# Vortex-Pair Excitation near the Superconducting Transition of $Bi_2Sr_2CaCu_2O_8$ Crystals

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The dissipation observed at the superconducting transition of the high-$T_c$ superconductor $Bi_2Sr_2CaCu_2O_8$ is explained quantitatively by the Kosterlitz-Thouless theory of vortex-antivortex pair excitations within the $CuO_2$ planes. This conclusion is drawn from the observation of an exponential square-root singularity in the resistivity and a power-law dependence of the resistivity on magnetic field. The Kosterlitz-Thouless phase transition ($T_c = 84.7$ K) is at the $\rho = 0$ point, and the mean-field Ginzburg-Landau transition ($T_{c0} = 86.8$ K) at $\sim 25\%$ of the normal-state resistivity.

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A characteristic structural feature of the new high-$T_c$ superconductors is the presence of Cu-O planes, suggesting strongly two-dimensional (2D) physical properties. In particular, the class of oxide superconductors $Bi_2(Sr,Ca)_3Cu_2O_x$ (Bi-Sr-Ca-Cu-O),[1,2] was found to show large anisotropic behavior in the normal-state resistivity,[3,4] upper critical fields,[5-7] and transport critical currents,[8] together indicating a system of superconducting $CuO_2$ planes which are weakly coupled. A key issue in understanding the mechanism of the high-$T_c$ superconductivity is the importance of 2D fluctuations near the transition. If the system remains strongly 2D, then the planes would show behavior analogous to the thermal fluctuations in thin films of conventional superconductors, where the dissipation is associated with the motion of thermally excited pairs of vortices with opposite circulation.[9,10]

An isolated superconducting sheet would be described by the Kosterlitz-Thouless (KT) theory of phase transitions in 2D systems,[11-14] where the vortex pairs remain bound below the phase transition temperature $T_c$, which lies below the mean-field Ginzburg-Landau (GL) transition $T_{c0}$. The evidence for two-dimensionality in Bi-Sr-Ca-Cu-O is found in an exponential square-root singularity in resistivity $\rho(T)$ near $T_c$, in regions with power-law magnetoresistance below $T_c$, and in the fluctuation conductivity above $T_{c0}$. The familiar power-law *I-V* curves observed in experiments on thin superconducting films,[15] interpreted as the breaking of vortex pairs by the current,[14] were not observed because of a lack of sensitivity to experimentally accessible currents ($I \lesssim 0.2$ A). As an alternate procedure, we use weak magnetic fields to induce vortex-pair breaking near $T_c$.

The superconducting sheets in Bi-Sr-Ca-Cu-O are assumed to be comprised of the double $CuO_2$ planes spaced $d_\perp = \frac{1}{2}c = 15$ Å apart along the $c$ axis.[2] Coupling between sheets can be treated as a perturbation, as discussed theoretically by several groups.[12,16,17] Writing the vortex-pair interaction energy as $U(r_{ij}) = 2\pi K k_B T \times \ln(r_{ij}/\xi)$, where $\xi$ is the $a$-$b$-plane GL coherence distance, renormalization of the KT coupling parameter $\pi K = \phi_0^2 \zeta_\perp / 16\pi^2 k_B T \lambda_{ab}^2$ leads to the theoretical universal relation for an isolated sheet $\pi K = 2$ at $T_c$.[11,12] Here, $\lambda_{ab} = (\lambda_a \lambda_b)^{1/2}$ is the $a$-$b$-plane penetration depth and $\zeta_\perp$ is the effective thickness of the fluctuating superconducting sheets.

