Current Induced Decomposition of Abrikosov Vortices in *p-n* Layered Superconductors and Heterostructures

A. L. Rakhmanov,[1,2] Sergey Savel'ev,[1] and F. V. Kusmartsev[1]

[1]*Department of Physics, Loughborough University, Loughborough LE11 3TU, United Kingdom*
[2]*Institute for Theoretical and Applied Electrodynamics Russian Academy of Sciences, 125412 Moscow, Russia*
(Received 3 June 2008; published 4 November 2008)

We describe the decomposition of Abrikosov vortices into decoupled pancake vortices in superconductors having both electron and hole charge carriers. We estimate the critical current of such a decomposition, at which a superconducting-normal state transition occurs, and find that it is very sensitive to the magnetic field and temperature. The effect can be observed in recently synthesized self-doped high-T_c layered superconductors with electrons and holes coexisting in different Cu-O planes and in artificial p-n superconductor heterostructures. The sensitivity of the critical current to a magnetic field may be used for sensors and detectors of a magnetic field, which can be built up from the superconductor heterostructures.

DOI: 10.1103/PhysRevLett.101.197002 PACS numbers: 74.25.Qt, 74.72.Hs

Recently, great progress has been made in atomic-layer synthesis of precise multilayers and superlattices with high-T_c superconductors having good characteristics and uniformity [1]. This opens up prospects for development of a variety of novel devices, combining different functionality. A possible realization of such a heterostructure is a multilayer consisting of alternating high-T_c and low-T_c superconductors (e.g., a stack of $YBa_2Cu_3O_7/Nb$ junctions [2]). In these multilayers, the charge carriers in the neighboring layers have different signs (holes (p) in high-T_c and electrons (n) in low-T_c layers) forming p-n superconductor heterostructures. Moreover, the synthesis has been recently reported of self-doped layered superconductors with the composition $Ba_2Ca_3Cu_4O_8F_2$ having p and n types of charge carriers in different Cu-O planes [3]. In these compounds, two neighboring Cu-O planes are electron doped while the next two planes are hole doped.

It can be expected that due to alternating p and n conductivity in the layers, the physical properties of the combined electron-hole superconductors should be very unusual for both the self-doped compounds and artificial heterostructures. These may include an anomalous intrinsic Josephson effect and a very peculiar Hall effect. It is well established that in the layered superconductors, an Abrikosov vortex is viewed as a stack of pancakes – two dimensional vortices, each of them located in a single superconducting layer. We expect that the properties of the single vortex and the behavior of the flux line lattice (FLL) should be very specific in the p-n superconductors. Here, an Abrikosov vortex experiences a sign-alternating component of the force acting on pancake vortices in p or n superconducting layers when the electric current is applied. This alternating force is a consequence of the dependence of the Hall electromotive force (emf) on the sign of the charge carriers. In highly anisotropic superconductors, even a weak alternating force arising at low current density can destroy the pancake vortex stack due to weak vortex rigidity. We should also note that the anomalous Hall effect with a sign change of Hall emf with magnetic field is observed in typical high-T_c materials [4]. Recent experiments [5] reveal that such an anomaly can be attributed to the presence of both p and n types of charge carriers rather than to the properties of the FLL in superconducting cuprates. Thus, the study of superconducting materials with coexisting p and n types of conductivity is of interest for both fundamental physics and applications.

In this Letter, we analyze consequences that follow from the Hall effect in p-n layered compounds [3]. The sign of the Hall emf depends on the sign of the charge carriers. A component of the force, F_H, acting on each of the pancake vortices in the direction collinear to the current flow also depends on the sign of the charge carriers. As a result, due to the alternating p and n types of conductivity in different layers in the p-n superconductor, the value F_H changes its sign with the period d_c. Here, d_c is either the crystal lattice constant in the **c** direction (for layered superconductors) or the interlayer distance (in the case of artificial heterostructures). The force F_H is small compared to the Lorentz force F_L, which is independent of the sign of the charge carriers and directed perpendicular to the current. However, the tilt modulus, C_{44}, of the FLL is small in layered superconductors and the sign-alternating force F_H can "slice" the stack of pancakes at current densities j much smaller than the depairing current density j_d. Such an effect gives rise to a decoupling of the pancake vortices in the vortex stacks and produces a melting of the pancake vortex lattice with a mechanism different from the vortex loop unbinding discussed, e.g., in Ref. [6]. In the present case, each vortex of the FLL is decomposed after melting into decoupled n- and p- pancakes forming an n-, p- pancake "plasma." The existence of such a "plasma" may give rise to either positive or negative Hall constant. The resulting sign of the Hall constant depends on the mobility and the concentration of the p and n charge carriers. Here, we consider the dissociation of pancake vortices into an n-p pancake "plasma" and analyze the effect of F_H on the value of the melting temperature, T_m, of the FLL. We found that the alternating Hall force significantly affects the FLL melting temperature in the layered *p-n* superconductors even at relatively low values of the applied current density. Therefore, we believe that this phenomenon can be observed experimentally and used in applications. We believe also that in the case of the self-doped cuprates, it can confirm independently the picture of alternating n and p Cu-O planes suggested for the self-doped superconductors in Ref. [3] on the basis of ARPES data.

