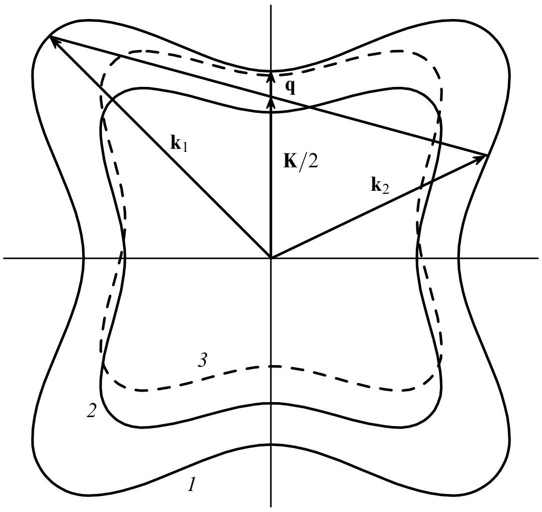

Figure 7. Fermi contours in the case where there is pairing with a finite momentum: *1*, $\varepsilon^{(-)}(k) = \varepsilon_{\mathrm{F}}$; *2*, $\varepsilon^{(+)}(k) = \varepsilon_{\mathrm{F}}$, and *3*, $\varepsilon^{(+)}(k+q) = \varepsilon_{\mathrm{F}}$.

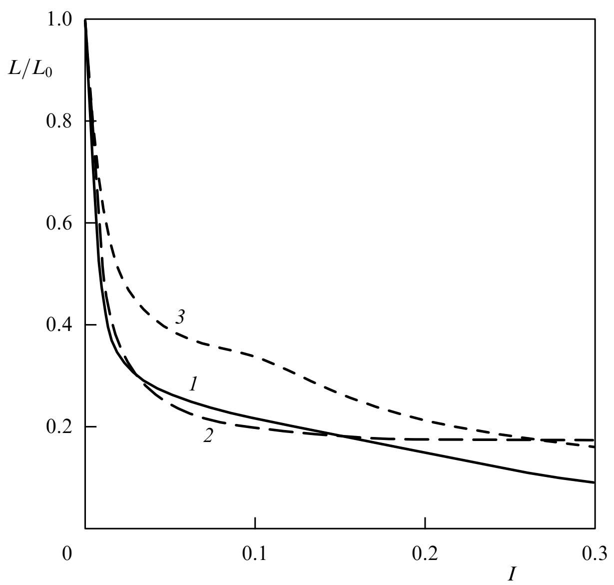

Figure 8. Dependence of the length L/L_0 of a part of the Fermi contour over which $\varepsilon(K/2+k) = \varepsilon(K/2-k)$, on the magnetization I for pairing with a finite momentum: *1*, $n = 0.86$; *2*, $n = 1.00$, and *3*, $n = 1.14$.

from the one discussed earlier for $K = 0$ (in the sense that this length will enter the self-consistency equation in a different way), the ratio L/L_0 still characterizes the suppression of superconductivity by ferromagnetism. The kinematic constraint region occupies a section of the Fermi surface within a certain sector directed along $\mathbf{K} = (0, K)$. Hence, when calculating L (and L_0), we must limit ourselves to this sector. As a result, the value of L_0 proves to be smaller than the value obtained above at $K = 0$, while the ratio L/L_0 is larger. The dependences of L/L_0 on I for different occupancies are shown in Fig. 8. Comparing Figs 5 and 8, we see that the situation of pairing with a finite $\mathbf{K}$ is preferable from the point of view of coexistence of ferromagnetism and superconductivity.

(6) Lee and Kim [13] studied the effect that a change in magnetization (caused by variations in niobium concentration) has on the magnetic properties of $Ru_{1-x}Nb_xSr_2Eu_{1.5}Ce_{0.5}Cu_2O_z$ samples. As magnetization decreases with increasing niobium concentration, a diamagnetic response appears when the sample is cooled in a magnetic field (the Meissner–Ochsenfeld effect). This result has been interpreted as the formation of a spontaneous vortex structure under growing magnetization [14]. Within the model examined in the present work, the emergence of a spontaneous vortex state is due to the appearance of a nonuniformity ($q \neq 0$) and is determined by condition (10). At lower magnetizations, a uniform vortex-free superconducting state, corresponding to ideal diamagnetism, proves to be preferable.

