Low Frequency Magneto-optical Properties of Josephson-Coupled Superconductors

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The low temperature transverse dielectric function is found for Josephson-coupled layered superconductors for polarization of an electric field perpendicular to the layers in the presence of a dc magnetic field oriented parallel and perpendicular to the layers. The plasma edge lays well below the superconducting gap in highly anisotropic systems and it is reduced to zero in strong parallel fields due to the overlap of nonlinear regions of Josephson vortices. In strong perpendicular fields the plasma edge is also reduced due to displacements of pancake vortices from straight lines caused by pinning centers. Corresponding characteristic fields for plasma edge reduction are found.

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It was shown recently that the collective plasma mode may lay well below the superconducting gap in highly anisotropic superconductors for orientation of the electric field $\mathbf{E}$ perpendicular to the layers (along c axis) [1–5]. The plasma frequency for $\mathbf{E} \parallel ab$ is quite large (≈ 1 eV) for high-T_c superconductors. In contrast, the frequency of the plasma mode for $\mathbf{E} \parallel \mathbf{c}$ with a momentum along the c axis, $\omega_c = c/\lambda_c\sqrt{\epsilon_0} = c/\lambda_{ab}\gamma\sqrt{\epsilon_0}$, may be very low because it is inversely proportional to the anisotropy parameter γ. Here λ_{ab} and λ_c are the London penetration lengths for currents along the a-b plane and c axis, respectively; ϵ_0 is the high frequency dielectric constant for the electric field along the c axis. In Bi-2:2:1:2 with γ in the range 300–1000 the plasma frequency ω_c should lay in the range 10–30 cm^{-1}, well below the superconducting gap $\Delta \approx 300$ cm^{-1}.

The measurements of optical reflectivity performed by Tamasaku, Nakamura, and Uchida [6] for $\mathbf{E} \parallel \mathbf{c}$ in $La_{1-x}Sr_xCuO_4$ show almost complete reflection at low temperatures in the frequency range below 20–50 cm^{-1}. This result was explained by Tachiki, Koyama, and Takahashi [2] as originating from the dependence of the transverse dielectric function for $\mathbf{E} \parallel \mathbf{c}$ on the low lying plasma frequency ω_c:

$$\epsilon_t(\omega) = \epsilon_0\left[1 - \frac{\omega_c^2}{\omega(\omega + i\omega_r)}\right], \qquad (1)$$

where ω_r is the relaxation rate.

In the following we show that in Josephson-coupled superconductors (Bi- and Tl-based systems and organic layered superconductors) the plasma edge ω_c in Eq. (1) at low temperatures is determined by interlayer Josephson coupling (ω_c^2 is proportional to the Josephson interlayer critical current density $J_c^{(c)}$) and thus may be strongly reduced by the dc magnetic field applied parallel [7,8] or perpendicular [9] to the layers. We find that $H_0 = \Phi_0/\gamma s^2$ is the characteristic field for this effect in parallel fields; s is the interlayer spacing. For the perpendicular field the characteristic magnetic field B_D which affects both $J_c^{(c)}$ and ω_c strongly depends on the pinning, and we express the decoupling field B_D in terms of the critical current density along the layers, $J_{ab}^{(c)}$. We find that for Bi-2:2:1:2 both fields H_0 and B_D lay in the interval $\approx 1 - 3$ T, and this system is the best candidate for observing magneto-optical effects.

We note that the decrease of the plasma edge in the longitudinal dielectric function $\epsilon_l(\omega)$ for parallel magnetic field was described in [4]. The behavior of $\epsilon_l(\omega)$ in the magnetic field may be studied by measurements of reflectivity on the superconductor in a waveguide, but not in optical measurements where the transverse dielectric function is important. Thus in the following we extend the results obtained in [4] to the optical properties which are sensitive to the transverse dielectric function.

