LOW TEMPERATURE SUPERFLUID DENSITY OF *d*-WAVE SUPERCONDUCTOR IN AN APPLIED MAGNETIC FIELD

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The temperature and field dependences of the superfluid density ρ_s in the vortex state of a d-wave superconductor are calculated using a microscopic model in the quasiclassical approximation. We show that, at temperatures below $T^* \propto \sqrt{H}$, the linear T dependence of ρ_s crosses over to a T^2 dependence differently from the behavior of the effective penetration depth, $\lambda_{\text{eff}}^{-2}(T)$. We point out that the expected dependences could be probed by a mutual-inductance technique experiment.

Introduction

While the mechanism of superconductivity and unusual nature of the normal state in high-temperature superconductors (HTSC) are not yet understood, there is a consensus that the zero field superconducting state has a d-wave superconducting energy gap, with nodes along the diagonals of the Brillouin zone [1]. The presence of the nodes results in a density of low energy quasiparticle excitations large compared with conventional s-wave superconductors even at the temperatures much smaller than the transition temperature, $T \ll T_c$. Although these excitations are reasonably well described by Landau quasiparticles, their presence brings in a qualitatively new quasiparticle phenomenology not encountered in conventional superconductors. Among many aspects of this new physics, a major role is played by these low energy excitations in the mixed (or vortex) state [2—5]. Since all HTSCs are extreme type-II superconductors, a huge mixed phase extends from the lower critical field, $H_{c1} \sim 10 \div 100\,\text{Gs}$ to the upper critical field, $H_{c2} \sim 100\,\text{T}$. An important property of the vortex state was pointed out in [6], where it was predicted that, in contrast to conventional superconductors, the density of states (DOS) in d-wave systems is dominated by the contribution from excited quasiparticle states rather than the bound states associated with vortex cores. It was shown that, in an applied magnetic field H, the extended quasiparticles DOS $N(\omega = 0, H) \sim \sqrt{H}$ rather than $\sim H$ as in the conventional case. This result was confirmed by specific heat measurements on high quality single crystals [7, 8]. Subsequently, the semiclassical treatment of [6] was incorporated into the Green's functions formalism extended to include the effects of impurity scattering [9]. Accounting for the impurity scattering which violates a simple $\sqrt{H}$ dependence has improved the agreement between the theory and measurements of the electronic specific heat (see Refs. in [4]).

Another interesting result was obtained by the μSR measurements of the temperature and field dependences of the effective in-plane penetration depth, λ_{eff}, in single crystals $YBa_2Cu_3O_{6.95}$ in [10]. At high magnetic fields, a flattening of $\lambda_{\text{eff}}^{-2}$ was observed (defined in these experiments as the width of magnetic field distribution) at low temperatures in contrast to the linear T behavior expected in a clean d-wave superconductor. If one assumes $\lambda_{\text{eff}}^{-2}$ is proportional to the superfluid density, then such a flattening could, in principle, indicate an opening of a secondary gap at the nodes of a d-wave superconducting gap, as was already suggested after the measurements of thermal conductivity [2].

It was argued in [11] that the simple relation $\lambda_{\text{eff}}^{-2} \sim \rho_s$ is not valid for the penetration depth extracted in μSR experiments at finite fields, so that, using the proper definition [12] of λ_{eff} which corresponds to the μSR measured penetration depth, the observed behavior of $\lambda_{\text{eff}}^{-2}$ can be explained by the nonlocal London model for a d-wave superconductor. The behavior of the superfluid stiffness, defined as a second derivative of the free energy of the system with respect to a vector potential, cannot be addressed in this case by the μSR.