Model.—We describe the FLL motion by the Nozières-Vinen model [7], which is appropriate for estimates (except the case of d-wave superconductors in the superclean limit [8]). According to this model, the dynamic equation for FLL moving with velocity v_L can be written as

$$\frac{n_s e}{c}(\mathbf{v}_s - \mathbf{v}_L) \times \mathbf{B_L} = -\eta \mathbf{v}_L, \quad (1)$$

where n_s, e, and v_s are the density, charge, and velocity of the superconducting charge carriers, $\mathbf{B_L}$ is the magnetic induction produced by the lattice of the pancake vortices, $\eta = \sigma_n H_{c2} B_L / c^2$ is the flux flow viscosity, and σ_n is the normal state conductivity. The supercurrent, $\mathbf{j} = e n_s \mathbf{v}_s$, is directed along the y axis, and the magnetic induction is along the z axis. Solving the last equation, we get for the FLL velocity components $v_{Ly} = \chi |v_s| / (1 + \chi^2)$ and $v_{Lx} = \chi \mathrm{sign}(e) v_{Ly}$, where $\chi = j B_L / c\eta |v_s| = cj / |v_s| \sigma_n H_{c2}$ and we assume that $j > 0$, while the direction of the condensate velocity v_s depends on the sign of the charge e. Thus, the direction of the vortex drift along the current depends on the sign of the charge carriers in the superconductor. From the equation $\mathbf{E} = -\mathbf{v}_L \times \mathbf{B}_L / c$, we find the electric field components $E_x = v_{Ly} B_L / c$ and $E_y = -v_{Lx} B_L / c$, and the Hall angle $\theta_H = E_y / E_x = -\chi \mathrm{sign}(e)$.

Here, we consider a layered superconductor with alternating p and n types of conductivity having the periodicity d_c in the z direction, Fig. 1. The sign-alternating force $F_{\mathrm{H}}(z)$ and the Lorentz force $F_{\mathrm{L}} = j B_L / c$ produce the FLL displacement $\mathbf{u}_L$ in the xy plane. Within the elastic limit, this displacement can be calculated from the elasticity equation for moving FLL [9]

$$\frac{1}{c}\mathbf{j} \times \mathbf{B}_L - \frac{n_s e(z)}{c}\mathbf{v}_L \times \mathbf{B}_L + \eta \mathbf{v}_L$$
$$= -\int [C_{66}(k_x^2 + k_y^2) + C_{44} k_z^2] \mathbf{u}_L(k_z) e^{i\mathbf{kr}} \frac{d^3 k}{(2\pi)^3}, \quad (2)$$

where $e(z) = |e| \mathrm{sign}(e)$ is a periodical sign-alternating steplike function with the absolute value $|e|$ and period d_c, and C_{66} and C_{44} are the shear and tilt modulus of the FLL. Here, we consider the case when j is higher than the critical current due to pinning. In this case, we neglect the pinning force in Eq. (2) and $\dot{u}_L \ll \chi |v_s|$, which can be averaged out for fast moving FLL [9].

