Field-Induced Suppression of the Phase Transition in $Bi_2Sr_2CaCu_2O_8$

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Magnetization measurements in field on a $Bi_2Sr_2CaCu_2O_8$ single crystal do not reveal any sharp phase transition between the normal and superconducting states. The $M(H)$ curves for temperatures about T_c in fields up to 5 T all display diamagnetic behavior. Above the mean-field transition temperature T_c^{MF}, quasi-two-dimensional diamagnetic fluctuations are responsible. Below T_c^{MF} strong deviations from the conventional mean-field (Abrikosov) behavior lead to an apparent $T_c(B)$ which *increases* with field. This effect clearly demonstrates the inadequacy of the mean-field theory to describe the vortex state below the crossover line $H_{c2}(T)$.

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It has been a general belief that the in-field transition between the normal and mixed states in a type-II superconductor is given by a second-order phase transition defined as $H_{c2}(T)$. However, recent theoretical work [1] regarding the high-T_c cuprates rather characterized $H_{c2}(T)$ as a crossover line between the normal state and a vortex liquid. A similar prediction has been formulated based on the analogy with quasi-two-dimensional (2D) magnetic insulators [2]. In the high-temperature superconductors (HTS), the vortex-liquid regime in the (B,T) phase diagram is quite extended, especially when the material, because of its large anisotropy, is quasi-2D. For example, for $Bi_2Sr_2CaCu_2O_8$ (Bi 2:2:1:2) with an anisotropic effective-mass ratio $\Gamma \gtrsim 3000$ [3], above a field $B_{2D} \approx \Phi_0/\Gamma s^2$ in the c direction the pancake vortices in adjacent CuO_2 double layers with spacing s are essentially decoupled [4]. Consequently, the vortex-lattice melting temperature in the limit of weak pinning is expected to be the same as for a 2D superconductor of thickness s ($s = 1.5$ nm); i.e., for $B > B_{2D} \approx 1$ T it melts at about 30 K. Below the melting line it is natural to consider the Abrikosov lattice as the ground state when investigating fluctuation effects [5]. For the liquid state it can no longer be expected that this is allowed [6]. The magnetic properties in the liquid state may therefore deviate from the well-known Abrikosov behavior. As we shall see below, considerable deviations are indeed observed due to fluctuations.

Note that the fluctuations suppress the melting line in a field by over 50 K from the mean-field transition $T_c \simeq 90$ K. This can be contrasted with the fluctuation effects in zero field [7] where the Kosterlitz-Thouless (KT) transition at T_c is suppressed by only 2 K below the Ginzburg-Landau (GL) transition temperature T_c^0 [8].

In this paper we report on magnetization measurements on a high-quality Bi 2:2:1:2 single crystal in a commercial SQUID system with the field along the c direction. The data are all taken in the vortex-liquid and normal regimes, i.e., in the fully reversible field-temperature regime above the irreversibility or melting line. The crystal weight and dimensions are 11.4 mg and $4 \times 4 \times 0.11$ mm^3 and its lattice parameters are $a = 0.5333$ nm, $b = 0.5485$ nm, and $c = 3.076$ nm. The composition as determined by microprobe analysis is $Bi_{2.2}Sr_{1.9}CaCu_2O_x$. Details of the preparation will be published in Ref. [9]. High-resolution electron microscopy confirmed the high purity of the crystal, for only minor traces of the 2:2:0:1 phase could be detected in a concentration of 2 out of 1000 CuO_2 double layers [10]. T_c has been determined from the ac susceptibility transition in a 2.8-μT ac field applied parallel to the a-b planes in a shielded dc environment of less than 0.3 μT. A linear extrapolation of the $\chi'(T)$ data to zero defined $T_c(0) = 88.1$ K, and that to $\chi' = -1$, the transition width of 1.5 K, demonstrating the uniformity of our sample.

In Fig. 1 our $M(T)$ data are shown for various fields ($\mathbf{H} \parallel \mathbf{c}$) between 0.1 and 5 T. The gradual decrease above 90 K indicates the effect of diamagnetic fluctuations [11]. Since this effect is negligibly small above 120 K, we corrected for a background signal by subtracting M at 120 K. The data sets between 70 and 87 K display an essentially linear behavior. Assuming linear $M(T,H)$ behavior to be valid near $H_{c2}(T)$, as is the case in Abrikosov's solution, extrapolation to $M = 0$ would determine $T_c(B)$. These intercepts are shown in the inset of Fig. 1. It is seen that the resulting "$T_c(B)$" increases with increasing B, opposite to the usual behavior. We stress here that we do not believe these data represent the mean-field $H_{c2}(T)$, but rather that they demonstrate that the simple linear construction fails, for reasons to be discussed below.