Nevertheless, the superfluid density in itself could also be extracted, for example, from the measurements of the low-frequency complex sheet conductance as done in two-coil mutual-inductance technique measurements on thin films [13, 14]. The superfluid density (stiffness) in this case is found from the inductive part $\sigma_2(\omega)$ of the conductivity [15] as

$$\frac{\rho_s(T)}{m} \equiv \frac{c^2}{4\pi e^2 \lambda_{\rm L}^2(T)} = \frac{1}{e^2} \lim_{\omega \to 0} \omega \sigma_2(\omega, T), \tag{1}$$

where e is the electron charge, m is its mass (or, more generally, a mass of the charge carrier), and c is the velocity of light. It is very important to make a distinction between the London penetration depth $\lambda_{\rm L}$ appearing in (1) and $\lambda_{\rm eff}$ deduced from the μSR measurements [12].

Thus, in this work, we investigate the influence of the magnetic field directly on the superfluid stiffness in HTSC using the quasiclassical approximation at $H_{c1} < H \ll H_{c2}$. Our result is that, at low temperatures, there is also a flattening of ρ_s which is not related to an opening of the gap and is just the result of a Doppler shift of the quasiparticle energies in the vortex state.

Although a quantum mechanical treatment is certainly needed for a consistent treatment of the vortex state [5], the semiclassical approximation always provides a good starting point when a new physical property of the system is involved. Furthermore, as argued in [4] for the parameter range relevant to the study of HTSC, this approximation reproduces the energy spectrum of the near-nodal quasiparticles in a vortex state to a high degree of accuracy allowing also to include the effect of impurity scattering into the analysis. In particular, the thermal conductivity was studied using the semiclassical approximation [16—18].

1. Superfluid Stiffness in the Presence of Impurities

Superfluid stiffness is given by [19]:

$$\frac{\rho_s^{\alpha\beta}(T,H)}{m} = \frac{\rho_{s0}^{\alpha\beta}}{m} - \frac{\rho_n^{\alpha\beta}(T,H)}{m}. \tag{2}$$

The normal fluid density, ρ_n in (2), calculated within the "bubble approximation" with dressed fermion propagators (i.e., with self-energy Σ due to the scattering by impurities included) but neglecting the vertex and Fermi liquid corrections, is

$$\frac{\rho_n^{\alpha\beta}}{m} = \int \frac{d^2k}{(2\pi)^2} \int_{-\infty}^{\infty} d\omega \tanh \frac{\omega}{2T} \frac{v_{F\alpha} v_{F\beta}}{4\pi i} \times$$

$$\times {\rm tr}[G_A(\omega,\mathbf{k}) G_A(\omega,\mathbf{k})) - G_R(\omega,\mathbf{k}) G_R(\omega,\mathbf{k})]. \tag{3}$$

As shown in [19], the vertex corrections can be neglected if the impurity scattering potential is isotropic in the $\mathbf{k}$-space. Likewise the Fermi liquid, corrections can be taken into account along the lines of [19]. In (3), $G_{R,A}(\omega,\mathbf{k})$ are the retarded and advanced Green's functions

$$G_{R,A}(\omega,\mathbf{k}) = \frac{(\omega \pm i\Gamma)\hat{I} + \tau_3 \xi(\mathbf{k}) - \tau_1 \Delta(\mathbf{k})}{(\omega \pm i\Gamma)^2 - \xi^2(\mathbf{k}) - \Delta^2(\mathbf{k})} \tag{4}$$

with the dispersion law $\xi(\mathbf{k})$, the d-wave superconducting gap $\Delta(\mathbf{k})$, and the Fermi velocity is $\mathbf{v}_F = \partial \xi(\mathbf{k})/\partial \mathbf{k}|_{\mathbf{k}=\mathbf{k}_F}$. The width $\Gamma = -{\rm Im}\Sigma(\omega)$ in (4) is the scattering rate due to impurities and other sources, e.g., a disordered vortex lattice (see below). Following [19], we assume that, around the nodes, Γ is frequency- and momentum-independent, so that the presence of impurities is modeled by a widening of δ-like quasiparticle peaks by a Lorentzian

$$A(\omega,\mathbf{k}) = \frac{1}{2\pi i}[G_A(\omega - i0,\mathbf{k}) - G_R(\omega + i0,\mathbf{k})] =$$

$$= \frac{\Gamma}{2\pi E}\left[\frac{E + \tau_3\xi - \tau_1\Delta}{(\omega - E)^2 + \Gamma^2} + \frac{E - \tau_3\xi + \tau_1\Delta}{(\omega + E)^2 + \Gamma^2}\right]. \tag{5}$$

with a constant width Γ (see, however, [20] where this assumption was not confirmed for dopants considered as impurity centers). Here $E(\mathbf{k}) = \sqrt{\xi^2(\mathbf{k}) + \Delta^2(\mathbf{k})}$ is the quasiparticle dispersion law.

