Vortex Lattice in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ Well Above the First-Order Phase-Transition Boundary

J. H. S. Torres *, R. Ricardo da Silva$^1$, S. Moehlecke, Y. Kopelevich

Instituto de Física “Gleb Wataghin”, Universidade Estadual de Campinas, Unicamp 13083-970, Campinas, São Paulo, Brasil

Abstract

Measurements of non-local in-plane resistance originating from transverse vortex-vortex correlations have been performed on a Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ high-T$_c$ superconductor in a magnetic field up to 9 T applied along the crystal c-axis. Our results demonstrate that a rigid vortex lattice does exist over a broad portion of the magnetic field – temperature (H-T) phase diagram, well above the first-order transition boundary $H_{FOT}(T)$. The results also provide evidence for the vortex lattice melting and vortex liquid decoupling phase transitions, occurring above the $H_{FOT}(T)$.

Key words: A. Superconductors, D. Flux pinning and creep, D. Phase transitions
PACS: 74.72.Hs, 74.60.Ge

The knowledge of the magnetic field – temperature (H-T) phase diagram is a cornerstone of the phenomenological description of superconductors. The H-T phase diagram of conventional type-II superconductors is well known. In the Meissner-Ochsenfeld state, the surface currents screen the applied magnetic field. Above the lower critical field $H_{c1}(T)$ the applied field penetrates the superconductor in the form of an Abrikosov vortex lattice which persists up to the upper critical field $H_{c2}(T)$, where the superconductivity vanishes in the bulk of the sample. On the other hand, in high-temperature superconductors (HTS), due to strong thermal fluctuations, the occurrence of the vortex lattice melting phase transition at $H_m(T) \ll H_{c2}(T)$ has been predicted (for review articles see Ref. [1] and Ref. [2] and references therein). Since then, considerable efforts have been dedicated to identifying the melting transition in experiments. At present, it is widely believed that the vortex

* Corresponding author.

Email address: henrique@ifi.unicamp.br (J. H. S. Torres).

1 Also at FAENQUIL, 12600-000, Lorena, SP, Brasil
lattice melting is related to a first-order transition (FOT) occurring, e.g. in Bi$_2$Sr$_2$CaCu$_2$O$_8$ (Bi2212) high-Tc superconductor, at a very low field $H_{FOT}(T)$ [3]. Thus, at $T \sim T_c/2$ the $H_{FOT} \sim 500$ Oe, which is several orders of magnitude smaller than $H_{c2} \sim 100$ T. At the FOT, the equilibrium magnetization jump $\Delta M_{eq}(H,T)$ takes place [3], which, together with the Clausius-Clapeyron equation, would imply the occurrence of the entropy jump $\Delta s(H,T)$ associated with the transition. However, an alternative explanation for the magnetization (entropy) jump due to first-order depinning transition is also possible [4]. At this floating-type transition, the vortex lattice does not melt but does the opposite, it decouples from the atomic lattice becoming more ordered. The entropy jump associated with the vortex lattice floating transition has recently been obtained in Monte Carlo simulations by Gotcheva and Teitel [5]. The occurrence of the floating vortex solid phase situated between a pinned vortex solid and a vortex liquid state has been proposed two decades ago by Nelson and Halperin [6]. The FOT in HTS is accompanied by a sudden increase in the electrical resistivity [7] which demonstrates the sharpness of the vortex depinning onset. A sharp resistive magnetic-field-induced depinning transition separating disordered (low-field) from ordered (high-field) vortex states has recently been reported for NbSe$_2$, suggesting the first-order nature of the depinning transition [8]. In Bi2212, the in-plane non-local resistance, the indication of a finite shear stiffness of the vortex matter, has been measured above the FOT by Eltsev et al. [9]. It has been concluded, however, that the strong transverse vortex-vortex correlations take place in the vortex liquid phase [9].