The authors would like to express their gratitude to V L Ginzburg for discussing the problems touched upon in this paper. The work was made possible by the financial support of the Russian Science and Education Program 'Integration' (projects AO 133 and AO 155), the Russian Foundation for Basic Research (grant No. 02-02-17133), and the Federal Science and Technology Target Program 'Research and Development in Priority Areas of Science and Technology' (state contract No. 40.072.1.1.1173).

References

1. Ginzburg V L *Zh. Eksp. Teor. Fiz.* **31** 202 (1956) [*Sov. Phys. JETP* **4** 153 (1957)]
2. Bardeen J, Cooper L N, Schrieffer J R *Phys. Rev.* **106** 162 (1957)
3. Sarma G *J. Phys. Chem. Solids* **24** 1029 (1963)
4. Anderson P W, Suhl H *Phys. Rev.* **116** 898 (1959)
5. Buzdin A I et al. *Usp. Fiz. Nauk* **144** 597 (1984) [*Sov. Phys. Usp.* **27** 927 (1984)]
6. Larkin A I, Ovchinnikov Yu N *Zh. Eksp. Teor. Fiz.* **47** 1136 (1964) [*Sov. Phys. JETP* **20** 762 (1965)]
7. Fulde P, Ferrell R A *Phys. Rev.* **135** A550 (1964)
8. Felner I et al. *Phys. Rev. B* **55** R3374 (1997)
9. Jorgensen J D et al. *Phys. Rev. B* **63** 054440 (2001)
10. Pickett W E, Weht R, Shick A B *Phys. Rev. Lett.* **83** 3713 (1999)
11. Shimahara H, Hata S *Phys. Rev. B* **62** 14541 (2000)
12. Belyavskiĭ V I, Kapaev V V, Kopaev Yu V *Zh. Eksp. Teor. Fiz.* **118** 941 (2000) [*JETP* **91** 817 (2000)]
13. Lee H K, Kim Y C *Int. J. Mod. Phys. B* **17** 3682 (2003)
14. Sonin E B, Felner I *Phys. Rev. B* **57** R14000 (1998)

PACS numbers: **74.45.+c**, **74.50.+r**
DOI: 10.1070/PU2004v047n09ABEH001877

Theory of magnetic contacts between clean superconductors

Yu S Barash, I V Bobkova, T Kopp

The problem of the coexistence of superconductivity and ferromagnetism is one of the most interesting in the physics of condensed matter and over the years has attracted a lot of attention, beginning with Ginzburg's paper [1]. A special case is the problem of the coexistence and the mutual effect of spatially separated and adjacent superconducting and ferromagnetic phases. Such a statement of the problem includes research into contacts between superconductors and magnetic substances and, in particular, superconducting contacts through ferromagnetic interlayers. At present it is a well-known fact that under certain conditions such contacts constitute what is known as π-contacts. The possibility of forming a π-contact due to the magnetic properties of the interlayer was recognized by theoreticians more than a quarter of a century ago [2], while the first specific example considered by Buzdin et al. [3] was a completely transparent superconductor – ferromagnet – superconductor junction (an SFS junction). The physical meaning of the term π-contact can be explained as follows. The Josephson current j is proportional to the first derivative of the thermodynamic potential Ω with respect to the phase difference χ between the superconductors:

$$j(\chi) = \frac{e}{\hbar} \frac{\partial \Omega}{\partial \chi} .$$

Here, χ is the phase difference of the superconducting order parameters. Thus, the thermodynamic potential reaches its minimum for a phase difference at which the current vanishes. In many cases, say in tunneling contacts, the dependence of the current on phase difference is rather simple:

$$j(\chi) = j_c \sin \chi .$$

In accordance with such a dependence and the relationship between the superconducting current and the thermodynamic potential, the ground state of the junction with $j_c > 0$ is realized when the phase difference $\chi = 0$. Such a contact is called a 0-contact. But if $j_c < 0$, the thermodynamic potential minimum occurs at $\chi = \pi$. In this case, we speak of a π-contact. Most often the 0-contacts are realized, as happens in ordinary isotropic superconductors separated by a nonmagnetic interlayer. Although the formation of π-contacts was first predicted theoretically for superconductors separated by a ferromagnetic metal, experimentally such junctions were first fabricated for ordinary nonmagnetic interlayers in corner tunneling contacts and in SQUIDs with high-temperature superconductors [4, 5]. In this case, the formation of a π-contact is related to the sign-changing nature of the order parameter in high-temperature superconductors as a function of the direction of the quasi-particle momentum. π-contacts with ordinary superconductors and ferromagnetic interlayers have been fabricated only recently and were actively investigated through experiments [6 – 9]. Real measurements usually involve 'dirty' superconductors. There are also theoretical investigations into such systems [10 – 13]. Contacts between clean superconductors, which are discussed below, are also of strong interest. The reason is not only that such contacts can be fabricated but also that the emerging effects are of general interest from the viewpoint of physics. The unusual properties of π-contacts have also attracted much attention in connection with the possibility of using them in superconducting electronics [8, 14 – 17]. We mean, of course, π-SQUIDs [8, 15] and the elements of quantum computers [16, 17].