The coupling between layers may be treated as a symmetry-breaking interaction with the transverse coupling parameter $K_\perp \ll K$. We estimate $K_\perp / K \approx 10^{-5}$, which is the $c$-axis to $a$-$b$-plane conductivity ratio,[3] assuming isotropic scattering times, as well as the square of the critical-current-density ratio.[8] Following the argument presented by Hikami and Tsuento,[16] logarithmically interacting vortex pairs within a plane are energetically favored over multiple-plane vortex rings only below a critical pair separation $r_0 = \xi (K/K_\perp)^{1/2}$. Therefore, independent vortex pairs may be the dominant topological excitations if the length scale probed by the experiment does not exceed $r_0$. Introducing the length scale parameter $l = \ln(r/\xi)$ employed in the KT theory, the imposed constraint is equivalently expressed as $l < l_0 = \ln(r_0/\xi) \approx 6$. Ito has presented modifications of the KT renormalization-group equations treating $K_\perp$ as a mean-field perturbation.[17] By analyzing the numerical solutions, we find $l_0$ drops smoothly from 5.5 to 2.7 as $\pi K$ is varied from 2 to 6. The cutoff at $l_0$ may therefore be smaller and more restrictive at lower temperatures (large $K$) than indicated by the unrenormalized estimate. Thus the KT fixed point $\pi K = 2$ corresponding to $T = T_c$ is in principle not experimentally observable because of the constraint $l < l_0$ with $l_0$ finite.

As a consequence of the finite $l_0$, we focus on a demonstration that phase fluctuations within the $CuO_2$ layers broaden the resistance transition in Bi-Sr-Ca-Cu-O when $l < l_0$ is realized experimentally. The system was probed at finite $l$ by stimulating vortex depairing with an applied magnetic field. Here, the experimental length scale is defined as $l_H = \frac{1}{2} \ln(\eta H_{c2}/H)$, where $\eta \approx 1$ is to be determined experimentally. In the absence of a magnetic field ($H \lesssim 0.01$ Oe) thermally excited depairing is found above $T_c$. We also find that the effective thickness $\zeta_\perp$ of the fluctuating superconducting sheets is much less than the crystal thickness.

Single crystals of nominal composition $Bi_2Sr_2CaCu_2O_8$

(estimated from Rutherford backscattering analyses) used in our study were grown in a copper-oxide flux in the form of thin platelets. The growth procedure was similar to that described by Liu *et al.*,[18] except that during growth the melt was cooled from 1200 °C to 600 °C at a rate of 2 °C/h. Results described here were obtained on a crystal of dimensions $L_a = 1.41$ mm, $L_b = 0.89$ mm, and $L_c = 2$ $\mu$m. X-ray diffraction showed the crystal to be untwinned with characteristic superlattice peaks only along [0$k$0] and a weak peak at (300) but not at (030). Four contacts (25-$\mu$m-diam Au wire and using Ag paste) were made on the crystal $a$-$b$ face aligned with the crystallographic $a$ and $b$ directions, in accordance with the method described previously.[3] Resistances were measured using a phase-sensitive ac technique at 11.3 Hz for $I < 0.35$ mA. For 0.35 mA $< I < 0.2$ A, 5–15-$\mu$s current pulses at 50–100 Hz were used in order to minimize Joule heating. The in-plane resistivity components $\rho_a$ and $\rho_b$ were computed from measured resistances using our extension of Montgomery's method.[3] The nearly point-contact geometry assures well defined current injection into the crystal. We recognize that this sacrifices a precise evaluation of the current dependence of the resistivity under conditions of uniform current flow.

Figure 1(a) shows the measured temperature dependence of $\rho_a$ for magnetic fields along the $c$ direction of $H = 0 \pm 0.01$, 100, 200, and 400 Oe. The dashed line indicates the normal-state resistivity $\rho_N$ extrapolated from measurements above $T = 110$ K. Note the distinct change in the form of the temperature dependence caused by the applied field. As for the current dependence, a power-law dependence of $\rho$ on $I$ like that in thin films[15] was not found. There is an apparent decrease in $\rho_a$ and $\rho_b$ as $I$ is increased from 10 to 100 mA, believed to be from systematic error in the high-current data analysis because the Montgomery deconvolution assumes spatially uniform resistivity. There is a temperature shift in the $\rho(T)$ curves for currents above 0.1 A, indicating a depression in $T_{c0}$ due to finite critical currents.[8] On the other hand, magnetic fields of $\sim 1$ G induce noticeable pair-breaking effects. This demonstrates that the scale $r_H = (\phi_0/2\pi H)^{1/2}$ set by the magnetic field is shorter than the spacing $r_I = 2\pi K k_B T/(J\phi_0 \zeta_\perp/c)$ between vortex pairs dissociated by a current density $J$.[14] Clearly, $\zeta_\perp$ is bounded by the thickness of the crystal $L_C$ and the microscopic spacing $d_\perp$. In the absence of sensitivity to currents, expressed as $r_I > r_H$ with $I = 0.1$ A and $H = 1$ Oe, we estimate an upper bound near $T_c$ of $\zeta_\perp < 130$ Å.

