Inhomogeneous state and the anisotropy of the upper critical field in layered superconductors with Josephson layer interaction

L. N. Bulaevskii

P. N. Lebedev Physics Institute
(Submitted April 27, 1973)

Anisotropy of the upper critical field \( H_{c2} \) in layered superconductors with Josephson interaction of the layers is investigated. It is shown that in pure (along the layer) superconductors at low temperatures an inhomogeneous state is realized for field directions that are close to the layer direction, provided that the Cooper pairs in the field \( H_{c2} \) are not in the lower Landau orbit. The angular dependence of \( H_{c2} \) at \( T = 0 \) and the temperature dependence of \( H_{c2} \) (1) are obtained for pure superconductors. The dependence of \( H_{c2} \) (1) on temperature and purity of the crystal is investigated. It is shown that in dirty superconductors the transition from the normal to superconducting state at low temperatures may be a first-order transition for field directions close to a parallel direction and a second-order transition for other directions. The experimental data for \( \text{TaS}_2(\text{Py})_{1/2} \) are analyzed and it is shown that Josephson interaction of layers occurs in this compound.

1. INTRODUCTION

It has already been noted \(^{[1,2]} \) that in the case of intercalation of layered compounds of the \( \text{TaS}_2 \) with molecules, superconductors with Josephson interaction between layers can be obtained. In the case of the hopping mechanism of conductivity between layers, an interaction of this type is realized if the following condition is satisfied:

\[
\hbar / \tau_s < \Delta (T),
\]

where \( \tau_s \) is the time between the hops of the electrons from one layer to the nearest neighboring one, and \( \Delta (T) \) is the superconducting gap at the temperature \( T \). Superconductors in which condition (1) is satisfied will be called layered superconductors with Josephson interaction of the layers (LSJI).

The value of the lower critical field \( H_{c1}(0) \) for the direction parallel to the layers and the structure of the vortical state in the LSJI were obtained earlier \(^{[1]} \). It is shown in the same reference that the anisotropy of \( H_{c1} \) and other features of the magnetic properties of LSJI in weak fields are due to the nonlinear dependence of the Josephson current between the layers on the vector potential \( A \). In the present article we investigate the anisotropy of the upper critical field of LSJI. An important factor in the mechanism whereby the nuclei of the superconducting phase are produced is the quasi-two-dimensional character of the LSJI \(^{[3]} \), which is also connected with the Josephson interaction of the layers. It was shown earlier \(^{[1]} \) that if the condition (1) is satisfied near \( T_c \) (when \( T_c - T \ll T_c \)), then the field \( H_{c2} \) can be obtained without allowance for the motion of the electrons between the layers. It is clear that this conclusion is also valid in the region of lower temperatures. Therefore, if the condition (1) is satisfied, then the motion of the electrons in the LSJI can be regarded as two-dimensional in the self-consistent field approximation in calculation of the parameters that determine the appearance of the superconductivity. We note that by going outside the framework of the self-consistent-field approximation, i.e., by taking the phase fluctuations into account, we arrive, in a truly two-dimensional system, to a destruction of the superconducting long-range order. This does not occur for LSJI, since the phase fluctuations are suppressed by the Josephson interaction in the layers. Within the framework of the two-dimensional motion of the electrons the self-consistent-field approximation turns out to be valid because the phase fluctuations exert an essential influence on the superconducting properties only in a very close vicinity of the transition point, at practically arbitrarily small interaction between the layer \(^{[4,5]} \).

For the case of a field parallel to the layers, in a pure superconductor (mean free path inside the layer \( l \gg \xi_0 \), the orbital motion of the electrons is essential, and \( H_{c2} \) is determined only by the paramagnetic effect. Then, at temperatures \( T < 0.55T_c \), an inhomogeneous state is realized \(^{[6]} \), and at \( T = 0 \) we have

\[
H_{c2} (1) = \sqrt{2J_2} - \Delta (0) / \mu_s, \text{ where } \mu_s = g \mu_B / 2.
\]

For all other field directions it is necessary when determining \( H_{c2} \) to take into account also the orbital effect connected with the two-dimensional motion of the electrons inside the layers.