The physical mechanism leading to a $0-\pi$ junction in transparent contacts through a ferromagnetic interlayer is based on the proximity effect in the interlayer. The exchange field in a ferromagnetic metal sandwiched by two superconductors induces oscillations against the background of the exponentially decaying wave function of Cooper pairs deep into the interlayer. Depending on the thickness of the ferromagnetic layer, the signs of the wave function of the Cooper pairs on its opposite boundaries may coincide or differ. In the first case we have a 0-contact, while in the second a π-contact. As a result, the critical Josephson current proves to be an oscillating function of the interlayer thickness [3]. Oscillatory effects emerge in most problems associated with the coexistence of superconductivity and magnetism. The paramagnetic interaction of the quasi-particle spins and the magnetic (exchange or external) field brings about a difference in the momenta of the quasi-particles of equal energy but with oppositely directed spins. For the electrons forming a Cooper pair this leads to the pair acquiring a center-of-mass momentum $\mathbf{q}$ and, hence, to an oscillating factor $\exp(i\mathbf{qR})$ in the pair's wave function. These oscillations are similar to the spatial oscillations of the order parameter in the FFLO (Fulde – Ferrell – Larkin – Ovchinnikov) superconducting state brought about by the paramagnetic interaction of the superconducting electrons with an external magnetic field [18, 19].

When there is an interlayer of a ferromagnetic insulator or semiconductor, the proximity effect shows itself much weaker than for a metal interlayer. It has been demonstrated through experiments that a ferromagnetic semiconductor creates a spin-dependent potential barrier for the quasi-particles in tunneling contacts, with the result that the probabilities of quasi-particles with spin up and down penetrating through this barrier differs [20, 21]. The proximity effect is also negligible for an interlayer of a ferromagnetic metal if there are tunneling potential barriers at the interlayer's boundaries. And the question now arises of the properties of the Josephson current in junctions through magnetic interlayers in the absence of a proximity effect, with the properties of the ferromagnetic interlayer being the same as in the case where the interlayer is surrounded by normal metals.

Assuming that the interlayer thickness is smaller than the superconducting coherence length, we will describe the properties of the interlayer within the framework of the S-matrix approach. When dealing with contacts, we are speaking of the reflection and transmission amplitudes of particles with spins up and down. In the case of uniform magnetization in the interlayer, the spin states parallel and antiparallel to the magnetization axis are good quantum numbers, i.e., they are conserved under reflection by or penetration through the interlayer. Hence, if we select the z-axis along the magnetization direction, the reflection and transmission amplitudes are diagonal in the spin space. Let us first examine the behavior of the Josephson current in magnetic contacts with low transparencies $D_\uparrow$ and $D_\downarrow$. In the approximation that is linear in D, the proximity effects have no influence on the Josephson current, with the result that the above oscillation effects cannot manifest themselves in the Josephson current. But will the behavior of the Josephson current in tunneling magnetic contacts be any different from that in the nonmagnetic case? The answer to this question is yes, and it contains an important physical characteristic of the magnetic interlayer to which the Josephson current is sensitive.

As is well known, the properties of the interlayer in the junction are ordinarily characterized by a single quantity in the expression for the Josephson current — that is, the transparency coefficient of the potential barrier. It turned out that in the case of a ferromagnetic interlayer the Josephson current is extremely sensitive to the gauge-invariant phase difference Θ of the amplitudes of reflection of quasi-particles with spins up and down from the interlayer: $r_{\uparrow(\downarrow)} = |r_{\uparrow(\downarrow)}| \exp(i\Theta_{\uparrow(\downarrow)})$, where $\Theta = \Theta_\uparrow - \Theta_\downarrow$. Tokuyasu et al. [22] and Fogelström [23] were the first to introduce Θ explicitly in the case of an impenetrable magnetic surface for the problems relation to the Josephson current through magnetic interlayers. There is also another parameter $\alpha = \pm 1$ characterizing the magnetic interlayer, whose sign directly determines the sign of the critical Josephson current [24]. In the nonmagnetic case, $\alpha = -1$. For a high 'purely