In the present work, we report the observation of in-plane non-local resistance in Bi2212 in an applied magnetic field up to $H = 9$ T. The obtained results demonstrate that the vortex lattice does exist over a broad portion of the H-T phase diagram above the FOT boundary.

The resistance measurements were performed on a $l \times w \times t = 1.94 \times 0.28 \times 0.03 mm^3$ size Bi2212 single crystal grown using the self-flux method. The crystal characterization details as well as dc magnetization measurements have been presented elsewhere [10]. The crystal zero-field superconducting transition temperature $T_c^0 = 87.7$ K has been determined from the maximum of the temperature derivative $dR/dT$. The resistance measurements were made using PPMS (9T magnet) Quantum Design commercial equipment in magnetic fields applied along the crystal c-axis. The FOT boundary has been obtained by means of dc magnetization $M(H,T)$ measurements ($H \parallel c$ - axis) with the SQUID magnetometer MPMS5 (Quantum Design).

We used the line-electrode geometry [11,12,13] to study the non-local in-plane resistance. Six silver epoxy electrodes with contact resistance $\sim 1\Omega$ were patterned on one of the main surfaces of the crystal, as shown in the inset of Fig.1, with a separation distance $s \sim 180 \mu m$. In the experiments, the dc current $I_{14}$ was applied between the current leads 1 and 4, and the voltage was measured
Fig. 1. Temperature dependences of “primary” $R_{14,23} = V_{14,23}/I_{14}$ and “secondary” $R_{14,56} = V_{14,56}/I_{14}$ resistances measured in applied magnetic field $H = 9$ T and $I_{14} = 1$ mA; dotted line is the “parasitic” contribution to $R_{14,56}$, resulting from the current distribution effect as estimated from Eq. (1). Inset shows the geometry of the experiment.

Figure 1 shows temperature dependences of both “primary” $R_{14,23} = V_{14,23}/I_{14}$ and “secondary” $R_{14,56} = V_{14,56}/I_{14}$ resistances obtained in applied field $H = 9$ T and for $I_{14} = 1$ mA. As can be seen from Fig. 1, $R_{14,56}$ is negative in the normal state, and shows a crossover to positive values below a certain temperature within the superconducting state.

The negative $R_{14,56}$, which develops with the increase of temperature is related to the current distribution through the crystal thickness, i.e. has a local origin. Using the equation derived by van der Pauw [14]

$$R_{14,56} = -(w R_{14,23}/\pi s) \ln[(a + b)(b + c)/b(a + b + c)],$$

where $a$, $b$, and $c$ are distances between electrodes 1 and 4, 4 and 5, 5 and 6, respectively, we obtain the $R_{14,56}(T)$, depicted in Fig. 1 by a dotted line. The agreement between calculated and measured $R_{14,56}(T)$ is rather good, taking into consideration the strong crystallographic anisotropy of Bi2212, the finite width of the electrodes, and a probable distortion of the electrical potential along the line contacts. Shown in Fig. 2(a) and Fig. 2(b) are $R_{14,23}(T,H)$ and
Fig. 2. “Primary” resistance $R_{14,23}(T)$ (a) and “secondary” resistance $R_{14,56}(T)$ (b) measured at various applied magnetic fields and $I_{14} = 1$ mA. Arrows (b) indicate the field-dependent temperature $T_{\text{max}}(H)$, where a single peak in $R_{14,56}(T,H)$ takes place. Symbols in (a) and (b) correspond to the same fields (b).

$R_{14,56}(T,H)$, respectively measured at various applied magnetic fields. As can be seen from Fig. 2(b), the positive “secondary” resistance $R_{14,56}(T,H)$ emerges and increases with field. For $H > 2$ T, the $R_{14,56}(T,H)$ shows a well defined peak at the temperature $T_{\text{max}}(H)$, which decreases with the field increase.

The results presented in Figs. 3–5 unambiguously demonstrate that the positive contribution to $R_{14,56}(T,H)$ originates from the non-local resistance, i.e. is not related to the current distribution effects.