We shall show below (see Sec. 3) that when account is taken of both the paramagnetic and the orbital effects, the inhomogeneous state is realized for field directions close to parallel, and in the LSJI it corresponds to the situation in which the Cooper pairs in the field \( H_{c2} \) are not on the lowest Landau orbit. In the same section, we obtain the angular dependence of \( H_{c2} \) at \( T = 0 \). The temperature dependence of \( H_{c2} \)(1) is investigated in Sec. 4. In Sec. 5 we obtain the dependence of \( H_{c2} (l) \) on the temperature and the degree of the purity of the LSJI.

Section 6 is devoted to dirty LSJI. In them, the inhomogeneous state is not realized, and the transition from the normal state to the superconducting state may turn out to be a first-order transition for field directions close to parallel, and a second-order transition for other directions. In this section we take into account the influence of the spin-orbit scattering on the type of transition and on the value of \( H_{c2} \). In the last section we analyze the experimental data for \( \text{TaS}_2(\text{Py})_{1/2} \) and show that this compound is of the LSJI type, and that an inhomogeneous state can apparently be observed in it.

2. EQUATION FOR THE DETERMINATION OF \( H_{c2} \)

If the condition (1) is satisfied, then in the calculation of \( H_{c2} \) we should take into account the paramagnetic effect and the orbital motion of the electrons inside the...
layers (in the \(x, y\) plane) in a field \(H_z = H \sin \theta\). Assuming the transition from the normal state to the superconducting state to be a second-order transition, we write down the linear equation for the order parameter \\

\[
\Delta(r) \approx \frac{T}{T_c} \approx \int d\tau' X(r, \tau') \Delta(\tau'), \quad r = (x, y),
\]

from which \(H_{c2}\) is determined as the maximum value at which a nontrivial solution for \(\Delta(r)\) exists \([9]\).

To describe the orbital motion of the electrons in a field \(H_z = H \sin \theta\), we use a quasiclassical approximation, since we shall show below that the condition for the applicability of this approximation

\[
e^{-\frac{\alpha H \sin \theta}{\Delta(0)}} \ll n \ll \alpha T
\]

is satisfied at practically all temperatures \(T \gg T_c^0/\varepsilon_j\).

In the lowest approximation in the (nonmagnetic) impurity concentration, we obtain for the kernel of Eq. (2) \([10]\)

\[
X(r, r') = \sum \frac{\pi T}{|\omega|} \delta'(r-r') - S_0(x, y),
\]

where \(\omega = (2n + 1)\pi T\) and the two-dimensional vector potential \(A = (A_x, A_y)\) corresponds to the field \(H_z = H \sin \theta\).

In analogy with the isotropic three-dimensional case \([10]\), it can be shown that \(S_0(x, y)\) is a function of only one operator \((\nabla - 2iA/Hc)^2\), and consequently the solutions of equations (2) and (4) will be the wave functions of the electron motion in the field \(H_z = H \sin \theta\). Choosing \(A\) in the form \(A_x = H_x \sin \theta \) and \(A_y = 0\), we obtain

\[
\Delta(r) = \exp \left\{ \int A(s) \delta s \frac{2eH \sin \theta}{\hbar c} (x + x') (y - y') \right\},
\]

where \(\Delta(r)\) and the two-dimensional vector potential \(A = (A_x, A_y)\) corresponds to the field \(H_z = H \sin \theta\).

Substitution of (5) in (2) and (4) yields

\[
\Delta(r) \approx \frac{T}{T_c} \approx \int d\tau' X(r, \tau') \Delta(\tau'), \quad r = (x, y),
\]

where \(\Delta(r)\) and \(X(r, r')\) are Hermite polynomials and the solutions (5) are degenerate with respect to the parameter \(\alpha_0\).