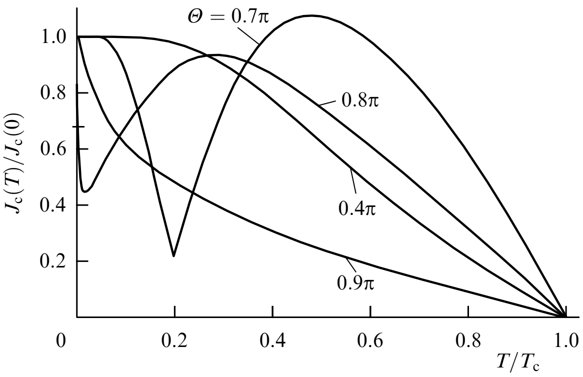

Figure 1. Critical Josephson current $J_c(T)$ for different values of Θ, normalized to its value $J_c(0)$ at zero temperature. The transparency coefficients $D_\uparrow = D_\downarrow = 0.1$ and $\alpha = -1$.

magnetic barrier', when $\hat{V} = h\sigma_z$, we have $\alpha = 1$. For the model of a rectangular magnetic barrier, where $\hat{V} = V + h\sigma_z$ in the interlayer, we obtain

$$\alpha = -\mathrm{sgn}\left[\left(V - \frac{p_{\mathrm{F},x}^2}{2} + h\right)\left(V - \frac{p_{\mathrm{F},x}^2}{2} - h\right)\right],$$

where p_{F} is the Fermi momentum. Thus, the equality $\alpha = 1$ is attained if the wave function for quasi-particles with one spin falls off exponentially within the barrier, while that for particles with the other spin oscillates. Generally, α can depend on the direction of the quasi-particle momentum.

The expression for the Josephson current $J(T, \chi) = J_c(T)\sin\chi$ flowing through a tunneling ferromagnetic interlayer has the form [24]

$$J_c(T) = -\alpha e|\Delta|\sqrt{D_\uparrow D_\downarrow}\left[\frac{\varepsilon_{\mathrm{B},0}(\Theta)}{|\Delta|}\tanh\frac{\varepsilon_{\mathrm{B},0}(\Theta)}{2T} - \frac{|\Delta|}{2T}\left(1 - \frac{\varepsilon_{\mathrm{B},0}^2(\Theta)}{|\Delta|^2}\right)\cosh^{-2}\frac{\varepsilon_{\mathrm{B},0}(\Theta)}{2T}\right], \qquad (1)$$

with $\varepsilon_{\mathrm{B},0}(\Theta) = |\Delta\cos(\Theta/2)|$. In the limit of a nonmagnetic interlayer, $\Theta = 0$, $\varepsilon_{\mathrm{B},0}(\Theta) = |\Delta|$, $D_\uparrow = D_\downarrow$, and equation (1) acquires the form of the standard Ambegaokar–Baratoff result. When $\alpha = -1$ and $D_\uparrow = D_\downarrow$, formula (1) coincides with the one derived by Chtchelkatchev et al. [25]. In the case of a magnetic interlayer, the temperature behavior of the critical Josephson current strongly depends on the magnitude of the parameter Θ. When $\Theta < \pi/2$, the temperature dependence is close to that describing the conventional behavior of $J_c(T)$. Expression (1) contains two terms with opposite signs, so that for $\Theta > \pi/2$ the current, as a function of temperature, is alternating. At a certain temperature, the tunnel current vanishes and there occurs a $0-\pi$ transition in the contact. Then, in the case of negative α, a π-contact is realized near T_c, while at low temperatures a 0-contact is formed.

One can go beyond the tunnel approximation in finding the Josephson current. This can be done analytically, as in Refs [24–26]. The corresponding behavior of the critical current is depicted in Fig. 1.

Since the absolute value of the critical current is usually measured in experiments, we have used this quantity in the figure. Clearly, the critical current at the $0-\pi$-transition point does not vanish completely, and the inflection point is actually a reflection of the abrupt change in the sign of the critical current. Fogelström [23] was the first to discover this behavior (it was done by numerical calculations). A distinct minimum in the function $J_c(T)$ occurs only for contacts with low transparency. In tunnel junctions, the minimum value of the critical current is proportional to the square of the barrier's transparency, and the current–phase relationship contains the first and second harmonics of the same order of magnitude. At the temperature where the term in the current that is linear in transparency vanishes, the Josephson current turns out to be proportional to $\sqrt{D_\uparrow D_\downarrow}(D_\uparrow + D_\downarrow)\sin\chi + \alpha D_\uparrow D_\downarrow \sin 2\chi$ [24].