We choose the coordinate system in such a way that the c axis of the crystal coincides with the z axis, light propagates along the x axis and alternating field, and E_z is along the z axis. The applied dc magnetic field is along the y axis for orientation parallel to the layers and along the z axis for perpendicular orientation.

The Maxwell equations read

$$\frac{\partial E_z}{\partial x} = \frac{1}{c}\frac{\partial B_y}{\partial t},$$
$$\frac{\partial B_y}{\partial x} = \frac{\epsilon_0}{c}\frac{\partial E_z}{\partial t} + \frac{4\pi}{c}\left[J_0\sum_n f_{n,n+1}(z) \times \sin\varphi_{n,n+1}(\mathbf{r},t) + \frac{E_z}{\rho_c}\right]. \qquad (2)$$

The first term in the square brackets describes the Josephson interlayer current, $J_0 = c\Phi_0/8\pi^2\lambda_c^2 s$ is the Josephson critical current density, $f_{n,n+1}(z) = 1$ if $ns < z < (n+1)s$ and 0 otherwise, and coordinates of layers are $z = ns$ and $\mathbf{r} = x, y$. We denote by $\varphi_{n,n+1}$ the gauge-invariant phase difference between layers n and $n+1$; it includes the part induced by the applied dc field B and another part induced by the field $E_z(t)$. The second term in the square brackets describes the current of quasiparticles in the context of the resistively shunted junction model [10], and $\omega_r \propto \rho_c^{-1}$. Note that concentration of quasiparticles tends to zero as $T \to 0$ due to the superconducting gap. Thus, $\rho_c^{-1} \to 0$ at $T \to 0$, and reflectivity at frequencies below the plasma edge strongly increase below T_c, especially at low temperatures (see also [5]).

The phase difference obeys the Josephson relation:

$$\frac{\partial \varphi_{n,n+1}(\mathbf{r},t)}{\partial t} = \frac{2es}{\hbar}\langle E_z(\mathbf{r},t)\rangle_{n,n+1}, \qquad (3)$$

where $\langle E_z(\mathbf{r})\rangle_{n,n+1}$ is the average electric field between layers n and $n+1$ at coordinate $\mathbf{r}$. In the following we assume that the electric field changes slowly in z direction and we replace $\langle E_z(\mathbf{r})\rangle_{n,n+1}$ by $E_z(\mathbf{r},z,t)$.

We consider the linear response of a superconductor to the electromagnetic field $E_z(t)$. We express $\varphi_{n,n+1}(\mathbf{r},t)$ as the sum of a time-independent part $\varphi_{n,n+1}^{(0)}(\mathbf{r})$, which is due to supercurrent induced by the dc field B and small ac part $\varphi_{n,n+1}^{(1)}(\mathbf{r},t)$, which obeys Eq. (3). Then in Eq. (2) we expand in $\varphi_{n,n+1}^{(1)}(\mathbf{r},t)$ and use the Fourier transformation for the time variable t. Finally, with the help of Eqs. (2) and (3) we get the equation for the electric field:

$$\left[\frac{\partial^2}{\partial x^2} + \frac{\omega^2}{c^2}\epsilon_t(\mathbf{r},z,\omega)\right]E_z(\mathbf{r},z,\omega) = 0, \qquad (4)$$

where the coordinate-dependent dielectric function is

$$\epsilon_t(\mathbf{r},z,\omega) = \epsilon_0 - \frac{\omega_c^2\epsilon_0}{\omega^2}\sum_n f_{n,n+1}(z)\cos\varphi_{n,n+1}^{(0)}(\mathbf{r}) + \frac{i}{\omega\rho_c}. \qquad (5)$$

Equations (4) and (5) describe propagation of an electromagnetic wave with frequency ω in a media with a coordinate-dependent dielectric function. This function is determined by the phase differences between layers in the presence of the applied dc magnetic field. In the Meissner state we get $\varphi_{n,n+1}^{(0)} = 0$ and $\epsilon_t(\omega)$ is given by Eq. (1). In a mixed state $\varphi_{n,n+1}^{(0)}(\mathbf{r})$ becomes nonzero and the dielectric function depends on the vortex lattice.