For quantitative evaluation of the consistency of the data with the KT theory, we use the expression $\rho = 2\pi\xi^2 n_f \rho_N$, in accordance with the procedure of Halperin and Nelson[14] for including the Bardeen-Stephen model of vortex motion.[19] The free vortices are created from thermally excited vortex pairs of density $n_f = C\xi^{-2}e^{-2l}y$, where deviation from the Bardeen-

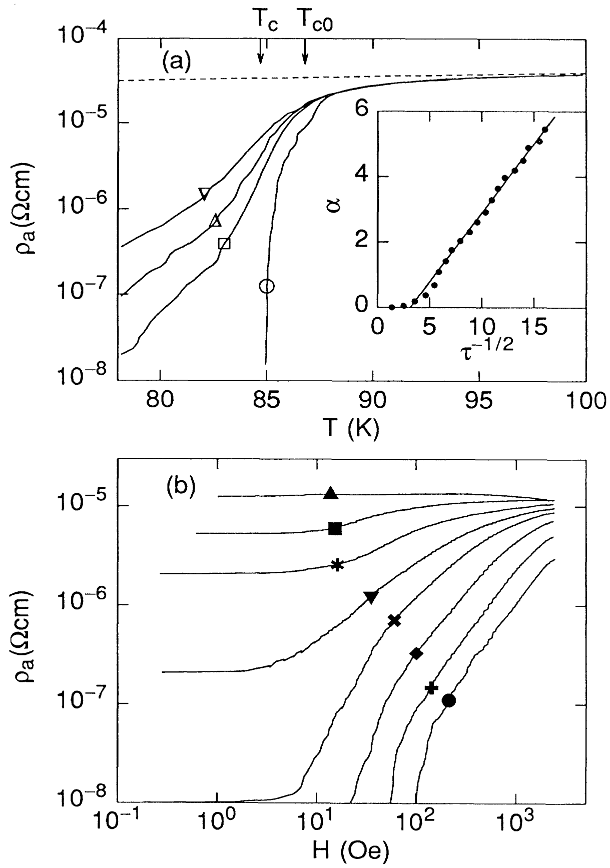


FIG. 1. (a) Temperature dependence of the resistivity component $\rho_a$ ($I = 6$ mA) with a magnetic field applied along the $c$ direction. Symbols denote $H$ (Oe): ○, 0; □, 100; △, 200; and ▽, 400. Dashed line is the extrapolated normal-state resistivity $\rho_N$. Inset: $\alpha = \ln(\rho_N/\rho)$ vs $\tau^{-1/2} = (T/T_c - 1)^{-1/2}$ obtained from the data for $H = 0$ (●) and theoretical linear fit. (b) Magnetoresistance isotherms at $T$ (K) as marked by symbols: ●, 78.2; +, 81; ◆, 83; ×, 84.4; ▼, 85; $*$, 85.7; ■, 86.6; and ▲, 87.5.

Stephen model is embodied in a constant $C \approx 1$. The vortex excitation probability $y(l)$ is calculated by renormalization theory[14]; e.g., asymptotically it is given by $2\pi y = (\pi K - 2)\exp[-(\pi K - 2)l]$ for $T \lesssim T_c$, and the coupling parameter $\pi K$ is also renormalized. For $T \lesssim T_c$, the magnetic field sets the scale parameter, denoted by $l_H$.