It has been recently pointed out [12] that the *linear*

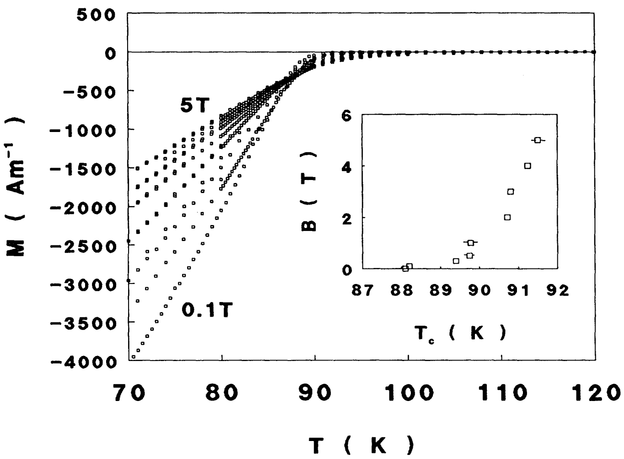

FIG. 1. Compilation of M-vs-T data of a Bi 2:2:1:2 single crystal in fields of 0.1, 0.3, 0.5, 1, 2, 3, 4, and 5 T (bottom to top) measured in both increasing and decreasing T runs. Inset: "$T_c(B)$" as determined by linearly extrapolating $M(T)$ back to $M=0$.

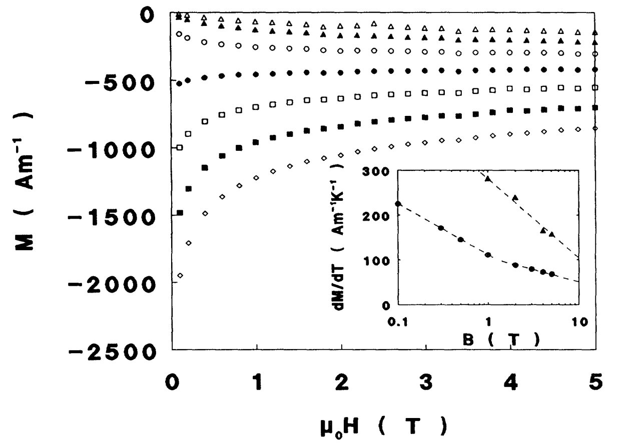

FIG. 2. Compilation of M-vs-H data above 0.1 T at 80 K (◇), 82 K (■), 84 K (□), 86 K (●), 88 K (○), 90 K (▲), and 92 K (△). Inset: A semilogarithmic plot of the slopes of the linear parts of the $M(T)$ data vs field (circles). Note the kink at about $B_{2D} \simeq 1.0$ T. For comparison the data of YBCO [13] are shown as well (triangles).

Abrikosov regime near H_{c2} is very narrow in T because of the large dH_{c2}/dT values of the HTS. Our measurements might therefore display the subsequent logarithmic regime in which $-M \propto (\Phi_0/\lambda^2)\ln(\beta H_{c2}/H)$, where β is a constant of order unity, $\lambda(t)=0.7\lambda_L(0)(1-t)^{-1/2}$ is the Ginzburg-Landau penetration depth, and $\lambda_L(0)$ is the London penetration depth at $T=0$. From the slope of dM/dT vs $\ln(H)$ the value of $\lambda_L(0)$ can be determined; see the inset of Fig. 2, where $YBa_2Cu_3O_7$ (YBCO) data [13] are also shown. The YBCO data give $\lambda_L(0) = 140 \pm 9$ nm, in good agreement with the literature [14]. Our data, however, show a kink at about 1 T which would lead to two values, namely, $\lambda_L(0)=170$ and 260 nm. In fact, a detailed comparison of our Bi 2:2:1:2 results with the computations of Ref. [12], taking $S \equiv -\mu_0 dH_{c2}/dT|_{T_c} = 1$ or 2 T/K, shows that neither the upturn of "$T_c(B)$" nor the field dependence of dM/dT can be satisfactorily explained.

Additional evidence for deviations from the mean-field theory is seen in Fig. 2 where the $M(H)$ data above 0.1 T are plotted for a series of temperatures about T_c. Results below 0.1 T and below 86 K (e.g., at 80 K) show an initial sharp decrease to $M=-7200$ A/m followed by a steep rise to -2700 A/m at 20 mT (without correcting for the demagnetization effect). This result seems in accord with the conventional magnetization about H_{c1}. The high-field behavior at 80 K also agrees with the Abrikosov prediction. Evidently, the results increasingly deviate from the canonical behavior when T is increased. A clearly detectable diamagnetic signal remains above $T_c(0)$. It shows a fast increase which levels off at about 1 T and remains almost constant up to 5 T.