The phase difference $\varphi_{n,n+1}^{(0)}(\mathbf{r})$ is determined [8,9] by

$$\nabla^2\varphi_{n,n+1} + \lambda_J^{-2}[\sin\varphi_{n+1,n+2} + \sin\varphi_{n-1,n} - (2 + s^2/\lambda_{ab}^2)\sin\varphi_{n,n+1}] = 0, \qquad (6)$$

where $\lambda_J = \gamma s$ is the Josephson length and $\nabla = \partial/\partial\mathbf{r}$. For the parallel field, solutions of this equation that are periodic in x and n should be found, with the area of the unit cell $a_0^2 = \Phi_0/B$. For the perpendicular field, the positions of pancake vortices $\mathbf{r}_{n\nu}$ (ν labels vortices in layer n) determine the positions of singularities for the phase difference $\varphi_{n,n+1}^{(0)}(\mathbf{r})$ according to the condition

$$(\nabla_x\nabla_y - \nabla_y\nabla_x)\varphi_{n,n+1}(\mathbf{r}) = \sum_\nu[\delta(\mathbf{r}-\mathbf{r}_{n\nu}) - \delta(\mathbf{r}-\mathbf{r}_{n+1,\nu})]. \qquad (7)$$

In this case the positions of pancake vortices should be obtained first, then $\varphi_{n,n+1}^{(0)}(\mathbf{r})$ may be found by solving Eq. (6) with boundary condition (7).

We consider first the effect of the dc field $\mathbf{B}\perp\mathbf{c}$. The structure of the lattice formed by Josephson vortices in weak fields $B \ll H_0$ is very similar to that of the Abrikosov vortices in anisotropic superconductors [8,11]. Centers of vortices form a triangular lattice. They form rows along the x axis where the distance between centers is a. ls is the distance between these rows along the z axis; l is an integer part of $a/s2\sqrt{3}\gamma$. The parameters a, l are determined by the condition that magnetic flux is Φ_0 per vortex. The difference with the standard Abrikosov lattice is that the center of each Josephson vortex is positioned between the layers (say, n and $n+1$), and thus the normal core is absent. Instead, there is a region of length λ_J along the x axis around the center of a vortex inbetween the layers n and $n+1$ where interlayer current density is of the order J_0 and the phase difference $\varphi_{n,n+1}^{(0)}(x)$ differs significantly from $2\pi m$ where m is an integer. In this region the nonlinear character of the Josephson current is important, and here $\cos\varphi_{n,n+1}^{(0)}(x)$ deviates significantly from unity. Far away from the vortex positioned at $x=0, n=0$ we can use the linear approximation for Eq. (6). Here we get

$$\varphi_{n,n+1}^{(0)}(x) \approx \frac{s}{l}\frac{(x/a)}{(x/a)^2 + (sn/l)^2}. \qquad (8)$$

For magnetic fields $B > 0.001$ T and frequencies ω in the range 0.01–50 cm^{-1}, the wavelength of the light is large compared with the intervortex distances a and l. Thus we can average $\epsilon_t(x,z,\omega)$ over coordinates, replacing $\cos\varphi_{n,n+1}^{(0)}(x)$ by its average value over the unit cell (with area $9al/8$ in the limit of large γ). The main contribution comes from the region where the phase difference is given by Eq. (8). Finally we get Eq. (1) with ω_c replaced by

$$\omega_c(B) = \omega_c\left[1 - \frac{\pi}{8}\frac{B}{H_0}\ln\frac{H_0}{B}\right], \qquad B \ll H_0. \qquad (9)$$