Choosing \(A\) in the form \(A_x = H_x \sin \theta \) and \(A_y = 0\), we obtain

\[
\Delta(r) = \exp \left\{ \int A(s) \delta s \frac{2eH \sin \theta}{\hbar c} (x + x') (y - y') \right\},
\]

where \(\Delta(r)\) is the Euler constant. Then, after integrating with respect to \(\Omega\), Eq. (7) takes the form

\[
(1 - \frac{T}{T_c}) \ln \frac{\tilde{C}(\Omega, T)}{\tilde{C}(\Omega, 0)} = \int d\tau' L(\tau) \ln \left[ 1 + \left( 1 - \frac{x}{\epsilon} \right)^{\tilde{C}(\Omega, T)} \right]
\]

where \(\tilde{C}(\Omega, T) = C + (2n + 1)\pi T\).

The numerically obtained plots of \(h_{0}(\alpha)\) for \(n = 0\) to 6 are shown in Fig. 1, from which it is seen that at \(\alpha > 1.25\) the solutions for \(\Delta(r)\) corresponding to \(H_c\) are functions of excited Landau orbits. The solid curve in Fig. 1 corresponds to the maximum (at a given \(\alpha\) value) of \(h_{0}\) and it yields the function \(h_{C2}(\alpha)\).

As \(\alpha \to \infty\), we have \(H_c \to 0\) in accordance with the results obtained in \([11]\) and according to the data for \(n = 12\) the value \(\alpha_{\text{max}}\), which determines \(h_{c2}\), increases linearly with increasing \(\alpha\), or, more accurately speaking, \(\alpha_{\text{max}}\) coincides approximately with the integer part of the quantity \(\alpha/1.8\). The function \(h_{C2}(\alpha)\) can now be obtained from (8), (9), and Fig. 1. The inhomogeneous state \((n > 1)\) is realized at an angle \(\theta < \theta_c\)

\[
= \sin^{-1}(H_{c2}\langle H \rangle/H_{c20}),
\]

Thus, whereas in a three-dimensional isotropic superconductor the inhomogeneous state can be realized only at sufficiently large value of \(H_{c20}/H_p\), in LSJII the inhomogeneous state is realized at small values of \(H_{c20}/H_p\), but in a narrow angle interval \(\theta < \theta_{c20}/H_{p}^2\). The dependence of \(H_{c2}\) on \(\theta\) in LSJII is characterized by two features connected with the realization of the inhomogeneous state, namely the sharp increase of \(H_{c2}\) (by an approximate factor 1.5) when \(\theta\) decreases from \(\theta_{c2}\) to zero, and a nonmonotonic dependence of \(H_{c2}\) on \(\theta\) at \(\theta < \theta_{c2}\). The presence of weak oscillations in the plot

3. INHOMOGENEOUS STATE AND ANISOTROPY OF \(H_{c2}\) IN PURE LSJII AT ZERO TEMPERATURE

To change over to the temperature \(T = 0\), we use the asymptotic digamma function. Then, after integrating with respect to \(\Omega\), Eq. (7) takes the form

\[
(1 - \frac{T}{T_c}) \ln \frac{\tilde{C}(\Omega, T)}{\tilde{C}(\Omega, 0)} = \int d\tau' L(\tau) \ln \left[ 1 + \left( 1 - \frac{x}{\epsilon} \right)^{\tilde{C}(\Omega, T)} \right]
\]

where \(\tilde{C}(\Omega, T) = C + (2n + 1)\pi T\).

The inhomogeneous state \((n > 1)\) is realized at an angle \(\theta < \theta_c\)
of $H_{C_2}$ against $\theta$ is due to the fact that when the angle $\theta$ decreases from $\theta_c$ to zero a transition to higher Landau orbits takes place in the system (see Fig. 1). However, the appearance of oscillations in the plot of $H_{C_2}$ against $\theta$ can apparently be observed only in superconductors with not too small values of $H_{C_2}/H_p$.