Figure 3 (a, b) shows $R_{14,56}(T)$ measured with $I_{14} = 100 \, \mu$A and $I_{14} = 1$ mA for $H = 2$ T [Fig. 3(a)] and $H = 4$ T [Fig. 3(b)]. In Fig. 4 we plotted $\Delta R_{14,56} = R_{14,56}(I_{14} = 100 \, \mu$A) $- R_{14,56}(I_{14} = 1$ mA) and $\Delta R_{14,23} = R_{14,23}(I_{14} = 100 \, \mu$A) $- R_{14,23}(I_{14} = 1$ mA) versus temperature. From Fig. 3 and Fig. 4 we note that (1) as the temperature approaches $T_{\text{max}}(H)$ from below, $R_{14,56}$
Fig. 3. “Secondary” resistance $R_{14,56}(T)$ measured with $I_{14} = 100 \mu$A (solid symbols) and $I_{14} = 1$ mA (open symbols) for $H = 2$ T (a) and $H = 4$ T (b).

becomes larger for smaller applied current, and the $R_{14,56}(T,H)$ peaks at $T \approx T_{max}$ (note that $\Delta R_{14,23} \approx 0$, i.e. $R_{14,23}$ is current-independent at the studied temperatures), (2) the current dependence of $R_{14,56}$ persists up to $\sim T_c$, i.e. is essentially related to the superconducting state. Note also that the current effect vanishes with field, so that it is negligible for $H \geq 6$ T. At $T < T_{max}(H)$ the ratio $R_{14,56}/R_{14,23}$ increases with temperature, as shown in Fig. 5, for several studied fields. All these experimental facts can hardly be understood within a local approach, indeed.

Certainly, the large enhancement of the non-local resistance with temperature over a broad temperature interval cannot be accounted for by vortex-vortex correlations in the vortex liquid. On the other hand, the long-range positive non-local resistance can arise from a correlated transverse motion of the vortex lattice [12]. The in-plane vortex-vortex correlations occurring on a millimeter scale have also been detected in Corbino-disk experiments [15,9].
Fig. 4. The difference $\Delta R_{14,56} = R_{14,56}(I_{14} = 100 \, \mu A) - R_{14,56}(I_{14} = 1 \, mA)$ (solid symbols) and $\Delta R_{14,23} = R_{14,23}(I_{14} = 100 \, \mu A) - R_{14,23}(I_{14} = 1 \, mA)$ (dotted line) versus temperature measured for $H = 2 \, T$.

Fig. 5. The “secondary” to “primary” resistance ratio $R_{14,56}/R_{14,23}$ vs. temperature obtained with $I_{14} = 1 \, mA$ illustrated for several measuring fields.
The ratio $R_{14,56}(T)/R_{14,23}(T) < 1$, see Fig. 5, can result from the vortex pinning effect which destroys the long-range positional order in the vortex lattice and, therefore, leads to a depression of the non-local resistance at large distances. As the temperature increases, both $R_{14,56}(T)$ and $R_{14,56}(T)/R_{14,23}(T)$ increase due to a vortex pinning efficiency decrease. Approaching $T_{\text{max}}(H)$, the ratio $R_{14,56}(T)/R_{14,23}(T)$ starts to decrease.