Employing the nodal approximation for the low temperature regime, $T \ll T_c$ [19], we obtain that

$$\frac{\rho_n^{\alpha\beta}(T)}{m} = \frac{v_F \delta_{\alpha\beta}}{\pi v_\Delta} J, \tag{6}$$

where $\mathbf{v}_\Delta = \partial \Delta(\mathbf{k})/\partial \mathbf{k}|_{\mathbf{k}=\mathbf{k}_F}$ is the gap velocity and

$$J = \frac{1}{\pi} \int_0^{\infty} d\omega \tanh \frac{\omega}{2T} \times$$

$$\times \left[\arctan \frac{\Gamma^2 - \omega^2}{2\omega\Gamma} - \arctan \frac{\Gamma^2 - \omega^2 + p_0^2}{2\omega\Gamma} - \right.$$

$$\left. - \frac{p_0 \Gamma}{(p_0 + \omega)^2 + \Gamma^2} + \frac{p_0 \Gamma}{(p_0 - \omega)^2 + \Gamma^2}\right] \tag{7}$$

with the cutoff energy p_0 (as estimated in [4], $p_0 \sim 1500$ K). Taking the no impurities limit, $\Gamma \to 0$, one can recover the known linear T dependence [19] from (6), (7):

$$\frac{\rho_n(T)}{m} = \frac{2\ln 2}{\pi}\frac{v_F}{v_\Delta}T, \tag{8}$$

where we set $\hbar = k_{\rm B} = 1$. It is essential, however, to take into account the presence of impurities, which modifies the low temperature T dependence of ρ_n from linear to $\sim T^2$ [21, 22]. Indeed, in the limit $\Gamma \gg T$, we get from (6), (7) that

$$\frac{\rho_n(T)}{m} = \frac{v_F}{\pi v_\Delta}\left[\frac{2\Gamma}{\pi}\ln\frac{p_0}{\Gamma} + \frac{\pi}{3}\frac{T^2}{\Gamma} - O\left(\frac{T^4}{\Gamma^3}\right)\right]. \tag{9}$$

For clean YBCO monocrystals, the estimated [23] value of the scattering due to impurities $\Gamma_0 \sim (1\div 2)$ K, but the quadratic temperature dependence of $\rho_s(T)$ is observed in thin films over a wide temperature range [14] implying that the value of Γ_0 could be bigger, so that we use $\Gamma_0 = 10$ K for our estimates. However, as we will see below, the obtained dependences on H are in fact even stronger for smaller values of Γ_0.

2. Doppler Shift

In the semiclassical approach to the vortex state, the presence of a superflow is accounted for by introducing the Doppler shift into the energy $\omega \to \omega - \epsilon(\mathbf{k},\mathbf{r})$ [6, 9], where $\epsilon(\mathbf{k},\mathbf{r}) = \mathbf{v}_s(\mathbf{r})\cdot\mathbf{k}$, and $\mathbf{v}_s(\mathbf{r})$ is the supervelocity field at a position $\mathbf{r}$ created by all vortices. There are several ways to treat the problem depending on the assumptions we made about the vortex lattice.