We can now justify the validity of the quasiclassical approximation used above. Since the paramagnetic effect only decreases $H_{C_2}$, it follows that $H_{C_2}<H_{C_{20}}/\sin \theta$ and at $T \gg T_c'/\epsilon_F$ the condition (3) is indeed satisfied.

4. THE TEMPERATURE DEPENDENCE OF $H_{C_2}(1)$ IN PURE LSJI

For a field parallel to the layers ($\sin \theta = 0$), the dependence of $\Delta(p)$ on the coordinates is of the form $\exp(1, \xi x)$, and Eq. (7) will go over into

$$\ln \frac{T}{T_c} = -\frac{1}{\pi} \int_0^{\infty} \frac{d\Omega}{2\Omega} \left[ \frac{\Omega^2}{24} + \frac{\Omega}{4\pi T_c} \right] - \ln \frac{1}{2} + \frac{\ln(1+2\Psi /4\pi T_c)}{4\pi T_c} \right). \tag{11}$$

Equation (11) determines the dependence of $H$ on $p$ and $H_{C_2}(1)$ corresponds to the maximum of this function with respect to the variable $p$.

Let us find first the asymptotic form of $H_{C_2}(1)$ at $T \ll T_c$. We introduce the dimensionless variables $\omega = \Omega /2\Delta(0), p = \Psi \Psi /2\Delta(0)$, and $\epsilon = \pi T /\Delta(0)$. The dependence of $h$ on $p$ at $T = 0$ was obtained earlier (11), and is determined from the equation

$$h' + |h' - p| = \frac{1}{2}(1 + p'). \tag{12}$$

The maximum of $h$ is reached at $p = 1$, and, as seen from (12), the function $h(p)$ is not analytic in the vicinity of this point. One can therefore expect the dependence of $h_{C_2}(1)$ on $T$ to be likewise nonanalytic as $T - 0$. We can now justify the validity of the quasiclassical approach used above. Since the paramagnetic effect can be neglected when $h_{C_2}(1)$ is calculated. We then obtain from (6) at $n = 0$

$$\ln \frac{1}{T} = \sum_{l} f \left[ \frac{1}{26 + 1} - \frac{\ln \Psi (\omega)}{1 - 2h + \Psi (\omega)} \right], \quad \epsilon = \frac{T}{T_c},$$

$$a = \epsilon (\epsilon^2 + 1 + \Psi), \quad \lambda = \frac{h}{2n T_c}, \quad \lambda = \epsilon \tan \epsilon \left[ 1 - \frac{2}{\sqrt{\pi}} \ln \epsilon \right] \epsilon^\epsilon dz. \tag{17}$$

Like Helfand and Werthamer (10), we plot $h^* = h/(\partial h / \partial \Omega \Omega) = 1$ against $\lambda$, where $(\partial h / \partial \Omega \Omega) = -1$ is the derivative of $h$ with respect to $\Omega$ in the temperature region where the interaction of the layers still remains of the Josephson type, i.e., in the region $\Omega \gg 1 - T / T_c$. Owing to the quasi-two-dimensional character, we have

$$\left( \frac{\partial h}{\partial \Omega \Omega} \right) \frac{1}{\Omega} = 2\pi \left( 1 - \frac{\nu \gamma}{2} + \frac{\Psi (\frac{1}{2} - \lambda)}{\Psi} \right) \epsilon^\epsilon. \tag{18}$$

The results of the numerical calculation are shown in Fig. 3, from which it is seen that $h^*(t)$ in LSJI is lower the smaller $\lambda$. In the isotropic three-dimensional case (10) the situation is reversed, and $h^*(t)$ is smaller the larger $\lambda$. At $T = 0$, with accuracy no worse than 3%, the plot of $H_{C_2}(1)$ against $\lambda$ is approximated by the formula

$$H_{c_2}(1) = H_{c_2}(1 + 1.13)$$

for $\lambda = 6$.