In the limit of strong fields, $B \gg H_0$, the intervortex distance along the x axis, $a = \Phi_0/sB$, becomes much smaller than λ_J, and nonlinear regions overlap. The vortex lattice in this limit is simple: vortices fill all interlayer spacings; the intervortex distance along the x axis is $a = \Phi_0/sB$. Vortices still form a triangular lattice, because, in the neighboring interlayer spacings, the positions of the vortices are shifted by $a/2$ (see [8]). The phase difference was found in [8] using perturbation theory with respect to the small parameter a^2/λ_J^2:

$$\varphi_{n,n+1}(x) = \frac{2\pi x}{a} + \pi n - \frac{a^2}{\pi^2\lambda_J^2}(-1)^n\sin\frac{2\pi x}{a}. \qquad (10)$$

Averaging $\cos\varphi_{n,n+1}$, we get Eq. (1) with ω_c replaced by

$$\omega_c(B) = \omega_c\frac{\sqrt{2}H_0}{\pi B}, \qquad B \gg H_0. \qquad (11)$$

We see that the plasma edge decreases linearly with a field at small B and vanishes $\propto B^{-1}$ as the field increases above H_0. For Bi-2:2:1:2 with $s = 15.6$ Å we get $H_0 \approx 1-3$ T and for $La_{1-x}Sr_xCuO_4$ with $\gamma \approx 20$ we estimate $H_0 \approx 50$ T. A preliminary result of the reflectivity experiment using a single crystal of $La_{1.9}Sr_{0.1}CuO_4$ clearly shows that the plasma edge shifts to low frequency when an external field is applied parallel to layers as we predicted [12].

In a perpendicular field, pancake vortices are arranged along straight lines $\mathbf{r}_\nu^{(0)}$ and these lines form a triangular lattice if pinning is absent. For this lattice $\varphi_{n,n+1}^{(0)}(\mathbf{r}) = 0$ as follows from Eq. (7). Randomly positioned pinning centers cause displacements of pancake vortices from straight lines [we consider here only the temperatures well below the irreversibility line $T_{\mathrm{irr}}(B)$ and ignore fluctuations]. In this case we replace $\cos\varphi_{n,n+1}^{(0)}$ in Eq. (5) by its value averaged over disorder, $\langle\cos\varphi_{n,n+1}^{(0)}\rangle$. The following formalism is close to that presented in [9] and [13]. To define positions of pancakes we assume that pinning is weak and that displacements of pancake vortices caused by pinning, $\mathbf{u}_{n\nu} = \mathbf{r}_{n\nu} - \mathbf{r}_\nu^{(0)}$, are small. We find $\mathbf{u}_{n\nu}$ taking into account the balance of pinning and elastic restoring forces.

The pinning force acting on a pancake at position $\mathbf{r}_{n\nu}$ is

$$\mathbf{F}_p(n, \mathbf{r}_{n\nu}) = -\int d\mathbf{r}\, U_n(\mathbf{r}) \frac{\partial p(\mathbf{r} - \mathbf{r}_{n\nu})}{\partial \mathbf{r}}. \quad (12)$$

Here $U_n(\mathbf{r})$ is the random pinning potential in the layer n, and $p(\mathbf{r})$ is the form factor of a pancake vortex. We suppose $U_n(\mathbf{r})$ is the Gaussian, with correlation function $\langle U_n(\mathbf{r})U_m(\mathbf{r}')\rangle = \gamma_U \delta(\mathbf{r} - \mathbf{r}')\delta_{nm}$, where the parameter γ_U characterizes disorder. We use $p(\mathbf{r}) = \xi_{ab}^2/(r^2 + \xi_{ab}^2)$, where ξ_{ab} is the superconducting correlation length in the a-b plane. The pinning force determines transverse displacements which are important in the following:

$$\mathbf{u}_t(q, \mathbf{k}) = \mathbf{u}(q, \mathbf{k}) - \frac{\mathbf{k}\mathbf{u}(q, \mathbf{k})}{k^2}\mathbf{k} = \frac{\mathbf{F}_{p;t}(q, \mathbf{k})}{c_{66}k^2 + c_{44}(q, \mathbf{k})Q^2},$$

$$\mathbf{F}_p(q, \mathbf{k}) = \sum_{n,\nu} F_p[n, \mathbf{r}_\nu^{(0)}] \exp[iqn + i\mathbf{k}\mathbf{r}_\nu^{(0)}]. \quad (13)$$