2.1. Averaging with Lorentzian

Following [17], we assume that the vortex lattice is *disordered*, which is not unreasonable for thin films [14]. Then, on the length scales large compared with the intervortex distance $a_v = \sqrt{\Phi_0/\pi B}$ (here $\Phi_0 = hc/2e$ is the flux quantum and B is the internal field), propagation of quasiparticles is described by the Green's function averaged over the vortex positions $\{\mathbf{R}_i\}$. This averaging can be done using the probability density

$$\mathcal{P}(\eta) = \langle\delta(\eta - \mathbf{v}_s(\mathbf{r})\cdot\mathbf{k})\rangle_{\{\mathbf{R}_i\}}, \tag{10}$$

so that the averaged Green's function

$$\langle G(\omega,\mathbf{k})\rangle = \int_{-\infty}^{\infty} d\eta\mathcal{P}(\eta)G^0(\omega-\eta,\mathbf{k}), \tag{11}$$

where $G^0(\omega,\mathbf{k}) = G(\omega,\mathbf{k})$ from Eq. (4) with $\Gamma = 0$. For typical accessible fields $H_{c1} < H \ll H_{c2}$, the magnetization due to the vortex lattice is small, and the internal field B can be replaced by the applied field, H, directed along the c-axis.

Assuming that the vortex positions are random and uncorrelated, one would get [24] a Gaussian distribution $\mathcal{P}(\eta) = (2\pi E_H^2)^{-1/2}e^{-\eta^2/2E_H^2}$ with

$$E_H^2 = k_\alpha k_\beta\langle v_s^\alpha(\mathbf{r})v_s^\beta(\mathbf{r})\rangle \approx \left(\frac{\hbar v_F}{2a_v}\right)^2 = \frac{\pi}{4}\hbar^2 v_F^2\frac{H}{\Phi_0}, \tag{12}$$

where we used the expression $|\mathbf{v}_s(\mathbf{r})| = \hbar/2mr$ (the Planck's constant $\hbar$ is restored in these expressions) for the supervelocity field created by a single vortex and took $r = a_v$. In what follows, we will use an estimate $E_H[\text{K}] \sim 30\sqrt{H[\text{T}]}$ obtained for the typical values of $v_F \approx (1.5 \div 2.5)\cdot 10^7$cm/s in YBCO. To make an analytical calculation possible, we follow [17] and replace the Gaussian probability distribution by a Lorentzian $\mathcal{P}(\eta) = \pi^{-1}E_H/(\eta^2 + E_H^2)$, so that

$$\langle G_R(\omega,\mathbf{k})\rangle = G_R^0(\omega + iE_H,\mathbf{k}). \tag{13}$$

Thus, within the employed approximation, the effect of quasiparticle scattering from the supervelocity field $\mathbf{v}_s(\mathbf{r})$ created by a completely disordered vortex lattice also results in a widening of the quasiparticle peaks with a constant width E_H. Although this property may not be valid for other forms of the distribution functions, we find it useful for a physical insight to the problem.

We note that the Lorentzian distribution used above has another property that the average of the product of Green's functions is equal to the product of averaged Green's functions, so that the individual Green's function can be averaged over the vortex positions. For any other type of the distribution function, this property is not valid and averaging should be done for the product of the Green's functions itself (see below).

Eq. (13) may be used to calculate the field-induced DOS, $N(\omega,H) = 1/(8\pi^2)\int d^2k\,\text{tr}A(\omega,\mathbf{k})$, so that $N(\omega = 0,H) \propto E_H\ln p_0^2/E_H^2 \sim \sqrt{H}\ln H_0/H$ with $H_0 \sim 2500$ T. This is just DOS from [19] with Γ replaced by E_H. Since $p_0 \gg E_H$, this result agrees with the dependence $N(\omega = 0,H) \sim \sqrt{H}$ obtained in [6]. Furthermore, as we will see below, the extra ln factor is caused by the assumption we made about the completely disordered vortex state. Note also that the first temperature-independent term of Eq. (9) may be identified as a DOS contribution to ρ_n.