The theoretical picture above $T_c$ is a thermally excited 2D vortex plasma, with a correlation length diverging at $T_c$, and a corresponding scale parameter $l_+ = [b(T_{c0} - T_c)/(T - T_c)]^{1/2}$, where $b \approx 1$. This expression contains a factor for the GL temperature dependence of the underlying superconducting state.[14] The mean-field transition temperature $T_{c0} = 86.80 \pm 0.5$ K and $d_\perp \approx 6$ Å were obtained by fitting the data above $T_{c0}$ with the Aslamosov-Larkin form for the fluctuation conductivity in 2D,[9,20] $\rho^{-1} - \rho_N^{-1} = (e^2/16\hbar d_\perp)(T_{c0} - T)^{-1}$. $T_{c0}$

corresponds to $\rho/\rho_N=0.25$, which is close to the midpoint at 87 K. The dissipation resulting from thermally activated dissociation of vortex-antivortex pairs just above $T_c$ can be written as[14] $\rho/\rho_N=a\exp[-2(b\tau_c/\tau)^{1/2}]$, where $\tau_c=(T_{c0}/T_c-1)$, $\tau=(T/T_c-1)$, and $a$ and $b$ are nonuniversal constants. Using the data obtained for $H=0$ we plot the quantity $\alpha=-\ln(\rho/\rho_N)$ as a function of the reduced temperature $\tau^{-1/2}$ in the inset of Fig. 1(a) and find good agreement with the expected behavior for $4\times10^{-3}\lesssim\rho/\rho_N\lesssim0.6$ and 85 K $\lesssim T\lesssim$ 87.5 K. From a best linear fit to the data near $T_c$ we obtain the KT phase transition temperature $T_c=84.70\pm0.2$ K, $a=3.78\pm0.24$, $b=1.34\pm0.02$, and a reduced temperature shift of $\tau_c=0.025\pm0.001$. The theoretical prediction based upon a dirty-limit calculation[10,13] is $\tau_c\approx0.002$. The discrepancy with the measured value indicates that a clean-limit approximation for $\lambda_{ab}$ is more valid, as confirmed below.

The measured field dependence of the resistivity $\rho_a$ is shown in Fig. 1(b) at selected temperatures in the vicinity of $T_c$. Within a range of fields we find a power-law dependence of the magnetoresistance. Deviation from power-law behavior is observed for $H\gtrsim10^3$ Oe where the motion of the large number of free vortices dominates over the vortex-antivortex depairing as the mechanism for causing the dissipation. For $T>T_c$ the field-activated depairing crosses over to thermally activated depairing when $l_+<l_H$. Within the restricted range of fields the observed power-law behavior of the resistivity can be used to obtain $\pi K$ from the expression

$$\rho/\rho_N=2\pi Ce^{-2l}y(l)\,, \tag{1}$$

together with the KT recursion equation $dy(l)/dl=(2-\pi K)y(l)$. Making the identification $l=l_H$ yields

$$\pi K=2\frac{d[\ln\rho]}{d[\ln H]}\,. \tag{2}$$

Figure 2(a) shows results for $\pi K$ obtained from slopes of the magnetoresiatnce data of Fig. 1(b) (▼) and from $\rho_a(T)$ of Fig. 1(a) (△). We note that the temperature where $\pi K=2$ is about 0.5 K lower than the $T_c$ obtained above. We view this as satisfactory consistency, considering that $T_c$ is not precisely defined theoretically.

The dotted line in Fig. 2(a) shows the expected behavior of $\pi K$ assuming a BCS temperature dependence for $\lambda_{ab}$ and neglecting renormalization. The shift $\tau_c$ can be combined with an estimate for $\zeta_\perp$ to obtain the renormalized $\lambda_{ab}$, using the expression derived from the universal jump at $T_c$, $\lambda_{ab}^2=\zeta_\perp\phi_0^2/32\pi^2k_BT_c$. The value $\zeta_\perp=15$ Å gives $\lambda_{ab}(0)\lesssim0.4$ $\mu$m. Although the renormalization correction ($\sim$20%) is not known accurately, the favorable comparison with muon-spin-relaxation measurements in Bi-Sr-Ca-Cu-O, $\lambda(0)\approx0.2$–$0.3$ $\mu$m,[21] confirms the assumption $\zeta_\perp\approx\frac{1}{2}c$ for the thickness of the 2D superconducting sheets.