Here we denote by $\mathbf{u}(q, \mathbf{k})$ the Fourier components of displacements:

$$\mathbf{u}(q, \mathbf{k}) = a_0^2 \sum_{n,\nu} \mathbf{u}_{n\nu} \exp(i\mathbf{k}\mathbf{r}_\nu + iqn). \quad (14)$$

The momentum $\mathbf{k}$ is in the a-b plane and we assume a circular Brillouin zone of radius $\sqrt{4\pi}/a_0$ for summation over $\mathbf{k}$. q is the momentum for the discrete variable n; $0 \le q \le 2\pi$ and $Q^2 = 2(1 - \cos q)/s^2$. The flux lattice shear modulus (per vortex unit length) is $c_{66} = \Phi_0^2/(8\pi\lambda_{ab})^2$. The tilt modulus is

$$c_{44} = \frac{B\Phi_0}{4\pi(1 + \lambda_c^2 k^2 + \lambda_{ab}^2 Q^2)} + \frac{\Phi_0^2}{32\pi^2\lambda_{ab}^4 Q^2} \ln\left(1 + \frac{a_0^2 Q^2}{4\pi}\right). \quad (15)$$

We neglect in c_{44} the contribution from the Josephson coupling, assuming $B \ll \Phi_0\gamma^2/\lambda_{ab}^2$.

Knowing displacements of pancakes, we can find the phase difference by solving Eqs. (6) and (7). For that we use a linear approximation replacing $\sin\varphi_{n,n+1}$ by $\varphi_{n,n+1}$ in Eq. (6). In this approach we increase three-dimensional effects, which tend to diminish the phase difference. Thus we underestimate the effect of magnetic field on $J_c^{(c)}$ and ω_c. Expanding over displacements $\mathbf{u}_{n\nu}$ we find

$$\varphi_{n,n+1}^{(0)}(\mathbf{r}) = a_0^{-2} \sum_{\mathbf{k},q} e^{-iqn}(1 - e^{iq})\mathbf{D}(\mathbf{k}, q)\mathbf{u}(q, \mathbf{k}), \quad (16)$$

where $\mathbf{D}(\mathbf{k}, q)$ is given by

$$\mathbf{D}(\mathbf{k}, q) = \frac{2\pi i}{a_0^2} \frac{k_y\hat{\mathbf{x}} - k_x\hat{\mathbf{y}}}{k^2 + 2\lambda_J^{-2}(1 - \cos q) + \lambda_c^{-2}}. \quad (17)$$

Here $\hat{\mathbf{x}}$ ($\hat{\mathbf{y}}$) are the unit vectors along the x (y) axis. According to Eqs. (16) and (17) the phase difference depends only on the transverse component of distortions.

Using Eqs. (12)–(17) we obtain

$$\varphi_{n,n+1}^{(0)} = \frac{i}{a_0^2} \sum_{\mathbf{k},q,\nu} \int d\mathbf{K}(1 - e^{iq})p(\mathbf{K}) \frac{U(q, \mathbf{k})}{\epsilon_t(q, \mathbf{k})} \times \left(\mathbf{K} - \frac{\mathbf{K}\cdot\mathbf{k}}{k^2}\mathbf{k}\right)\mathbf{D}(q, \mathbf{k}) \times \exp[i(\mathbf{K} - \mathbf{k})\cdot\mathbf{r}_\nu^{(0)} + iqn], \quad (18)$$

