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

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

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

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

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.

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High-temperature superconductivity today

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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 diagrams, corresponding to the maximum value of T_c, is assumed to be equal to 0.16 holes per unit cell. Actually, as a comparison between the experimental data obtained from ARPES measurements [1] and the results of theoretical calculations [2] has shown, the real number of carriers in optimally doped systems is much larger. It most likely corresponds to a number of carriers equal to $1 - x = 0.84$. Moreover, according to the results of optical measurements [3 – 5] and ARPES experiments [1], the total number of carriers in the existence domain of the superconducting phase changes to a much lesser extent than the standard phase diagram implies. All these data show that high-T_c systems exhibiting superconductivity are 'located' far from the doped Mott insulator.

As the doping drops below the optimal level, the electronic structure in the system undergoes substantial changes and an anisotropic pseudogap emerges. The scale of this pseudogap at low doping levels (but within the existence domain of the superconductivity) may be much larger than the scale of the superconducting gap, but it is much smaller than that of the insulating gap in the absence of doping. There is a series of experimental indications [6] pointing to the presence of a certain pseudogap in the electron excitation spectrum and to the fact that, among other things, this gap lowers the density of states at the Fermi surface. The nature of this pseudogap is the subject of many discussions up till now; furthermore, from the experimental angle, the observations also contradict each other. For instance, no traces of a pseudogap have been discovered in most tunneling experiments [7]. There are also no changes in the optical sum rules; such changes would indicate that there is a redistribution of electronic states from the pseudogap region to the region of higher energies [4]. Currently, experiments in scanning tunnel spectroscopy (STS) reported by Pan et al. [8] and McElroy et al. [9, 10] have revealed a new and somewhat unexpected aspect in the pseudogap problem. The researchers have found that in underdoped superconducting cuprates there is a spatial separation of the regions where the pseudogap and superconducting gap are observed. More than 30 years ago Nagaev [11] predicted the possibility of such 'separation' in semiconducting antiferromagnetic systems with doping. The separation in superconducting cuprates, however, does not fit into such a simple scheme, since it is also observed at high enough doping levels, where there are no traces of antiferromagnetism.

In discussing the properties of the normal state in the report, I focused on the recent experimental evidences of the existence of a fairly strong electron – phonon interaction (EPI) in superconducting cuprates. The dispersion of electronic excitations has been measured in ARPES experiments (the dispersion manifests itself in the momentum dependence of the electronic excitations) and it has been found that the electron mass is renormalized because of the interaction of the electron with some sort of bosons whose energy is on the order of phonon energies. The discussion that followed this discovery and concerned the nature of these bosons was extremely intense. Cuk et al. [13] and Zhou et al. [14] proved without any doubt that the nature of the renormalization of the electron mass in superconducting cuprates is related precisely to EPI. The electron – phonon coupling constant found by these researchers is large enough and amounts to about 1, decreasing as the doping level increases. Earlier indications of the existence of a strong EPI were obtained from processing the experimental data on the optical spectra of high-T_c systems. These results, discussed in detail in the review [15], show that EPI must be taken into account in any consistent theory of superconductivity in high-T_c systems.

One of the most important achievements in the studies on the nature of the superconducting state in cuprates (the achievements obtained in the period that followed the publication of Ref. [15]) is the experimental evidence of the proximity of the superconducting state in high-T_c systems to the Bardeen – Cooper – Schrieffer (BCS) model. McElroy et al. [9] (STS experiments) and Matsui et al. [16] (ARPES experiments) measured the excitation spectra of quasi-particles in the superconducting state and found that these excitations constitute a coherent mixture of electrons and holes. It is this type of quasi-particle excitation that Bogolyubov et al. [17] derived from exact diagonalization of the Hamiltonian of the BCS model. Of course, this does not mean that the superconducting state in cuprates is identical to that of the simple BCS model. First, we know that the superconducting order parameter in cuprates is highly anisotropic. Second, since a very strong EPI exists in these systems (as shown earlier), the order parameter is also energy-dependent and the Cooper pairs have a finite lifetime. Unfortunately, so far there is no consistent and realistic theory of superconductivity that would allow for all these factors. Further research is required in this area and new and fruitful discoveries in this field are sure to appear.

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