Assuming also a Matthiessen-type rule for the impurity and vortex contributions to the lifetime [24],

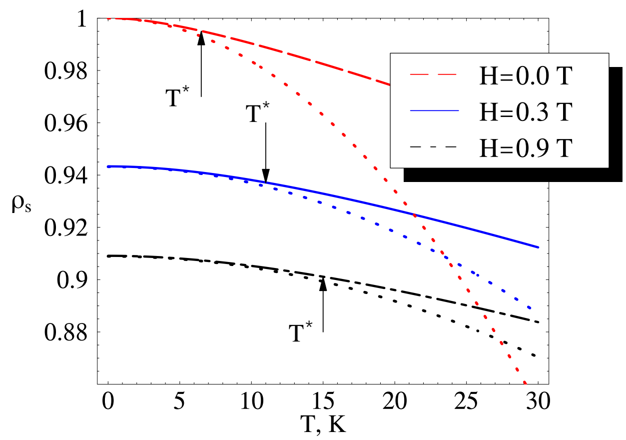

Fig. 1. The temperature dependences of $\rho_s(T,H)/\rho_s(0,0)$ calculated from Eqs. (2), (6), (7) and $\Gamma = \Gamma_0 + E_H$. The dotted lines are calculated using the T^2 asymptotics (9)

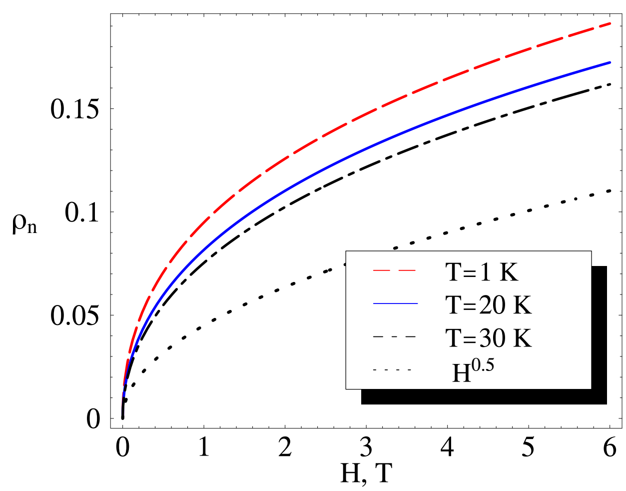

Fig. 2. The dependence $(\rho_n(T,H)-\rho(T,0))/\rho_s(0,0)$ on the applied field calculated from Eqs. (6), (7) and $\Gamma = \Gamma_0 + E_H$. The dotted line is obtained using the first term of Eq. (17), i.e. for $\rho_n = v_F/(\pi v_\Delta)E_H$

we finally obtain that the field dependence of the normal fluid density is described by Eqs. (6), (7) with $\Gamma = \Gamma_0 + E_H$, and its asymptotics for $T \gg \Gamma$ and $T \ll \Gamma$ are given by Eqs. (8) and (9), respectively. These results are presented in Fig. 1 (we used some typical values of the parameters for computation [25], e.g., $v_F/v_\Delta = 15$), where one can clearly see a crossover from the T^2 dependence of $\rho_s(T)$ to the linear T dependence at $T^* \sim \Gamma_0 + E_H$. For a clean system ($\Gamma_0 = 0$) and a single vortex averaging, this crossover of $\rho_s(T)$ was also obtained in [26]. Thus, this behavior of $\rho_s(T)$ differs indeed from the effective penetration depth $\lambda_{\rm eff}^{-2}(T)$ [11] which also crosses over from the linear T dependence, but to the T^3 dependence at (for a clean system) $T^* \sim \sqrt{H}$. Note also that since the nodal approximation was used to arrive at Eqs. (6), (7), one cannot approach the region $\rho_s(T,H)/\rho_s(0,0) \ll 1$, where the dependence of the superconducting gap on temperature should be taken into account.

In Fig. 2, we show the dependence of ρ_s on the applied field H. It appears to be not very different from the $\sim \sqrt{H}$ dependence obtained in [27] for a periodic vortex lattice at $T = 0$. The reason for this coincidence becomes more clear after considering a vortex liquid case.