where we introduced the Fourier transforms of the form factor $p(\mathbf{K})$ and pinning potential $U(q, \mathbf{K})$. Averaging over disorder we get $\langle\cos\varphi_{n,n+1}\rangle = \exp(-S)$, where $S = \langle\varphi_{n,n+1}^2\rangle/2$ is given by

$$S = \frac{\gamma_U}{a_0^6} \sum_{\mathbf{G},\mathbf{k},q} (1 - \cos q)p^2(\mathbf{G} + \mathbf{k}) \times \left|\left(\mathbf{G} - \frac{\mathbf{G}\cdot\mathbf{k}}{k^2}\mathbf{k}\right)\mathbf{D}(q, \mathbf{k})\right|^2. \quad (19)$$

Here $\mathbf{G}$ are reciprocal vectors of the vortex lattice. Note that for a perpendicular field ω_c may be expressed in terms of the effective interlayer critical current density $J_c^{(c)}$ because both depend on $\langle\cos\varphi n, n + 1\rangle$ [9]:

$$\omega_c^2(B) = J_0 \frac{8\pi^2 cs}{\epsilon_0\Phi_0} \langle\cos\varphi_{n,n+1}\rangle = J_c^{(c)}(B) \frac{8\pi^2 cs}{\epsilon_0\Phi_0}. \quad (20)$$

Our goal is to estimate the field B_D above which the Josephson coupling is strongly suppressed due to pancake vortices and, correspondingly, the plasma edge becomes reduced. We do not account for the renormalization of γ in the tilt modulus c_{44}, Eq. (15), and in $\lambda_J = \gamma s$ and $\lambda_c = \gamma\lambda_{ab}$, Eq. (17), which would lead to a stronger effect of pancake vortices on the Josephson coupling and plasma edge for fields above B_D (see [9]). To calculate $S(B)$ we approximate $p(\mathbf{G} + \mathbf{k}) \approx p(\mathbf{G})$ and replace summation over $\mathbf{G}$ by integration:

$$\sum_{\mathbf{G}} p^2(G)G^2 \approx \frac{2}{3}\pi a_0^2. \quad (21)$$

The main contribution to S comes from small q and k because $(a_0/\lambda_c)^2 \ll 1$. Integrating over q and k we get $S(B) \approx (B/B_D)^{3/2}$, where

$$B_D \approx \frac{\Phi_0^{11/3}}{(4\pi)^3(2\pi\gamma\gamma_U)^{2/3}s^2\lambda_{ab}^4}. \qquad (22)$$

To estimate B_D we should know the pinning parameter γ_U. That can be obtained from the value of critical current density $J_{ab}^{(c)}$ along layers in the framework of the collective pinning theory (see [13,14]). van der Beek *et al.* [14] found that in Bi-2:2:1:2 the regime of weak collective pinning of single vortex takes place in the range of fields below $\approx$5 T. In this regime the correlated volume consists of one pancake of dimensions $a_0^2 s$; such a regime takes place in the fields below $\approx 10(J_{ab}^{(c)}/\tilde{J}_0)H_{c2}$, where $\tilde{J}_0 = 4c\Phi_0/3\sqrt{3}\xi_{ab}(4\pi\lambda_{ab})^2$ is the depairing current along the a-b plane and $H_{c2} = \Phi_0/2\pi\xi_{ab}^2$. The critical current density in the single pancake pinning regime is

$$J_{ab}^{(c)} = \tilde{J}_0 \frac{\gamma_U^{1/2}\xi_{ab}}{4c_{66}} = \frac{4c\gamma_U^{1/2}}{3\sqrt{3}\Phi_0}.$$

For $J_{ab}^{(c)}$ in the range 5×10^5–5×10^6 A/cm^2 [14], we obtain γ_U in the interval 2×10^{-4}–2×10^{-2} erg^2/cm^4. Taking the parameter $\lambda_{ab} = 1700$ Å, we estimate the decoupling field B_D at $T \ll T_{\rm irr}$ to be in the interval 0.4–4 T. Similar scale for suppression of $J_c^{(c)}$ by a perpendicular field at low temperatures was found in Bi-2:2:1:2 by Cho *et al.* [15] from the *I-V* curve measurements.