Instability current at $T = 0$.—At first, we consider the case of low temperatures. In this case, the displacement

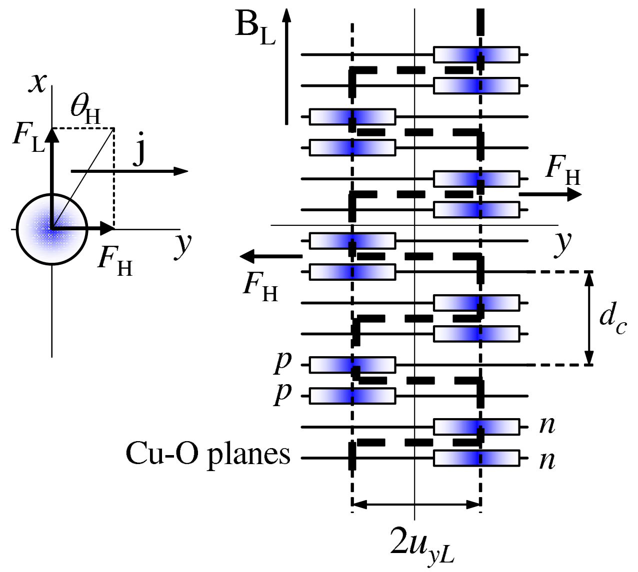

FIG. 1 (color online). Problem geometry: forces and displacements in the stack of the pancake vortices.

$\mathbf{u}_L(x, y, z)$ depends only on the z coordinate and has only a y component. Then, we derive from Eq. (2)

$$\frac{j B_L \chi e(k_z)}{c|e|} = C_{44}(k_z) k_z^2 u_{Ly}(k_z), \quad (3)$$

where $e(k_z)$ is the Fourier transform of $e(z)$. If $B_L > H_{c1}$, the main contribution to the FLL elastic energy for layered superconductors comes from the magnetic interaction of the pancakes in different layers, while the terms due to the Josephson coupling [9] as well as due to polarization effects [10] are small. In this case, we can ignore the specific symmetry of the *p-n* superconductor order parameter and electrical polarization at the respective *p-n* interfaces and express the FLL elastic modulus in the form [9]

$$C_{44} k_z^2 = U_{44} \approx \frac{B_L \Phi_0}{32\pi^2 \lambda_{ab}^4} \ln[1 + k_z^2 a_L^2], \quad C_{66} = \frac{B_L \Phi_0}{(8\pi\lambda_{ab})^2}, \quad (4)$$

where $a_L = \sqrt{2\Phi_0/\sqrt{3}B_L}$ is the FLL constant, Φ_0 is the flux quantum, and λ_{ab} is the in-plane London penetration depth. Neglecting logarithmic dependence of U_{44} on k_z and using Eqs. (3) and (4), we can approximate u_{Ly} as

$$u_{Ly}(j) \approx \pm \frac{8\pi c^2}{|v_s| \sigma_n \ln(\Phi_0 / B_L d_c^2)} \left(\frac{j}{j_d}\right)^2, \quad (5)$$

where different signs correspond to the layers with different types of charge carriers, $j_d = c\Phi_0 / 2\sqrt{2}\pi \xi_{ab} \lambda_{ab}^2 = c H_{c2} \kappa_{\mathrm{GL}} / \sqrt{2} \lambda_{ab}$ is the depairing current density in the Ginzburg-Landau (GL) approximation, ξ_{ab} is the coherence length in the **ab** plane, and $\kappa_{\mathrm{GL}} = \lambda_{ab} / \xi_{ab} \gg 1$ is the GL parameter.

If the value of the total displacement, $2|u_{Ly}(j)|$, becomes of the order of the lattice constant a_L, the FLL will be destroyed ("sliced", see Fig. 1) by the current induced sign-alternating force $F_{\rm H}$. For example, if $2|u_{Ly}(j)| = a_L/2$, the minimum of the magnetic energy corresponding to the triangular FLL disappears and some instability giving rise to the FLL decoupling occurs. We can estimate the current density j^* at which this "slice" instability arises using Lindemann criterion, $2|u_{Ly}(j^*)| = c_{\rm L} a_L$, where $c_{\rm L} \sim 0.1$ is the Lindemann constant. As a result, using Eq. (5), we obtain

$$\frac{j^*}{j_d} \approx 0.14\sqrt{\frac{c_{\rm L}|v_s|\sigma_n\sqrt{\Phi_0}\ln(\Phi_0/B_L d_c^2)}{c^2\sqrt{B_L}}}. \qquad (6)$$

We take for estimates $|v_s| = 3\times10^7$ cm/s (characteristic Fermi velocity), $\sigma_n = 10^{16}$ s^{-1}, $d_c = 3$ nm (both values characteristic for layered high-T_c superconductors), and $c_{\rm L} = 0.1$. The calculated dependence $j^*(B_L)$ is shown in Fig. 2 by a solid line. As it is seen from the figure, the FLL will be destroyed if the transport current is of the order of a few percent of the depairing current.