The physics of the Josephson effect is determined by the participation of interface Andreev bound states in the transport of the superconducting current across the contact [27–29]. In conventional symmetric Josephson junctions, the entire Josephson current is carried by these states. When magnetic contacts with uniform magnetization are involved, the Andreev bound states have a spin polarization that is either parallel or antiparallel to the magnetization direction. The corresponding energies of these spin-polarized Andreev states, which are degenerate for nonmagnetic interlayers, differ in the case of ferromagnetic interlayers, where $\Theta \neq 0$. The appearance of a $0-\pi$ transition in magnetic contacts between clean superconductors can be considered as the result of competition between the contributions of Andreev bound states with opposite spin polarizations. These contributions to the formation of the current–phase relations are shown separately in Fig. 2.

When the conditions for a minimum in the critical current are met, the energies of the Andreev states with oppositely directed spins have different signs. At zero temperature, only states with negative energies are occupied, with the result that only Andreev states with appropriate spin polarizations contribute to the Josephson current. The contribution of Andreev states with the other spin polarization increases with a rise in temperature together with the growth of the occupation of levels with positive energies. This contribution becomes important at temperatures on the order of the bound-state energy and results in a Josephson current flowing in the opposite direction compared to the case of the contribution of states with negative energies. The competition between Andreev states with opposite spin polarizations leads to a minimum in the Josephson current at temperatures lower than T_c and, hence, to a $0-\pi$ transition in the contact. It has proven useful to compare such behavior of the Josephson current in SFS junctions with that in conventional junctions, in which at low temperatures the contribution of Andreev levels with negative energy (which, in addition, are spin-degenerate) is also predominant. The occupation of Andreev bound states with positive energies increases with temperature, which leads to a reduction in the total Josephson current and, in addition to the temperature dependence of the order parameter, provides an important contribution to the formation of the temperature dependence of the current. In contrast to SFS junctions, the competing current carried by Andreev states with positive energies never exceeds the contribution of levels with negative energies in the nonmagnetic cases, and the two contributions vanish at the transition temperature.

The contributions of Andreev bound states with positive and negative energies to the Josephson current may be observed separately in nonequilibrium conditions. In particular, the nonequilibrium occupation of states that is stimulated by a microwave field may lead to an enhancement (reduction) in the contribution to the Josephson current when this occupation exceeds (is smaller than) the equilibrium

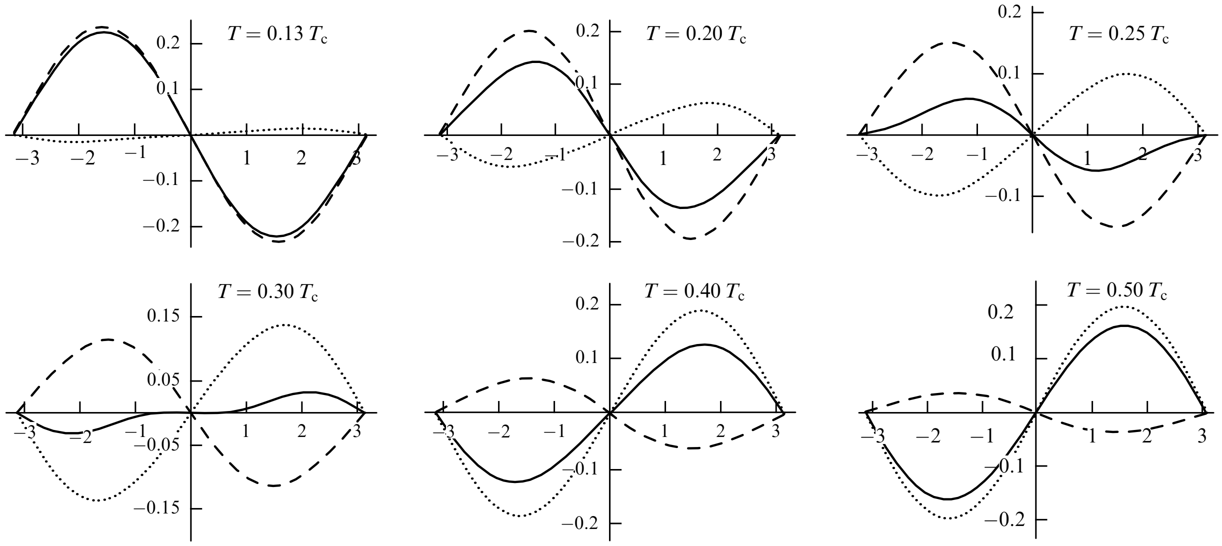

Figure 2. Current – phase relations at different temperatures for the total Josephson current (solid curves) and for separate contributions to the current from Andreev states with spin up (dashed curves) and spin down (dotted curves). The values of the parameters are $D_{\uparrow} = D_{\downarrow} = 0.05$ and $\Theta = 2\pi/3$. The current is normalized to the magnitude of the critical current in the nonmagnetic case $\Theta = 0$ at zero temperature.