As the temperature increases, the effect of pinning decreases due to the thermal depinning. However, thermal fluctuations of pancake vortices reduce the Josephson coupling [9] and plasma edge $\omega_c(B)$. According to the data obtained by Cho *et al.* [15], the critical current along the c axis, $J_c^{(c)}$, vanishes along the decoupling line which lays slightly above the irreversibility line in the plane (B,T). On this line ω_c should vanish also.

In summary, we have shown that low frequency optical properties of Josephson-coupled superconductors for polarization of the electric field along the c axis at low temperatures are strongly affected by the dc magnetic fields both parallel and perpendicular to the layers. Such strong magnetic fields shift the plasma edge to zero frequency, the Josephson character of the interlayer current being responsible for this effect. Thus reflectivity measurements in magnetic fields provide information on the effective interlayer critical current $J_c^{(c)}$ averaged over the sample, Eq. (20). In the case of the perpendicular field they also provide information on pancake vortices which strongly affect $J_c^{(c)}$ and ω_c. Note that $J_c^{(c)}$ determined by such measurements may be more accurate than that determined from *I-V* curve (dissipation) data. In the latter case an arbitrary voltage criterion is used to define $J_c^{(c)}$, while in the optical method this problem does not exist because $J_c^{(c)}$ is extracted from the real part of the dielectric function.

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[1] H. A. Fertig and S. Das Sarma, Phys. Rev. B **44**, 4480 (1991); T. Mishonov, Phys. Rev. B **44**, 12 033 (1991).
[2] M. Tachiki, T. Koyama, and S. Takahashi, Phys. Rev. B **50**, 7065 (1994).
[3] S. N. Artemenko and A. G. Kobel'kov, Pis'ma Zh. Eksp. Teor. Fiz. **58**, 435 (1993) [JETP Lett. **58**, 445 (1993)].
[4] L. N. Bulaevskii, M. Zamora, D. Baeriswyl, H. Beck, and J. R. Clem, Phys. Rev. B **50**, 12 831 (1994).
[5] S. V. Pokrovskii and V. L. Pokrovskii (to be published).
[6] K. Tamasaku, Y. Nakamura, and S. Uchida, Phys. Rev. Lett. **69**, 1455 (1992).
[7] K. B. Efetov, Zh. Eksp. Teor. Fiz. **76**, 1781 (1979) [Sov. Phys. JETP **49**, 905 (1979)].
[8] L. N. Bulaevskii and J. R. Clem, Phys. Rev. B **44**, 10 234 (1991); A. F. Volkov, Physica (Amsterdam) **192C**, 306 (1992).
[9] L. I. Glazman and A. E. Koshelev, Phys. Rev. B **43**, 2835 (1991); Physica (Amsterdam) **173C**, 180 (1991). L. L. Daemen, L. N. Bulaevskii, M. P. Maley, and J. Y. Coulter, Phys. Rev. Lett. **70**, 1167 (1993); Phys. Rev. B **47**, 11 291 (1993).
[10] B. G. Orr, J. R. Clem, H. M. Jaeger, and A. M. Goldman, Phys. Rev. B **34**, 3491 (1986).
[11] L. J. Campbell, M. M. Doria, and V. G. Kogan, Phys. Rev. B **38**, 2439 (1988).
[12] S. Kimura (unpublished).
[13] G. Blatter, M. V. Feigel'man, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur (to be published).
[14] C. J. van der Beek, P. Kes, M. Maley, M. J. V. Menken, and A. A. Menovsky, Physica (Amsterdam) **195C**, 307 (1992).
[15] J. H. Cho, M. P. Maley, S. Fleshler, A. Lacerda, and L. N. Bulaevskii, Phys. Rev. B **50**, 6493 (1994).