6. TYPE OF PHASE TRANSITION AT $H_{C_2}$ IN DIRTY LSJI

We consider first the case of dirty LSJI ($\Omega \ll \Omega_0$) without allowance for spin-orbit scattering. We can then show that the inhomogeneous state is not realized in the limit of dirty LSJI. Indeed, if we wish to regard the transition from the normal state into the superconduct-
different. An analysis of (22) at \( T = 0 \) \cite{15} shows that at a second-order transition at \( \theta < \theta_c \) and a first-order transition at \( \theta_c < \theta < \theta_{c0} \), where \( \theta_{c0} \) is the field \( \theta_{c0} = \theta_c \) at low temperatures \cite{16}, we can obtain the upper bound of the resistivity \( \rho_l \) across the layers \cite{18}, we can obtain an upper bound of \( h/\tau_l \), using the relation

\[
\frac{1}{D} = \frac{\tau_r}{\tau} = \frac{2\rho}{N(0)},
\]

where \( D \) is the distance between layers, equal to 12 \( \AA \) for \( \text{TaS}_2 \) \( \text{(Pyh12)} \). From \( \rho_l > 10^2 \Omega \cdot \text{cm} \) (at temperatures below 10 \( ^\circ \text{K} \)) we obtain \( h/\tau_l < 0.17 \text{K} \). Comparing this estimate with the value \( T_c = 3.25 \text{K} \), we verify that the condition (1) is satisfied in \( \text{TaS}_2 \) \( \text{(Pyh12)} \) at least for temperatures \( 1 - T/T_c > 3 \times 10^{-4} \) and, consequently, the interaction of the layers is of the Josephson type at practically all the temperatures.

We note that the value of \( h/\tau_l \) can be obtained by measuring the field \( H_{c2}(l) \) for, in accordance with the results of \cite{17} we have

\[
H_{c2}(l) = \frac{c}{4\pi e\hbar} \ln \frac{\lambda_{c}}{d} + \frac{2\pi e}{4\pi e\hbar} \cdot \frac{c}{4\pi e\hbar} \cdot \frac{1}{N(0)},
\]

where \( \lambda_{c} \) is the Riemann function and \( \lambda_{c} \) and \( \lambda_{c} \) are the respective depths of penetration of the field perpendicular and parallel to the layers.

The values of the coherence length \( \xi_{0} \) and the mean free path \( l \) inside the layer can be obtained if we know (in addition to \( N(0) \) and \( T_c \)) the resistance \( \rho_{l||} \) along the layers and the critical field \( H_c(l) \). When \( \xi_{0} \) and \( l \) are determined from the equations

\[
\frac{\xi_{0}}{\tau} = \frac{\hbar}{\pi e^2 \rho_{l||} N(0)} \ln \frac{\lambda_{c}}{d}, \quad \lambda_{c} = \frac{m_{c} c}{4\pi e\hbar}, \quad \lambda_{c} = \frac{c}{4\pi e\hbar} \cdot \frac{1}{N(0)},
\]

and the quantities \( \lambda_{c} \), \( \hbar^{*} \), and \(-d\hbar^{*}/dt\) are defined in Sec. 5.