Alternative estimate of instability current.—The Lindemann criterion used above gives an approximate value of j^* and an alternative estimate of the instability current is of significance. We can do it using the following approach. The increase of the FLL elastic energy ΔF_L due to the periodical displacement $u_{Ly}(z)$ can be presented as $\Delta F_L = \sum_{k_z} U_{44}(k_z)|u_{Ly}(k_z)|^2/2$, where $k_z = 2\pi n/d_c$. Using Eq. (2) with $e(k_z)/|e| = 2/\pi n$, we derive with the same accuracy as Eq. (5)

$$\Delta F_L(j) = \frac{1}{3}\frac{B_L\Phi_0}{\ln(\Phi_0/B_L d_c^2)}\frac{c^4}{v_s^2\sigma_n^2\lambda_{ab}^4}\left(\frac{j}{j_d}\right)^4. \qquad (7)$$

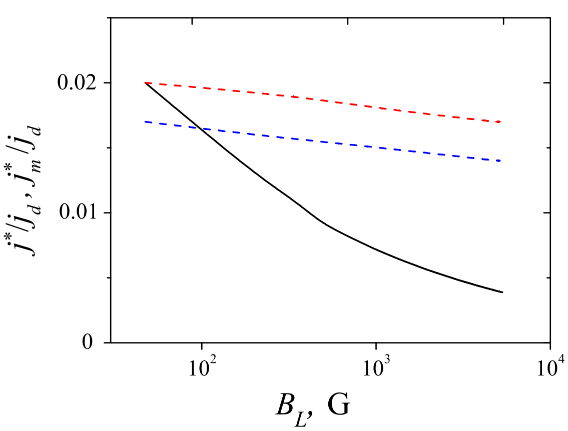

FIG. 2 (color online). The dependence of the instability currents on the magnetic induction B_L at $|v_s| = 3\times10^7$ cm/s, $\sigma_n = 10^{16}$ s^{-1}, and $d_c = 3$ nm; solid (black) line is the value j^*/j_d calculated using Eq. (6); upper (red) dashed line is the value j_m^*/j_d calculated using Eq. (8) at $\Delta M = 0.06$ G and lower (blue) dashed line is j_m^*/j_d at $\Delta M = 0.03$ G.

The melting to liquid of the FLL consisting of the ordered stacks of pancake vortices gives rise to a change ΔM in the sample magnetization [11]. The destruction of the ordered FLL by the current should also be accomplished by a magnetization change and, as a consequence, by the change of the magnetic energy $\Delta F_M = \Delta M B_L$. Then, the current that destroys FLL can be estimated using the condition $\Delta F_L(j_m^*) \sim \Delta M B_L$. In doing so, we obtain from Eq. (7)

$$\frac{j_m^*}{j_d} = \zeta\left[\frac{3\Delta M v_s^2\sigma_n^2\lambda_{ab}^4\ln(\Phi_0/B_L d_c^2)}{\Phi_0 c^4}\right]^{1/4}, \qquad (8)$$

where ζ is a constant of the order of unity. The dependence of j_m^*/j_d on B_L at $\zeta = 1$ is shown by dashed lines in Fig. 2. To calculate this function, we use the same values as for j^*, $\lambda_{ab} = 200$ nm, and $\Delta M = 0.06$ and 0.03 G, which corresponds to ΔM measured in Refs. [11,12] for FLL melting in B2212 in a wide temperature range. As is seen from the figure, values j^* and j_m^* are of the same order; however, their dependence on B_L is different. This result confirms the reliability of the crude Lindemann criterion for our goals.