The data depicted in Figs. 1 and 2 are the major results of this paper. They suggest that strong fluctuation effects play an important part in the experimental $T_c(B)$. The fluctuation diamagnetization has been reviewed in Ref. [1]. The scaling of M and χ' follows from the singular part of the free energy per unit volume, or per unit area per coupled layer when it is used in 2D:

$$F_s \propto k_B T \xi_+^{-D} F(B/\bar{B}), \quad (1)$$

where $\bar{B} \equiv \Phi_0/\xi_+^2$ and D is the dimensionality. The scaling function F is different for each of the four regimes: 2D or 3D, weak or strong fluctuations (the latter being the critical regime). The coherence length ξ_+ diverges at T_c in the critical regime between T_c and T_c^0; for $D=2$ this divergence is exponential, while for $D=3$, $\xi_+ \propto (T-T_c)^{-\nu}$, with $\nu \approx \frac{2}{3}$. The coherence length behaves classically in the weak-fluctuation regime (WFR) above T_c^0, namely, as $\xi_+ \approx \xi_{GL}(0)\tau_0^{-1/2}$, with $\tau_0 = (T-T_c^0)/T_c^0$. In 2D a convenient interpolation formula has been proposed [8]: $\xi_+ \simeq b^{-1/2}\xi_c \sinh[(b\tau_c/\tau)^{1/2}]$, with b a dimensionless constant of order unity, $\tau \equiv (T-T_c)/T_c$, $\tau_c \equiv (T_c^0 - T_c)/T_c$, and $\xi_c \approx \xi_{GL}(0)\tau_c^{-1/2}$. The susceptibility in zero field in the 2D WFR is [8]

$$-\chi_s^{2D} = (\pi\mu_0 k_B T/3\Phi_0^2)\xi_+^2 . \quad (2)$$

The scaling of M is given by

$$-M_s = \partial F_s/\partial \bar{B} \propto (k_B T \xi_+^{2-D}/\Phi_0) F'(B/\bar{B}) . \quad (3)$$

The scaling function F' saturates at large arguments for $D=2$ and grows proportional to $(B/\bar{B})^{1/2}$ in 3D. For small B, F' is linear in the argument. Low-field contours of constant M would thus reveal the prevailing dimensionality, because they are equivalent to contours of constant $B\xi_+^{4-D}$. For $D=3$ this leads to $B \propto (T-T_c^0)^{1/2}$ in the WFR and $B \propto (T-T_c)^{\nu}$ in the critical regime. Both would then have downward curvature of such contours for 3D. For $D=2$, on the other hand, the contours are straight in the WFR and curve upwards in the critical regime.

The inset of Fig. 3 shows contours of $-M=10$, 30, and 100 A/m at T above 86 K. The curvature shows that $D=2$. The dashed lines indicate that $T_c^0=88.4\pm0.4$ K in agreement with our experimental definition of $T_c(0)$. For 2D, $M_s(B)$ should saturate to about $-\overline{M}=k_BT_c^0/\Phi_0 s$, so $M=-390$ A/m for Bi 2:2:1:2: This agrees reasonably well with the value $M=-310$ A/m at 88 K and 5 T (see Fig. 2). Note also that down to 0.3 T, M indeed is constant. In Fig. 1 this corresponds to the point at which all the lines coincide. The combination of $M(B)$ being a nonzero constant at T_c^0 and the essentially linear behavior of M vs T with increasing slope for smaller fields causes "$T_c(B)$" as defined above to increase with B.

In the WFR the coupling strength between the superconducting layers was explicitly taken into account by Gerhardts [15] and Klemm, Beasley, and Luther [16]. The dimensionality is represented by a parameter $r=(8/\pi^2)\Gamma^{-1}[\xi_{\rm GL}(0)/s]^2$ which, with $\Gamma=3000$ and $\xi_{\rm GL}(0)\approx 2$ nm, is estimated to be $r=4.6\times10^{-4}$. We now compare our data with the theoretical $M(B)$ curve at T_c depicted in Fig. 5 of Ref. [15]. In the regime between $B=1$ and 5 T the magnetization is given by $-M=0.3k_BT_c/s\Phi_0$, which yields $M=-120$ A/m. This value corresponds to our experimental data slightly above $T=92$ K suggesting that this might be the mean-field (BCS) transition temperature $T_c^{\rm MF}$. According to a remark in Ref. [8], $T_c^{\rm MF}$ is shifted with respect to T_c^0 by an amount of order $\tau_cT_c^0\ln[\xi_{\rm GL}(0)/\xi_c]=0.5\tau_cT_c^0\ln\tau_c$. With $\tau_c=2.5\times10^{-2}$ [7] this amounts to $T_c^{\rm MF}-T_c^0=4.0$ K, so that $T_c^{\rm MF}=92$ K is in accord with the above analysis.