The measured values of \( \rho_{l||} \) and \( H_{c2}(l) \) for \( \text{TaS}_2 \) \( \text{(Pyh12)} \) are given by Morris and Colemain in \cite{17}. Unfortunately, during the course of the measurements with one and the same crystal, the value of \( \rho_{l||} \), varied in the course of time from \( 10^2 \) \( \Omega \cdot \text{cm} \) at the start of the experiment to \( 6 \times 10^{-7} \Omega \cdot \text{cm} \) at the end. Apparently, the measure-

L. N. Bulavskii

mements of $H_{c2}(\theta)$ at $T=2\,^\circ$K are among the earliest ones, since they yield the largest value of $T_c$. We can therefore take for the yield at the start of the experiments $H_{c2}(\theta=0)=1.4$ kOe and $\rho_0=10^{-5}$ $\Omega$-cm at $T=2\,^\circ$K. In this case we obtain with the aid of (28) $\xi_b=4.4\times10^{-8}$ cm, $\ell=3.5\times10^{-6}$ cm, and $\nu F=1.1\times10^{10}$ cm/sec. Accordingly as in the quasi-two-dimensional case\(^{11}\) we have $n(0)=\frac{m}{2\pi a}$, we get for the effective mass the value $m^*\approx9\times10^{-27}$ g and $k_F\approx10^{12}$ cm\(^{-1}\).

The carrier density $n$ is connected with $k_F$ by the relation $n=\frac{k_F}{2\pi a}$, and for TaS$_2(Py)_{1/2}$ we obtain $n\approx10^{19}$ cm\(^{-3}\). If each Ta atom gives one electron to the conduction band, then the electron density should be of the order of $0.85\times10^{26}$ cm\(^{-3}\). According to the results of Thompson, Gamble, and Koehler\(^{11}\), the concentration of the conducting electrons should be somewhat more than half of this value, since the volume per Ta atom in this material is twice as large as in TaS$_2$ and since, according to the results of Thompson, Gamble, and Koehler, pyridine (Py) adds approximately an additional 0.25 electron to the conduction band, increasing the carrier density by approximately $0.2\times10^{19}$ cm\(^{-3}\). The good agreement of all these estimates for $n$ shows that the values of $\xi_b$ and $\ell$ presented above are not far from the real ones.

Thus, in the study of Morris and Coleman\(^{11}\), the TaS$_2(Py)_{1/2}$ crystal was of intermediate purity at the start of the measurements ($\lambda=1.1$), but by the end of the experiments it must be regarded already as a dirty superconductor ($\lambda=4-6$). Since the spin-orbit scattering is weaker by at least 2 orders of magnitude than the usual scattering, it follows that $\lambda^\circ<0.05$, and the spin-orbit scattering for the investigated samples is immaterial. Therefore, when the field direction approaches parallel ($\theta=0$) and at temperatures $T<\frac{1}{2}T_c$ there should be realized below $H_{c2}$, depending on the degree of purity of the crystal, either an inhomogeneous state or else a first-order transition from the normal state to the superconducting state\(^{10,11}\).

Morris and Coleman\(^{11}\) measured the anisotropy of $H_{c2}$ at 2.84 and 1.4$\,^\circ$K. In the former case the temperature is close to $T_c$ ($\theta=0.88$), the transition from the normal state to the superconducting state is a second-order transition, and in accordance with the results of $^{11}$ the dependence of $H_{c2}$ on $\theta$ is given by the expression

\[
dots H_{c2}(\theta) = H_{c2}(\theta)_{\text{max}} \left[ 1 - \frac{1}{2\pi} \frac{\sin\theta}{\sin\theta_{\text{max}}} \right]^2.
\]

At $T=2.84\,^\circ$K, the value of $H_{c2}$ at $\theta=4^\circ$ is $\approx7$ kOe, and $H_{c2}(\theta=0)=0.5$ kOe (according to the results of Sec. 5, the value of $H_{c2}(\theta=1)/4$ kOe at $T=2\,^\circ$K corresponds to a field $H_{c2}(\theta=0)=0.45$ kOe at $T=2.84\,^\circ$K). From (29) we obtain $H_{c2}(\theta=36.2)$ kOe, and for the angles $\theta=4^\circ$ we have $x\ll1$ and $H_{c2}(\theta)/H_{c2}(\theta=0)\approx1/\sin^2\theta$. At $\theta=3$ and $2^\circ$ we obtain from (29) for the same ratio the values 315 and 560, respectively, whereas the experimental values are 335 and 530.