Instability current at $T \neq 0$.—Here, we estimate the FLL meting temperature, $T_m(j)$, in the p-n superconductor using the Lindemann criterion as for $T = 0$. The temperature fluctuations contribute to the FLL displacement, and if the mean squared thermal displacement in the y direction, $\langle u_{Ly}^2\rangle_{\rm th}^{1/2}$, becomes comparable to the spacing between the pancakes in different layers, $a_L - 2|u_{Ly}(j)|$, then these fluctuations significantly affect the transition current j^*. The crystal is almost isotropic in the **ab** plane. Consequently, the x- and y-components of the mean squared thermal displacement are equal, $\langle u_{Ly}^2\rangle_{\rm th}^{1/2} = \langle u_{Lx}^2\rangle_{\rm th}^{1/2}$. In the elastic limit, the FLL mean squared thermal displacement in the y direction can be expressed as [13]

$$\langle u_{Ly}^2\rangle_{\rm th} = \frac{k_{\rm B}T}{2}\int\frac{d^3k}{(2\pi)^3}\frac{1}{C_{66}(k_x^2+k_y^2)+U_{44}(k_z)}, \qquad (9)$$

where $k_{\rm B}$ is the Boltzman constant. Using Eq. (4) for the elastic modulus and performing integration in the limits $0 < k_{x,y} < 2\pi/a_L$ and $0 < k_z < 2\pi/d_c$, we derive with the same accuracy as Eq. (5)

$$\langle u_{Ly}^2\rangle_{\rm th} \approx \frac{k_{\rm B}T}{16\pi^2 C_{66} d_c}\ln\left(1 + \frac{4\pi^2 C_{66}}{U_{44}a_L^2}\right), \qquad (10)$$

where U_{44} is taken at $k_z = 2\pi/d_c$. At low magnetic fields, when $4\pi^2 C_{66}/U_{44}a_L^2 = \sqrt{3}\pi^2 B_L\lambda_{ab}^2/\Phi_0\ln(\Phi_0/B_L d_c^2) \ll 1$, the thermal displacement is defined by the tilt modulus only, $\langle u_{Ly}^2\rangle_{\rm th} \approx k_{\rm B}T/4U_{44}a_L^2 d_c$, while both shear and tilt deformations are of importance at higher fields.

The temperature dependence of the crossover current density $j^*(T)$ is determined by the Lindemann equality $\langle u_{Ly}^2\rangle_{\rm th}^{1/2} + 2|u_{Ly}[j^*(T)]| = c_L a_L$, which generalizes our

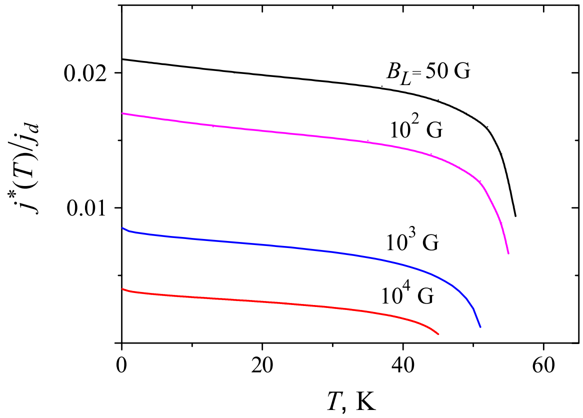

FIG. 3 (color online). The dependence $j^*(T)/j_d$ on T at different B_L; the parameters are the same as in Fig. 2 and $\lambda_{ab}^2(T) = \lambda_{ab}^2(0)/(1 - T/T_c)$, $\lambda_{ab}(0) = 200$ nm, $T_c = 60$ K.

previous result for low temperatures. From the latter equation, we get

$$\left(\frac{j^*(T)}{j_d}\right)^2 \approx \left(\frac{j^*}{j_d}\right)^2 - \frac{|v_s|\sigma_n \ln(\Phi_0/B_L d_c^2)}{16\pi c^2} \langle u_{Ly}^2(T)\rangle_{\rm th}^{1/2}, \tag{11}$$

where $\langle u_{Ly}^2(T)\rangle_{\rm th}$ is given by Eq. (10). The curves $j^*(T)$ are shown in Fig. 3 for different magnetic fields. In calculations, we use the same parameters as for the curves in Fig. 2 and $\lambda_{ab}^2(T) = \lambda_{ab}^2(0)/(1 - T/T_c)$, $\lambda_{ab}(0) = 200$ nm, $T_c = 60$ K. We assume also that v_s and σ_n are independent of T. As it follows from the figure, the effect of temperature is of importance even at relatively weak magnetic fields.