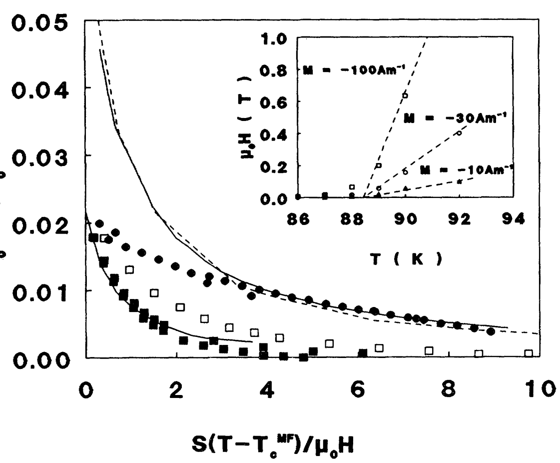

FIG. 3. Comparison of $-M\Phi_0/k_BT(\Phi_0/4\pi\Gamma B)^{1/2}$ vs $S(T-T_c^{\rm MF})/\mu_0H$ plots for experimental data at 0.3 T (●), 2 T (□), and 5 T (■) with the WFR theory [15] for layered compounds with $r=10^{-4}$ and $B/B_s=0.3$ and 5 (solid lines) and anisotropic 3D behavior (dashed line). The experimental data fit is obtained with $S=1.1$ T/K, $T_c^{\rm MF}=92.2$ K, and $B_s=1$ T. Inset: Contours of constant M above 86 K for $-M=10$, 30, and 100 A/m. The dashed lines display the expected 2D WFR behavior and define T_c^0.

In addition, we compare with the temperature dependence of M at $B=0.3B_s$ and $B=5B_s$, shown in Fig. 7(b) of Ref. [15]. For $r<10^{-2}$, the scaling field B_s is given by $B_s=3.7\Phi_0/\Gamma s^2=3.7B_{\rm 2D}$, yielding for Bi 2:2:1:2, $B_s=1.07$ T. In Fig. 3 the theory for $r=10^{-4}$ is represented by the solid lines. The symbols depict our data at 0.3, 2, and 5 T. For comparison a theoretical anisotropic 3D result is depicted by the dashed line. As was concluded in Refs. [15] and [16] the shape of the curves is not unique, but the scaling with B_s is. This horizontal scaling is very sensitive to the choice of $T_c^{\rm MF}$, and not so much to the value of S. Given the value of $B_s=1.07$ T, the best fit to the 0.3- and 5-T data was obtained with $S=1.1$ T/K and $T_c^{\rm MF}=92.2$ K (solid lines in Fig. 3). The deviation from the 0.3-T data occurs within a 1-K interval above $T_c^{\rm MF}$. For 5 T this interval is 15 times smaller. The 2-T data show behavior similar to the 5-T results in accord with the 2D nature of the fluctuations.

Both comparisons with the weak-fluctuation theory point to $T_c^{\rm MF}=92.2\pm0.5$ K. Combining this with the trend in the $T_c(B)$ plot in Fig. 1 we estimate $S_{\rm MF}=3$ to 7 T/K. Such a large value seems to disagree with the best fit in Fig. 3. However, it would be in line with Palstra's result $S=7$ T/K for YBCO [17]. The corresponding coherence length would point to local pairing, since $\xi(0)\leq1$ nm.

The temperature dependence of the Ettinghausen effect in YBCO [17] about T_c could be well described by Ullah and Dorsey [18]. Starting from the Lawrence-Doniach model [19] and incorporating the interaction between fluctuations, good agreement with the data was achieved for $S=7$ T/K. However, a breakdown of scaling below T_c could not be explained by the mean-field approach. We think this breakdown is similar to the deviating behavior of our in-field $M(T)$ data in Fig. 1. Recently, Ikeda and Tsuneto [20] computed the effect of fluctuations on $M(T)$ and found reasonable agreement with the YBCO results [13] using $S=3.7$ T/K. However, a similar computation for Bi 2:2:1:2 by only changing Γ yields $M(T)$ curves which clearly disagree with our data, especially because "$T_c(B)$" decreases with B. On the other hand, Ikeda, Ohmi, and Tsuneto [21] were able to fit their theory nicely to the in-field $R(T)$ data on Bi 2:2:1:2 single crystals [22]. Because the interpretation of $R(T)$ curves may also involve thermally activated flux motion, this agreement may be fortuitous. We therefore argue that the magnetization data provide a stricter test for theoretical models [23].

In summary, the apparent increase of "$T_c(B)$" with B actually reflects the 2D nature of the fluctuations. As the system is cooled in a field, a gradual transition occurs from the normal state to a kind of mixed-state behavior (the vortex liquid) that clearly deviates from the mean-field behavior. These deviations are greatly enhanced over what is seen in YBCO because of the extreme anisotropy of Bi 2:2:1:2. We leave it for further theoretical research to clarify these intriguing features.

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