Thus, the agreement between the theoretical results\(^{11}\) and the experimental data\(^{11}\) for $T=2.84\,^\circ$K is good, and confirms the premise that the interaction of the layers in TaS$_2(Py)_{1/2}$ is of the Josephson type. We note that Morris and Coleman\(^{11}\) have observed a very slow approach of the resistance to the normal value when the field $H$ was increased above $H_{c2}$ at small $\theta$ and at $T=2.84\,^\circ$K. This variation of the resistance with the field is apparently connected with the superconducting fluctuations above $H_{c2}$. The paramagnetic effect facilitates the appearance (owing to the fluctuations) of superconducting regions with characteristic dimensions on the order of $\hbar\nu F/\mu_0 H$ (it is precisely for this reason that an inhomogeneous state is realized in pure LSJI at $T<0.55T_c$ and small $\theta$). Therefore the influence of the fluctuations on the conductivity above $H_{c2}$ (but below $T_c$) should be stronger than above $T_c$.

We consider now the data for $H_{c2}(\theta)$ at $T=1.4\,^\circ$K. At this temperature we have $H_{c2}(\theta=4.9)$ kOe, and for this value of the field we obtain $l=10^{-6}$ cm. Obviously, at such a small mean free path and at a temperature $T=0.43T_c$ the inhomogeneous state is not realized. We apply the theory of dirty LSJI to this case. Then, according to the results of Sec. 6 and the book by St. James et al.\(^{12}\), for angles $\theta>2^\circ$ the transition from the normal state to the superconducting state is a second-order transition. The calculated theoretical values of $H_{c2}(\theta)/H_{c2}(\theta=0)$ for the angles $\theta=6^\circ$ turn out to be in this case lower than the experimental ones (they differ by factors 3.3, 2.5, and 1.5 for the respective angles $2^\circ$, $4^\circ$, and $6^\circ$). This difference can apparently be attributed to the uncertainty in the experimental determination of $H_{c2}(\theta)$ at small $\theta$, owing to the weak and nonmonotonic growth of the resistance with increasing magnetic field above $H_{c2}$. At $\theta=0$ and $T=1.4\,^\circ$K, the transition to the superconducting state from the normal state should be a first-order transition and should occur in a field $-0.9Hp=54$ kOe. Yet experimentally, at $\theta=0$, the resistance amounts to only a small fraction of the resistance in the normal state, even in fields that are 2.5 times larger than $Hp=60$ kOe. This effect is possibly connected with the realization of a metastable state. At $T=0$, the metastable superconducting state with a homogeneous order parameter can be preserved in fields $H<\frac{1}{2}H_p(\theta=0)$ (12). The region of metastable states with an inhomogeneous order parameter is apparently even broader.

In conclusion, the author thanks the participants in V. L. Ginzburg's seminar, and also A. I. Larkin for a useful discussion of the work.

\(^{11}\)We note that according to ($^1$) the critical temperature $T_c$ in LSJI is determined by the density of states $N(0)=\frac{m}{2\pi a}$, where $a$ is the thickness of the conducting layer.

\(^{11}\)L. N. Bulaevskii, Zh. Eksp. Teor. Fiz. 64, 2241 (1973) [Sov. Phys.-JETP 37, 1133 (1973)].


\(^{11}\)L. N. Bulaevskii and Yu. A. Kukharenko, ibid. 60, 1519 (1971) [5, 288 (1966)].


\(^{11}\)L. P. Gor'kov, ibid. 34, 735 (1958); 36, 1918 (1959) [7, 505 (1958); 9, 1364 (1959)].

\(^{11}\)E. Helfand and N. R. Werthamer, Phys. Rev. 147, 268 (1966).
13. D. St. James et al., Sverkhprovodimost' vtorogo roda (Type-II Superconductivity), Pergamon.

Translated by J. G. Adashko

129