population. This significantly disturbs the balance of currents from positive and negative levels already in the nonmagnetic case [30 – 32]. When the positive and negative levels in an SFS junction have different spin polarizations, the stimulation of transitions between these levels must be accompanied by spin flip. Observations of stimulated transitions between levels with only spin flip would make it possible to identify the above-noted difference in the spin polarizations of Andreev states with positive and negative energies, change the population of these states, and establish their relative role in the formation of a Josephson current.

Now let us go back to the above-established basic condition $\Theta > \pi/2$, which must be met for tunneling π-contacts to exist in SFS systems. The parameter Θ is a characteristic of the ferromagnetic material if the characteristic length of quasi-particle decay does not exceed the interlayer width. This is usually the case with interlayers made of ferromagnetic insulators and semiconductors. Achieving fairly large values of the parameter Θ then becomes a task for materials science. Simple estimates show that such values of Θ can be achieved under rather restrictive, and yet realistic, conditions. For instance, if a homogeneous ferromagnetic insulator is used as the interlayer, the condition $\Theta > \pi/2$ can be realized only if the exchange field h, the energy gap ε_g, and the Fermi energy ε_F are of the same order of magnitude. More promising are the conditions under which Θ depends not only on the properties of the ferromagnetic material of the interlayer but also on the geometric parameters (on the interlayer thickness, in the simplest case). Such behavior is characteristic of interlayers fabricated from ferromagnetic metals with, possibly, potential barriers. For instance, for the interlayer of a ferromagnetic material or even more complex structures consisting of two ferromagnetic layers separated by a nonmagnetic insulating barrier, we have $\Theta = 2lh/v_F$. Here l is the total thickness of the ferromagnetic metallic interlayer. This implies that $\Theta \sim \pi$ for $l \sim \xi\Delta/h \ll \xi$, provided that $\Delta \ll h \ll \varepsilon_F$, viz. conditions which are often met.

In an FIF interlayer consisting of two layers of a ferromagnetic metal separated by a nonmagnetic insulating barrier, the magnetizations that lie in the plane of the layers may form an angle φ with each other, and the Josephson current is very sensitive to this angle. At $\varphi = 0$, when the magnetizations of the two ferromagnetic layers are parallel, the problem is equivalent to that of a junction with a uniform magnetization, which we considered earlier. Thus, the tunneling contact with $\varphi = 0$ for $\Theta > \pi/2$ is a π-contact. But if $\varphi = \pi$, the magnetizations of the two layers are antiparallel, and the junction in question will never change to a π-state, no matter what the temperature and barrier transparency are [33]. This statement holds even for dirty superconductors, where the Andreev states are completely smeared [12, 34]. As a result, for $\Theta > \pi/2$, the $0-\pi$ transition in the junction in question proceeds not only under temperature variations at small values of φ, but at a fixed temperature under misorientation-angle variations as well [33].

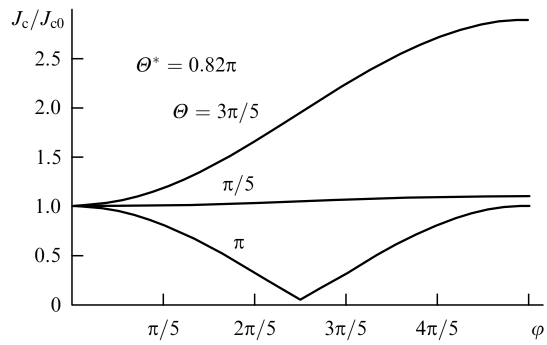

Figure 3. The critical Josephson current as a function of angle φ, normalized to its value at $\varphi = 0$. In the particular case of parallel magnetizations, the junction at a given temperature T is in a π-state when $\Theta > \Theta^*(T, D)$. The cusp in the curve corresponds to a $0-\pi$ transition that occurs at a certain value of the misorientation angle.

Figure 3 shows the dependence of the critical Josephson current at $T = 0.1T_c$ on the magnetization-misorientation angle φ for different values of Θ and at a junction transparency $D = 0.01$.