On the other hand, the occurrence of the maximum in $\Delta R_{14,56}(T)$ at $T \approx T_{\text{max}}(H)$, see Fig. 4, rules out any trivial (i.e. within a local approach) explanation of the $R_{14,56}(T)$ reduction above $T_{\text{max}}(H)$. A non-trivial origin of the maximum in $R_{14,56}(T,H)$ is supported by the observation of a splitting of this maximum into two peaks which takes place below $\sim 1$ T, see Fig. 6. The occurrence of two peaks (or hollow) in $R_{14,56}(T,H)$ at low fields can be understood assuming the “reentrant” enhancement of the vortex pinning efficiency in the temperature interval $T_{p1}(H) < T < T_{p2}(H)$, which resembles a phenomenon known as “peak effect” (PE) [16,17]. In agreement with the PE occurrence, the increase of the “primary” $R_{14,23}(T,H)$ resistance slows down at $T > T_1 \sim T_{p1}$, as Fig. 7, where $dR_{14,23}(T,H)/dT$ vs. $T$ is plotted, illustrates for $H = 0.3$ T. A similar phenomenon takes place at the temperature $T_1(H)$ just below $T_{\text{max}}(H)$ in the high-field limit, see Fig. 7 (the second peak in the derivative $dR_{14,23}(T,H)/dT$ occurring at temperature $T_2(H)$ is related to the superconductor-normal metal transition). In the low-field limit, we have also observed a current-induced suppression of $R_{14,56}(T,H)$ at $T > T_{\text{min}}(H)$, see inset in Fig. 6. This implies a similar vortex state occurring above $T_{\text{max}}(H)$ (high fields) and above $T_{\text{min}}(H)$ (low fields).

The above results are summarized in the magnetic field – temperature (H-T) diagram (Fig. 8) which we discuss now.

The FOT boundary obtained by means of dc magnetization measurements (not presented here) is plotted in Fig. 8 together with data from Ref. [7], measured for a similar Bi2212 crystal. It is evident from Fig. 8 that at $T \sim T_c/2$, the $H_{\text{FOT}}(T)$ is about 100 times smaller than the $H_{\text{max}}(T)$, implying that the long-range transverse vortex-vortex correlations persist well above the FOT boundary. This fact has a natural explanation, assuming that the floating (depinning) transition is associated with the FOT [4,5,6]. At $H < H_{\text{FOT}}(T) \sim \Phi_0/\lambda_{ab}$, the vortex lattice shear modulus decreases exponentially with field $c_{66} \approx (\varepsilon_0/\lambda_{ab}^2)(H\lambda_{ab}/\Phi_0)^{1/4} \exp[-(\Phi_0/H\lambda_{ab})^{1/2}]$, whereas at $H > H_{\text{FOT}}(T)$, $c_{66} \approx (\varepsilon_0/4\Phi_0)H$, i.e. $c_{66}$ linearly increases with field [1,2], where $\varepsilon_0 = (\Phi_0/4\pi\lambda_{ab})^2$. This implies that, at $H < H_{\text{FOT}}(T)$, the interaction between vortices and the quenched disorder overwhelms the vortex-vortex interaction, leading to a stronger vortex pinning in the low-field regime (note that at $H << H_{\text{FOT}}(T)$ the pinned vortex liquid is expected [18]). With the field increase the $c_{66}(H,T)$ and hence the inter-vortex interaction increase, and the vortex lattice de-couples from the atomic lattice at $H_{\text{FOT}}(T)$. Usually, the
Fig. 6. The “secondary” resistance $R_{14,56}(T)$ measured for several low (see text) fields with $I_{14} = 1$ mA. The inset shows $R_{14,56}(T)$ measured at $H = 0.1$ T with $I_{14} = 100$ $\mu$A (solid symbols) and $I_{14} = 1$ mA (open symbols).

Fig. 7. The derivative $dR_{14,23}/dT$ vs. $T$ demonstrating the slowing-down of the “primary” resistance increase at $T \geq T_1(H)$ which is situated just below the $T_{\text{max}}(H)$ at high fields ($H \geq 2$ T) and coincides with the $T_{p2}(H)$ at low fields ($H \leq 0.5$ T). At the temperature $T_2(H)$, transition to the normal state takes place.
Fig. 8. Magnetic field – temperature (H-T) diagram constructed on the base of experimental results (see text). The first order transition boundary $H_{FOT}(T)$ measured in both this work (□) and Ref. [7] (x) is also shown.

