $|n\rangle$ to other bands, and ζ is the intra-atomic SO energy.[5] Evidence was given that for the case of CdS, where $E \simeq 10^7$ V/cm is the electric field that the cadmium ions create on the sulphur ions, ζ is the SO energy of the $2p$ shell of sulphur, and $\alpha_V(\mathrm{CdS}) \simeq 2\times 10^5$ cm s^{-1}, this equation is quantitatively correct.[5] Since a metal inside the double electric layer is similar to a polar metal, the equation is qualitatively applicable. The evaluation of the band structure with allowance for all features of the crystal structure near an interface is a difficult problem. However, for our purpose, it is sufficient to approximate the matrix elements of coordinate and momentum as well as the interband energies by their normal, atomic values, i.e., to take these to be essentially the same as in CdS. Electric fields in contacts are also of the same order of magnitude as in CdS, but in the case of contacts with heavy metals, such as Au or W, the SO coupling ζ is a few tens of times as many as $\zeta(S)$. Accordingly, α_v in these contacts should exceed $\alpha_v(\mathrm{CdS})$ of the same degree. So one gets $\alpha m \sim 10^{-2}$. The polarization per one electron can, therefore, amount to $\alpha m T_c/\epsilon_F - 10^{-5} - 10^{-4}$ in a domain of width ξ_0 on the S side of the interface at a superfluid velocity of the order of T_c/k_F. (On the N side, the effect should be smaller due to reduced value of the singlet gap.) The spin polarization might reveal itself through the nuclear-magnetic-resonance (NMR) frequency shift in N/S sandwiches. If the sandwich is subject to a magnetic field **B**, the Knight shift, being determined by the sum of the magnetization induced by **B** ($\mathbf{M_0} \sim \mathbf{B}$) and that induced by the supercurrent ($\mathbf{M_s} \sim \mathbf{c}\times\mathbf{J_s}$), should have a term proportional to the mixed product $\mathbf{B}\cdot\mathbf{c}\times\mathbf{J_s}$, which is an odd function of the field **B**. Since the different metals have different NMR frequencies, the supercurrent induced magnetization on either side of the contact can be investigated separately. It should be emphasized that the MEE does not depend on the value of the triplet interaction, i.e., it should exist even at pure s-wave pairing when $\lambda_p = 0$. In the general case, there are no reasons for λ_p/λ_s to be anomalously small. So $\Delta_{(t)}/\Delta_{(s)} \sim (\alpha m)[\lambda_p N(0)]$ can well appear to be of the order of 10^{-3} under favorite conditions. In 3D polar superconductors, $\Delta^{(t)}_{\alpha\beta}$ leads to different bound energy of the Cooper pairs with opposite helicities.[10] Possible effects of $\Delta_{(t)}$ on contact phenomena are not a concern here.

In conclusion, we argue that there should be an interface SO coupling that can be considered as the material manifestation of the mirror-symmetry breaking at the plane of the contact. The coupling is inherent in contacts of any conductors and, therefore, should be relevant to many spin-dependent phenomena. In the case of N/S contacts, the coupling has been shown to locally change the superconductivity character on both sides of the contact—the Cooper pair wave function acquires the triplet part. This gives rise to the MEE. The effect makes it possible to measure the value of α by means of the NMR technique as well as to control the nuclear spin polarization at no magnetic field dc current should result in the permanent polarization and ac current should induce spin-flip processes.

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[1] V.M. Edelstein, Phys. Rev. Lett. **75**, 2004 (1995).

[2] V.M. Edelstein, J. Phys.: Condens. Matter **8**, 339 (1996).

[3] G.A. Prinz, Phys. Today **48** (4), 58 (1995); B.E. Kane, Nature (London) **393**, 133 (1998); M. Johnson, Science **260**, 320 (1993).

[4] E.I. Rashba and V.I. Scheka, Fiz. Tverd. Tela (Leningrad) **2**, 162 (1958); R.C. Gasella, Phys. Rev. Lett. **5**, 371 (1960).

[5] R. Romestain, S. Geshwind, and G.E. Devlin, Phys. Rev. Lett. **39**, 1583 (1977).

[6] F.J. Ohkawa and Y. Uemura, J. Phys. Soc. Jpn. **37**, 1325 (1974).

[7] G. Lommer, F. Malcher, and U. Rössler, Phys. Rev. Lett. **60**, 728 (1988).

[8] Note that in the article by V.L. Geshkenbein and A.I. Larkin [JETP Lett. **43**, 395 (1986)] a term of the form of Eq. (1) was hypothesized to be present in the interface transmission amplitude owing to a difference in SO coupling on either side of the interface. If the hypothesis were correct, it would mean that there is the additional reason for the Hamiltonian (1) to exist. However, the hypothesis was never conclusively proved even for a model system.

[9] L.N. Bulaevskii, A.A. Guseinov, and A.I. Rusinov, Zh. Éksp. Teor. Fiz. **71**, 2356 (1976) [Sov. Phys. JETP **44**, 1243 (1976)].

[10] V.M. Edelstein, Zh. EksP. Teor. Fiz, **95**, 2151 (1989) [Sov. Phys. JETP **68**, 1244 (1989)]; Physica C **185**, 1469 (1991).

[11] V.M. Edelstein, Pis'ma Zh. EksP. Teor. Fiz, **69**, 363 (1999) [JETP Lett. **69**, 392 (1999)].

[12] L.P. Gor'kov and E.I. Rashba, Phys. Rev. Lett. **87**, 037004 (2001).

[13] L.P. Gor'kov, Zh. Éksp. Teor. Piz. **37**, 1407 (1960) [Sov. Phys. JETP **10**, 998 (1960)].

[14] See, for example, P.-G. de Gennes, *Superconductivity of Metals and Alloys* (Benjamin, New York, 1966).

[15] P. Deift and E. Trubowitz, Commun. Pure Appl. Math. **32**, 121 (1979).

[16] H. Kleinert, Fortschr. Phys. **26**, 565 (1978).

[17] See, for example, C.-R. Hu and V. Korenman, Phys. Rev. **178**, 684 (1969).