2.2. *Averaging for Vortex Liquid*

While the averaging with Lorentzian is more simple and allows one to establish a link between two independent problems of scattering from impurities and scattering from random supervelocity field, it is more appropriate to rely on the method of averaging for the vortex liquid used in [4]. In this case, one introduces a Doppler shift $\epsilon_n(\mathbf{r}) = \mathbf{v}_s(\mathbf{r}) \cdot \mathbf{k}_n$ at the node $\mathbf{k}_n$ ($\mathbf{k}_1 = -\mathbf{k}_3$ and $\mathbf{k}_2 = -\mathbf{k}_4$ for a d-wave superconductor, so that $\epsilon_1 = -\epsilon_3$ and $\epsilon_2 = -\epsilon_4$) to approximate the Doppler shift for the entire nodal region. Then if we know how to express a physical quantity F in terms of the Green function $G(\omega,\mathbf{k})$, we can compute its local value $F(\mathbf{r})$ with the local "Doppler shifted" Green's function, $G(\omega,\mathbf{k};\epsilon(\mathbf{k},\mathbf{r})) = G(\omega - \epsilon_n(\mathbf{r}),\mathbf{k})$. We then approximate the field-dependent measured value $F(H)$ by the spatial average which depends on the magnetic field H through ϵ_n as [9] $A^{-1}\int d^2\mathbf{r} F(\epsilon_1(\mathbf{r}),\epsilon_2(\mathbf{r})) = \int_{-\infty}^{\infty} d\epsilon_1 d\epsilon_2 F(\epsilon_1,\epsilon_2)\mathcal{L}(\epsilon_1,\epsilon_2)$, where the first integral is taken over the part of a unit cell of the vortex structure (with the area A) in real space where the Doppler shift is much smaller than the gap maximum and, in the second integral, we use the distribution function

$$\mathcal{L}(\epsilon_1,\epsilon_2) = A^{-1}\int d^2\mathbf{r}\,\delta(\epsilon_1 - \mathbf{v}_s(\mathbf{r})\mathbf{k}_1)\delta(\epsilon_2 - \mathbf{v}_s(\mathbf{r})\mathbf{k}_2), \tag{14}$$

where $\mathbf{k}_1$ and $\mathbf{k}_2$ label two nearest nodes. If $\mathcal{L}(\epsilon_1,\epsilon_2)$ depends on a single variable $\epsilon_1^2 + \epsilon_2^2$, then a simpler distribution function $\mathcal{P}(\epsilon) = \int_{-\infty}^{\infty} d\epsilon_1 \mathcal{L}(\epsilon,\epsilon_1)$ may be used.

The distribution function is aimed at correcting the fact that the single vortex averaging [26] underestimates the number of points where the Doppler shift vanishes.

In a vortex lattice, there exist points where $|\mathbf{v}_s(\mathbf{r})| = 0$: the high symmetry locations such as midpoints between the centers of two neighboring vortices. Consequently, for vortex lattices, $P(0)$ is larger than it is in the single vortex picture. The precise form of the distribution $\mathcal{P}(\epsilon)$ depends on the assumptions made about the vortex lattice structure. Among several candidates for $\mathcal{P}(\epsilon)$ suggested in [4], we choose $\mathcal{P}(\epsilon) = 1/2E_H^2(\epsilon^2 + E_H^2)^{-3/2}$ which is the most convenient for analytical calculations. Using this function, one can obtain the following field-induced DOS: $N(\omega = 0, H) \sim \sqrt{H}$ in the strong field, $\Gamma_0 \ll |\epsilon_n|$ limit and $N(\omega = 0, H) \sim H \ln H_0/H$ in the impurity dominated, $\Gamma_0 \gg |\epsilon_n|$ regime. This shows that this distribution function which allows one to reproduce the Volovik's original result ($\sim \sqrt{H}$) is indeed better for the description of a more ordered vortex glass or even lattice state.