depinning transition in HTS is rather sharp [7]. We stress that besides theoretical expectations of a sharp depinning transition [19,20], an experimental evidence of a jumpy-like magnetic-field-induced floating transition has recently been reported [8]. At $H > H_{FOT}(T)$ and low enough temperatures, $c_{66}(H,T)$ is weakly temperature-dependent. In this regime, the vortex lattice becomes more ordered when the temperature is increased due to the suppression of the vortex pinning efficiency by thermal fluctuations, resulting in the increase of the non-local in-plane resistance with temperature. This observation is in excellent agreement with the second-order diffraction in small-angle neutron scattering experiments [21] which revealed the formation of a more ordered vortex lattice with the temperature increase for intermediate temperatures and magnetic fields. The $c_{66}(H,T)$ rapidly decreases, however, approaching either the upper critical field $H_{c2}(T)$ or the melting phase transition boundary $H_m(T) < H_{c2}(T)$. In both cases, the vortex lattice can better adjust the pinning potential [16,17] leading to the reduction of the non-local signal. There are two plausible scenarios which allow us to account for the occurrence of the minimum in $R_{14,56}(T,H)$ in the low-field regime, see Fig. 6. The first possibility is that thermal fluctuations smear out the pinning potential, improving the vortex lattice which leads to the reentrant increase of $R_{14,56}(T,H)$ with the temperature increase at $T > T_{min}(H)$ [22]. At $T = T_{p2}(H)$, the vortex lattice melts or the superconducting order parameter diminishes because of strong fluctuations in its amplitude; both effects will suppress the non-local resistance. On the other hand, the minimum in $R_{14,56}(T,H)$ occurring at the
\( T_{\text{min}}(H) \) can coincide with the melting transition temperature \( T_m(H) \) in the presence of quenched disorder [23,24,25]. Then, in a narrow temperature interval above the \( T_m(H) \), a shear viscosity due to a finite crossing energy \( U_x(H,T) \) of the entangled vortex liquid [26,27] can lead to the restoration of the non-local resistance. As temperature increases further, the \( U_x(H,T) \) vanishes [27], and the non-local resistance will be suppressed together with the entangled vortex state at the “decoupling” transition temperature \( T_D(H) = T_{p2}(H) > T_m(H) \), above which vortex fluctuations have a two-dimensional (2D) character. This second scenario agrees with the observed suppression of \( R_{14,56}(H,T) \) at \( T > T_{\text{min}}(H) \) by the applied current, assuming the occurrence of current-induced vortex cutting [28]. There is also a striking correspondence between the experimental results, see Fig. 9, and the low-field portion of the H-T phase diagram proposed by Glazman and Koshelev [29] for layered superconductors. Indeed, the \( H(T_{\text{min}}) \) can be described perfectly by a theoretical 3D “melting line” [1,2,29]\
\[ H_m(T) \approx \frac{\Phi_0 c_L^4}{(k_B T)^2} \gamma^2, \] (2)
where \( \gamma = \lambda_c / \lambda_{ab} \) is the anisotropy factor, \( \lambda_c \) is the out-of-plane penetration depth, and \( c_L = 0.1 - 0.4 \) is the Lindemann number. The Eq. (2) can be re-written in the form\
\[ H_m(T) = B(1 - t^2)^2 / t^2, \] (3)
where \( t \equiv T_m / T_{c0}, T_{c0} \) is the mean-field transition temperature, and \( B = \Phi_0 c_L^4 / 256 \pi^4 k_B T_{c0}^2 \gamma^2 \lambda_{ab}^4(0) \). The fitting gives \( B = 1.5 \) T (see Fig. 9). Taking a dimensional crossover field [29] for our crystal \( H_{3D-2D} \) \( \approx \Phi_0 / (\gamma d)^2 \sim 0.5 \) T which separates 3D (\( H < H_{3D-2D} \)) and quasi-2D (\( H > H_{3D-2D} \)) vortex fluctuation regimes, we obtain \( \gamma \approx 40 \) (here \( d = 15 \) Å is the distance between weakly coupled CuO\(_2\) bi-layers). Then, with \( \lambda_{ab}(0) \sim 1000 \) Å, one gets a reasonable value for the Lindemann number \( c_L = 0.23 \).