Conclusions.—We show that the unusual "slicing" instability of the pancake vortex lattice can occur in recently synthesized self-doped high-T_c layered superconductors with electron and hole charge carriers alternating in different Cu-O planes and in p-n superconductor heterostructures, e.g., in the artificial hight-T_c or low-T_c multilayers. The critical current which gives rise to this instability is significantly smaller than the depairing current j_d and the effect of the current density on the FLL melting temperature is very strong. The critical current has a pronounced dependence on magnetic field in a very broad range remaining at the same time smaller than j_d. For example, even at moderate magnetic fields ($B_L \lesssim 1$ kG), the critical current is less than 1% of j_d, and it can reduce the FLL melting temperature to zero (see Figs. 2 and 3). This allows one to seek the quantum melting of the FLL in such materials. The predicted effect may be used to shed light on the distinction between Bragg and vortex glass states in the layered superconducting heterostructures. Its existence is a signature of the vortex glass, since vortices there are mobile (rather than in the Bragg glass) and are more easily decoupled. The predicted effect can be of importance for applications. In particular, when we apply a current to a system in a magnetic field oriented perpendicular to the p-n superconducting layers, we can switch it from a superconducting to a resistive state when the current reaches the critical value found in the present Letter. This is accomplished by a pronounced jump on the current-voltage characteristics. Thus, the p-n heterostructure can be used as a nonlinear circuit element (switch, fault-current limiter). The sensitivity of such a device can be easily tuned by changing magnetic field in a broad range, from gauss to tesla, or by temperature. Our finding may also serve as an additional method for study of the electronic properties of the p-n high-T_c layered superconductors.

We acknowledge partial support from RFBR 06-02-16691, ESF AQDJJ programs, and the EPSRC via No. EP/D072581/1, EP/F005482/1.

[1] I. Bozovic, IEEE Trans. Appl. Supercond. **11**, 2686 (2001); J. Chakhalian, J. W. Freeland, H.-U. Habermeier, G. Cristiani, G. Khaliullin, M. van Veenendaal, and B. Keimer, Science **318**, 1114 (2007).
[2] H. J. H. Smilde , Ariando, D. H. A. Blank, G. J. Gerritsma, H. Hilgenkamp, and H. Rogalla, Phys. Rev. Lett. **88**, 057004 (2002).
[3] Y. Chen, A. Iyo, W. Yang, X. Zhou, D. Lu, H. Eisaki, T. P. Devereaux, Z. Hussain, and Z.-X. Shen, Phys. Rev. Lett. **97**, 236401 (2006); S. Shimizu, H. Mukuda, Y. Kitaoka, A. Iyo, Y. Tanaka, Y. Kodama, K. Tokiwa, and T. Watanabe, Phys. Rev. Lett. **98**, 257002 (2007).
[4] M. Galffy and E. Zirngiebl, Solid State Commun. **68**, 929 (1988); Y. Iye, S. Nakamura, and T. Tamegai, Physica C (Amsterdam) **159**, 616 (1989).
[5] D. LeBoeuf *et al.*, Nature (London) **450**, 533 (2007).
[6] A. K. Nguyen, A. Sudbø, and R. E. Hetzel, Phys. Rev. Lett. **77**, 1592 (1996).
[7] P. Nozières and W. F. Vinen, Philos. Mag. **14**, 667 (1966).
[8] N. B. Kopnin and G. E. Volovik, Phys. Rev. Lett. **79**, 1377 (1997).
[9] E. H. Brandt, Rep. Prog. Phys. **58**, 1465 (1995); A. E. Koshelev and P. H. Kes, Phys. Rev. B **48**, 6539 (1993); T. R. Goldin and B. Horovitz, Phys. Rev. B **58**, 9524 (1998); A. E. Koshelev, Phys. Rev. Lett. **83**, 187 (1999).
[10] D. I. Khomskii and F. V. Kusmartsev, Phys. Rev. B **46**, 14245 (1992); J. Kolacek, P. Lipavsky, and E. H. Brandt, Phys. Rev. Lett. **86**, 312 (2001).
[11] E. Zeldov, D. Majer, M. Konczykowski, V. B. Geshkenbein, V. M. Vinokur, and H. Shtrikman, Nature (London) **375**, 373 (1995).
[12] K. Kadowaki and K. Kimura, Phys. Rev. B **57**, 11674 (1998).
[13] G. Blatter, V. Geshkenbein, A. Larkin, and H. Nordborg, Phys. Rev. B **54**, 72 (1996).