Clearly, there are two distinct regimes for the critical Josephson current as a function of the misorientation angle. The two regimes are separated by the characteristic value $\Theta^*(T, D)$ which depends on the temperature and transparency of the junction. For $\Theta < \Theta^*$, the current is a monotonic function of the misorientation angle and reaches its maximum value when the magnetizations are antiparallel. At the same time, for $\Theta > \Theta^*$ the current is not a monotonic function of φ. It has a distinct minimum at a certain intermediate value of φ (at which the $0-\pi$ transition takes place), and a maximum at $\varphi = \pi$. When $\Theta = \pi$, the currents at $\varphi = 0$ and $\varphi = \pi$ are equal. The parameter Θ^* is related to the properties of the junction at $\varphi = 0$. In a junction with parallel magnetizations of the ferromagnetic layers ($\varphi = 0$) and with $\Theta = \Theta^*(T, D)$, the $0-\pi$ transition can be shown to occur exactly at the temperature T considered. Hence, for $\Theta > \Theta^*(T, D)$, the equilibrium state of the junction with $\varphi = 0$ at the temperature T is a π-state, while for $\Theta < \Theta^*(T, D)$ it is a 0-state.

The dependence of the Josephson current on the misorientation angle φ becomes especially simple in the tunneling limit. In tunneling quantum-point contacts with the interlayer under consideration, the Josephson current has the form $J(T, \varphi, \chi) = J(T, \varphi) \sin \chi$, where

$$J(T, \varphi) = J^{(\mathrm{p})}(T) \cos^2 \frac{\varphi}{2} + J^{(\mathrm{a})}(T) \sin^2 \frac{\varphi}{2} \,. \tag{2}$$

The quantity $J^{(\mathrm{p})}(T) \equiv J(T, \varphi = 0)$ is described by expression (1), while $J^{(\mathrm{a})}(T) \equiv J(T, \varphi = \pi)$ has the following form

$$J^{(\mathrm{a})}(T) = \frac{eD|\Delta|}{\cos(\Theta/2)} \tanh \frac{|\Delta| \cos(\Theta/2)}{2T} \,. \tag{3}$$

Thus, $\left|J^{(\mathrm{p})}(T)\right|$ and $\left|J^{(\mathrm{a})}(T)\right|$ are the critical currents in tunneling contacts with parallel and antiparallel orientations of the exchange fields in a three-layer interface.

Expression (2) describing the dependence of the Josephson current on the magnetization-misorientation angle can be applied in the tunneling approximation under very broad assumptions, but cannot be used for highly transparent junctions. It was first derived on the assumption that no Andreev bound states emerge in the system, when the quantities $J^{(\mathrm{p})}(T)$ and $J^{(\mathrm{a})}(T)$ always have the same sign and, hence, there is no $0-\pi$ transition [35]. Expression (2) shows that under a change in φ a $0-\pi$ transition occurs if $J^{(\mathrm{p})}(T)$ and $J^{(\mathrm{a})}(T)$ have opposite signs. Only in this case will the Josephson current be a nonmonotonic function of φ.

A $0-\pi$ transition triggered by variations in the misorientation angle could be used for switching a junction from the π-state to the 0-state at a fixed temperature. Indeed, if the coercive force in one ferromagnetic layer is much larger than in the other, the mutual orientation of the magnetizations can be changed by switching on an external magnetic field, rotating it through a certain angle, and then switching it off. Here, the value of the field should be selected in such a way that it can rotate the magnetization of only one ferromagnetic layer, i.e., the layer with the lower coercive force.

The present work was made possible by the financial support of the Russian Foundation for Basic Research (grant No. 02-02-16643), the Physical Sciences Division of the Russian Academy of Sciences, the Ministry of Science and Education of the Russian Federation (Yu S B and I V B), the 'Dynasty' Foundation (I V B), and Forschungszentrum Jülich (Landau Scholarship; I V B). T Kopp is grateful for the financial support from SFB (grant No. 484) and the BMBF Project No. 13N6918A.