On the other hand, \( H(T_{p2}) \) can be best approximated by the linear dependence (see Fig. 9)\
\[ H_D(T) = C(T_c - T) / T, \] (4)
which describes the thermally induced 3D-2D vortex liquid decoupling transition in a vicinity of \( T_c \) [29]. Here \( C = \alpha_D \Phi_0^3 / d k_B T_c^4(4\pi \lambda_c)^2 \), and \( \alpha_D \) is some constant. With the fitting parameter \( C = 4.4 \) T, one has \( \alpha_D \sim 1 \). The apparent crossing of \( H_m(T) \) and \( H_D(T) \) lines seen in Fig. 9 originates from the entering into the critical superconducting fluctuations region (see, e.g., Ref. [30]).

For \( H > 0.5 \) T, \( H(T_{p1}) \) and \( H(T_{p2}) \) start to merge and, for \( H > 2 \) T, a single
transition in the vortex matter takes place at $T_{\text{max}}(H)$. For $H \gg H_{3D-2D}$, the theory [29] predicts that $H_m(T)$ approaches the melting temperature of an isolated superconducting CuO$_2$ bi-layer $T_m^{2D} \approx (k_B 8\pi \sqrt{3}/d\varepsilon_0)^{-1/d\varepsilon_0}$ according to the equation:

$$H_m(T) \approx H_{3D-2D} \exp\left\{b[T_m^{2D}/(T-T_m^{2D})]^\nu\right\},$$

(5)

where $b \sim 1$, and $\nu = 0.37$. Figure 10 demonstrates a good agreement between Eq. (5) and the experimental $H(T_{\text{max}})$ boundary at $H > 4$ T. The fitting gives $H_{3D-2D} = 0.74$ T and $T_m^{2D} = 46.3$ K ($\lambda_{ab}(0) \approx 1200 \text{ Å}$).

Thus, taking the overall data together, we are led to conclude on a possible occurrence of vortex lattice melting and “decoupling” phase transitions associated with $H(T_{\text{min}})$ and $H(T_{\rho \phi})$ low-field boundaries, respectively, as well as on the melting of a quasi-2D vortex solid which takes place along the $H(T_{\text{max}})$ boundary at $H \gg H_{3D-2D}$. The current effect on the non-local resistance measured at both $T \geq T_{\text{min}}(H)$ and $T \geq T_{\text{max}}(H)$, and its vanishing with the field increase, suggests the occurrence of the entangled vortex liquid for low and intermediate fields. We stress that the results obtained here suggest an enhancement of the vortex pinning in the vortex liquid state, being in
agreement with Refs. [23,24,25,29].

To summarize, results of the present work provide an experimental evidence for the vortex lattice existence in Bi$_2$Sr$_2$CaCu$_2$O$_8$ well above the first-order transition boundary $H_{FOT}(T)$. For the first time, the H-T phase diagram of the high-T$_c$ superconductor is constructed on the basis of direct probe of transverse vortex-vortex correlations.

We gratefully acknowledge valuable discussions with G. Carneiro, P. Esquinazi, E. Granato, V. M. Vinokur, and E. Zeldov.

This work was supported by FAPESP, CNPq and CAPES Brazilian agencies.

References

[1] G. Blatter, M. V. Feigel’man, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Rev. Mod. Phys. 66, 1125 (1994).

[2] E. H. Brandt, Rep. Prog. Phys. 58, 1465 (1995).