Figure 2(b) shows the scale parameter $l_H$ for a field of 200 Oe. We solved Eq. (1) for $l$ by integrating the renormalization equations.[17] The theoretical temperature dependence of $l_H=\frac{1}{2}\ln(\eta H_{c2}/H)$ is shown by the dotted curve. The constants $2\pi C=1.5$ and $\eta=0.3$ were used in Fig. 2(b); these are not precisely known theoretically, but are expected to be on the order of unity. We used the critical field slope $-dH_{c2}/dT=7.5\times10^3$ Oe/K determined by Palstra *et al.*[5]

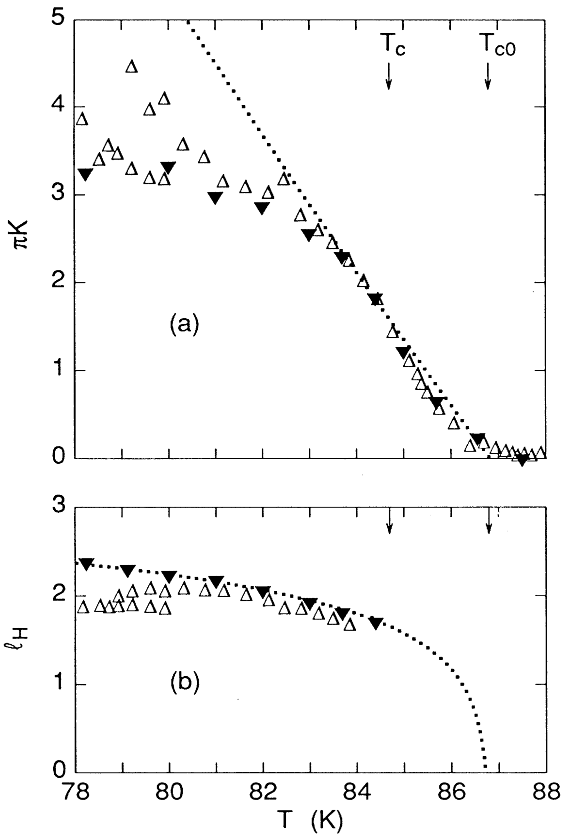


FIG. 2. (a) Vortex-pair coupling parameter $\pi K$ and (b) length scale parameter $l_H$ deconvolved from data shown in Fig. 1(a) (△) and Fig. 1(b) (▼). Dotted curves show theoretical GL temperature dependences of unrenormalized $\pi K$ and $l_H$ for $H=200$ Oe.

The implicit assumption in the KT model of a vortex fluid state in a magnetic field near $T_c$ is consistent with the recently reported vortex-lattice melting temperature of 30 K.[22] Neglecting pinning effects near $T_c$ is justified by the absence of a flux-flow critical current[23] and the weak activation for vortex motion.[24,25] In $YBa_2Cu_3O_7$ the anisotropy is $10^{-2}$, so that coupling between the planes may be too large to show vortex excitation in individual $CuO_2$ planes.[26] Thus the classic 2D behavior has been demonstrated only for thin epitaxial films of $YBa_2Cu_3O_7$.[27]

In conclusion, detailed measurements of the superconducting transition reveal several new features involving dissipation arising from excitations of two-dimensional vortex-antivortex pairs. It is shown that the Kosterlitz-

Thouless theory is applicable under the experimental conditions because the interplanar coupling is particularly weak in Bi-Sr-Ca-Cu-O. Magnetic-field–induced vortex depairing is directly observed in terms of a power-law dependence of the resistivity. An upper limit of 120 Å is established for the effective thickness $\zeta_\perp$ of the fluctuating 2D sheets. From an estimate of the penetration depth one concludes further that $\zeta_\perp$ is comparable to the spacing between $CuO_2$ double layers. Thus the resistivity approaches zero at the Kosterlitz-Thouless transition $T_c = 84.7$ K. The mean-field Ginzburg-Landau transition is estimated to be at $T_{c0} = 86.8$ K.

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