References

1. Ginzburg V L *Zh. Eksp. Teor. Fiz.* **31** 202 (1956) [*Sov. Phys. JETP* **4** 153 (1957)]
2. Bulaevskiĭ L N, Kuziĭ V V, Sobyanin A A *Pis'ma Zh. Eksp. Teor. Fiz.* **25** 314 (1977) [*JETP Lett.* **25** 290 (1977)]
3. Buzdin A I, Bulaevskiĭ L N, Panyukov S V *Pis'ma Zh. Eksp. Teor. Fiz.* **35** 147 (1982) [*JETP Lett.* **35** 178 (1982)]
4. Van Harlingen D J *Rev. Mod. Phys.* **67** 515 (1995)
5. Tsuei C C, Kirtley J R *Rev. Mod. Phys.* **72** 969 (2000)
6. Ryazanov V V et al. *Phys. Rev. Lett.* **86** 2427 (2001)
7. Kontos T et al. *Phys. Rev. Lett.* **89** 137007 (2002)
8. Guichard W et al. *Phys. Rev. Lett.* **90** 167001 (2003)
9. Bauer A et al. *Phys. Rev. Lett.* **92** 217001 (2004)
10. Buzdin A I, Vujicic B, Kupriyanov M Yu *Zh. Eksp. Teor. Fiz.* **101** 231 (1992) [*Sov. Phys. JETP* **74** 124 (1992)]
11. Koshina E A, Krivoruchko V N *Pis'ma Zh. Eksp. Teor. Fiz.* **71** 182 (2000) [*JETP Lett.* **71** 123 (2000)]; *Phys. Rev. B* **63** 224515 (2001)
12. Golubov A A, Kupriyanov M Yu, Fominov Ya V *Pis'ma Zh. Eksp. Teor. Fiz.* **75** 223 (2002) [*JETP Lett.* **75** 190 (2002)]
13. Golubov A A, Kupriyanov M Yu, Fominov Ya V *Pis'ma Zh. Eksp. Teor. Fiz.* **75** 709 (2002); **76** 268 (2002) [*JETP Lett.* **75** 588 (2002); **76** 231 (2002)]
14. Terzioglu E, Beasley M R *IEEE Trans. Appl. Supercond.* **8** (2) 48 (1998)
15. Schulz R R et al. *Appl. Phys. Lett.* **76** 912 (2000)
16. Ioffe L B et al. *Nature* **398** 679 (1999)
17. Blatter G, Geshkenbein V B, Ioffe L B *Phys. Rev. B* **63** 174511 (2001)
18. Larkin A I, Ovchinnikov Yu N *Zh. Eksp. Teor. Fiz.* **47** 1136 (1964) [*Sov. Phys. JETP* **20** 762 (1965)]
19. Fulde P, Ferrell R A *Phys. Rev.* **135** A550 (1964)
20. Tedrow P M, Tkaczyk J E, Kumar A *Phys. Rev. Lett.* **56** 1746 (1986)
21. Hao X, Moodera J S, Meservey R *Phys. Rev. B* **42** 8235 (1990)
22. Tokuyasu T, Sauls J A, Rainer D *Phys. Rev. B* **38** 8823 (1988)
23. Fogelström M *Phys. Rev. B* **62** 11812 (2000)
24. Barash Yu S, Bobkova I V *Phys. Rev. B* **65** 144502 (2002)
25. Chtchelkatchev N M et al. *Pis'ma Zh. Eksp. Teor. Fiz.* **74** 357 (2001) [*JETP Lett.* **74** 323 (2001)]
26. Cuevas J C, Fogelström M *Phys. Rev. B* **64** 104502 (2001)
27. Furusaki A, Tsukada M *Physica B* **165–166** 967 (1990)
28. Furusaki A, Tsukada M *Phys. Rev. B* **43** 10164 (1991)
29. Beenakker C W J, van Houten H *Phys. Rev. Lett.* **66** 3056 (1991)
30. Gorelik L Y et al. *Phys. Rev. Lett.* **75** 1162 (1995)
31. Gorelik L Y et al. *Phys. Rev. Lett.* **81** 2538 (1998)
32. Lundin N I *Phys. Rev. B* **61** 9101 (2000)
33. Barash Yu S, Bobkova I V, Kopp T *Phys. Rev. B* **66** 140503(R) (2002)
34. Krivoruchko V N, Koshina E A *Phys. Rev. B* **64** 172511 (2001)
35. Bergeret F S, Volkov A F, Efetov K B *Phys. Rev. Lett.* **86** 3140 (2001); *Phys. Rev. B* **64** 134506 (2001)

PACS numbers: **74.20.–z**, **74.72.–h**
DOI: 10.1070/PU2004v047n09ABEH001878

High-temperature superconductivity today

E G Maksimov

In the present paper, I briefly discuss the main ideas underlying the problems covered in the report.

First, on the basis of the existing experimental optical and angle-resolved photoemission spectroscopic (ARPES) data it has been shown that a standard phase diagram of high-T_c superconducting compounds contains a totally inaccurate representation of the number of carriers along the horizontal axis. In particular, the optimum number x of carriers in such