Hence, the local quantity J from Eq. (7) has to be replaced by its local value

$$J(\epsilon_n) = \frac{1}{\pi}\int_0^\infty d\omega \tanh\frac{\omega}{2T}\left[\arctan\frac{\Gamma^2 - (\omega-\epsilon_n)^2 + p_0^2}{2(\omega-\epsilon_n)\Gamma} - \arctan\frac{\Gamma^2 - (\omega-\epsilon_n)^2}{2(\omega-\epsilon_n)\Gamma} + \frac{p_0\Gamma}{(p_0+\omega-\epsilon_n)^2+\Gamma^2} - \frac{p_0\Gamma}{(p_0-\omega+\epsilon_n)^2+\Gamma^2}\right] \qquad (15)$$

and

$$\frac{\rho_n}{m} = \frac{v_F}{\pi v_\Delta}\int_{-\infty}^{\infty} d\epsilon J(\epsilon)\mathcal{P}(\epsilon). \qquad (16)$$

One can obtain that the impurity-dominated limit ($T \ll E_H \ll \Gamma_0$) is described by Eq. (9) with $\Gamma = \Gamma_0$, while, in the field-dominated ($T \ll \Gamma_0 \ll E_H$) regime, we get

$$\frac{\rho_n}{m} = \frac{v_F}{\pi v_\Delta}\left[E_H + \frac{2\Gamma_0}{\pi}\ln\frac{2p_0}{E_H} + \frac{\pi T^2}{3}\frac{\pi}{2E_H}\right]. \qquad (17)$$

Thus, irrespectively of the vortex structure, we assumed the last $\sim T^2$ term of (17) differs only by a factor $\pi/2$ from Eq. (9) taken with $\Gamma = \Gamma_0 + E_H \sim E_H$ for the field-dominated regime, $E_H \gg \Gamma_0$. It also agrees with [26], where, however, there is no second term related to impurities. The first term of (17) is already $\sim \sqrt{H}$, not $\sim \sqrt{H}\ln H_0/H$, exactly as one would expect from the DOS calculation and obtains for a single vortex averaging [26] and the vortex lattice as $T = 0$ [27].

Conclusion

We have shown using the quasiclassical approximation and the simplest form for the impurity scattering that the temperature $T^* \sim \Gamma_0$ of the crossover from the linear to T^2 dependence of the superfluid stiffness $\rho_s(T,H)$ increases in the presence of magnetic field to the temperature $T^* \sim \Gamma_0 + E_H$. The dependence of $\rho_s(T)$ is different from the effective penetration depth, $\lambda_{\text{eff}}^{-2}(T)$, studied in [11]. Concluding, we mention that the preliminary results of the experimental measurements of $\rho_s(T,H)$ studied in [28] confirm [29] that, for $T < 0.4$ K, the field dependence of $\rho_s(H)$ is $\sim \sqrt{H}$ in the low field region.

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НИЗКОТЕМПЕРАТУРНА НАДПРОВІДНА ГУСТИНА У d-ХВИЛЬОВОМУ НАДПРОВІДНИКУ У ЗОВНІШНЬОМУ МАГНІТНОМУ ПОЛІ

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Р е з ю м е

Використовуючи мікроскопічну модель у квазікласичному наближенні, одержано температурну та польову залежності надпровідної густини ρ_s у d-хвильовому надпровіднику. Показано, що при температурах, нижчих за $T^* \sim \sqrt{H}$, лінійна залежність ρ_s переходить у квадратичну, що відрізняється від поведінки ефективної глибини проникнення, $\lambda_{\rm eff}^{-2} \sim T$. Зазначено, що одержані залежності можна перевірити експерементально.

НИЗКОТЕМПЕРАТУРНАЯ СВЕРХПРОВОДЯЩАЯ ПЛОТНОСТЬ В d-ВОЛНОВОМ СВЕРХПРОВОДНИКЕ ВО ВНЕШНЕМ МАГНИТНОМ ПОЛЕ

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Р е з ю м е

Используя микроскопическую модель в квазиклассическом приближении, получены температурная и полевая зависимости сверхпроводящей плотности ρ_s в d-волновом сверхпроводнике. Показано, что при температурах ниже $T^* \sim \sqrt{H}$ линейная зависимость ρ_s переходит в квадратичную, что отличается от поведения эффективной глубины проникновения, $\lambda_{\rm eff}^{-2} \sim T$. Отмечено, что полученные зависимости могут быть проверены на эксперименте.