[3] E. Zeldov, D. Majer, M. Konczykowski, V. B. Geshkenbein, V. M. Vinokur, and H. Shtrikman, Nature 375, 373 (1995).
[4] Y. Kopelevich and P. Esquinazi, Solid State Commun. 114, 241 (2000).
[5] V. Gotcheva and S. Teitel, Phys. Rev. Lett. 86, 2126 (2001).
[6] D. R. Nelson and B. I. Halperin, Phys. Rev. B 19, 2457 (1979).
[7] D. T. Fuchs, E. Zeldov, D. Majer, R. A. Doyle, T. Tamegai, S. Ooi, and M. Koczykowski, Phys. Rev. B 54, 796 (1996).
[8] Y. Paltiel, E. Zeldov, Y. Myasoedov, M. L. Rappaport, G. Jung, S. Bhattacharya, M. J. Higgins, Z. L. Xiao, E. Y. Andrei, P. L. Gammel, and D. J. Bishop, Phys. Rev. Lett. 85, 3712 (2000).
[9] Yu. Eltsev, K. Nakao, S. Shibata, and N. Koshizuka, Proc. Int. Conf. On Materials and Mechanisms of Superconductivity, High Temperature Superconductors VI, Houston, Texas, 2000: Physica C 341-348, 1107 (2000).
[10] Y. Kopelevich, S. Moehlecke, J. H. S. Torres, R. Ricardo da Silva, and P. Esquinazi, J. Low Temp. Phys. 116, 261 (1999).
[11] H. J. Mamin, J. Clarke, and D. J. Van Harlingen, Phys. Rev. B 29, 3881 (1984).
[12] R. Wortis and D. A. Huse, Phys. Rev. B 54, 12413 (1996).
[13] Y. Kopelevich, F. Ciovacco, P. Esquinazi, and M. Lorenz, Phys. Rev. Lett. 80, 4048 (1998).
[14] L. J. van der Pauw, Philips Research Reports 13, 1 (1958).
[15] D. López, W. K. Kwok, H. Safar, R. J. Olsson, A. M. PETrean, L. Paulius, and G. W. Crabtree, Phys. Rev. Lett. 82, 1277 (1999).
[16] A. B. Pippard, Philos. Mag. 19, 217 (1969).
[17] A. I. Larkin, M. C. Marchetti, and V. M. Vinokur, Phys. Rev. Lett. 75, 2992 (1995).
[18] D. R. Nelson, Phys. Rev. Lett. 60, 1973 (1988).
[19] A. E. Koshelev and P. H. Kes, Phys. Rev. B 48, 6539 (1993).
[20] O. S. Wagner, G. Burkard, V. B. Geshkenbein, and G. Blatter, Phys. Rev. Lett. 81, 906 (1998).
[21] E. M. Forgan, M. T. Wylie, S. Lloyd, S. L. Lee, and R. Cubitt, Proc. LT-21: Czechoslovak Journal of Physics 46, 1571 (1996).
[22] Vanishing of the critical current above the PE region but slightly below the Hc2(T) was observed in Nb-O classical superconductor; Y. Kopelevich and S. Moehlecke, Phys. Rev. B 58, 2834 (1998).
[23] S. Bhattacharya and M. J. Higgins, Phys. Rev. Lett. 70, 2617 (1993); Phys. Rev. B 52, 64 (1995).
[24] C. Tang, X. Ling, S. Bhattacharya and P. M. Chaikin, Europhys. Lett. 35, 597 (1996).
[25] E. Granato, T. Ala-Nissila, and S. C. Ying, Phys. Rev. B 62, 11834 (2000).

[26] D. R. Nelson and H. S. Seung, Phys. Rev. B 39, 9153 (1989).

[27] C. Carraro and D. S. Fisher, Phys. Rev. B 51, 534 (1995).

[28] H. Safar, P. L. Gammel, D. A. Huse, S. N. Majumdar, L. F. Schneemeyer, D. J. Bishop, D. López, G. Nieva, and F. de la Cruz, Phys. Rev. Lett. 72, 1272 (1994).

[29] L. I. Glazman and A. E. Koshelev, Phys. Rev. B 43, 2835 (1991).

[30] A. K. Nguyen and A. Sudbo, Phys. Rev. B 58, 2802 (1998).