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XY-like critical behavior of the thermodynamic and transport properties of $YBa_2Cu_3O_{7-x}$ in magnetic fields near T_c

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The heat capacity, magnetization, and electrical conductivity of single-crystal samples of $YBa_2Cu_3O_{7-x}$ have been measured and are shown to support the existence of an intermediate critical regime in the vicinity of T_c, governed by the XY-like critical exponent $\nu \simeq \frac{2}{3}$. Clear evidence is found for the divergence of the ohmic conductivity along the line $H_m(T) \propto (1-T/T_c)^{2\nu}$, the vortex melting line. The glass exponents along that line satisfy $z_g(\nu_g - 1) \simeq 6$.

In the mean-field phase diagram of a type-II superconductor, first elucidated by Abrikosov,[1] superconductivity disappears via the unusual[2] continuous melting of the vortex lattice along the line $H_{c2}(T)$. As has been pointed out repeatedly,[3-5] however, the magnetic field introduces a transverse length scale that reduces the effective dimensionality of the superconductor from dimensionality d to $d-2$. Fluctuations, which are ignored in Abrikosov's solution, are then so greatly enhanced in three-dimensional (3D) systems that $H_{c2}(T)$ no longer marks a line of phase transitions. An early treatment of the fluctuations in the Gaussian limit by Lee and Shenoy[3] was followed by many extensions,[6] mainly using perturbation methods in the "high-field" limit $|T-T_{c2}(H)| \ll H/|dH_{c2}/dT|$. A more ambitious renormalization-group calculation carried out in the same limit[2] led to an unbounded free energy and the conclusion that the transition is first order at all values of H, including zero; the latter conclusion is not supported by simulations.[7]

In this paper, we present fluctuation diamagnetism, $M(H,T)$, data taken *on the same sample* of $YBa_2Cu_3O_{7-x}$ used previously[8] for heat-capacity measurements, $C_p(H,T)$. Those data, combined with magnetoresistance data from a second sample, are shown to lend weight to a phase diagram quite different from the modified mean-field picture described above. This is a consequence of the high transition temperature, small zero-temperature coherence length, and large Ginzburg-Landau parameter of this material.[9] The zero-field critical temperature T_c is considered to be a multicritical point at the juncture of the Meissner line $H_{c1}(T)$ and the vortex melting transition $H_m(T)$.[10] We demonstrate that both $M(H,T)$ and $C_p(H,T)$ behave as expected[8,10] when the fluctuation behavior is dominated by a critical point belonging to the 3D, XY universality class. A similar conclusion followed from a crossover analysis of the zero-field heat capacity on a comparable sample.[11] The data show no distinct feature associated with $H_{c2}(T)$. To locate the melting line $H_m(T)$, we have studied the temperature-field scaling behavior of the ohmic conductivity. This analysis complements earlier work[12] on the current-temperature scaling at fixed field. The conductivity tends to diverge along $H_m(T)$, with the same exponents deduced in studies of the nonohmic properties.[12] Experimental details of the magnetic measurement on the 40-μg sample[13] will be reported separately.

Figure 1 shows the field-cooled magnetization in the vicinity of $T_c = 90.3$ K. The temperatures indicated by arrows on each isochamp correspond to the same value of the scaled temperature, as we describe below. These lie roughly at the limit of reversibility; however, zero-field-cooled data were not systematically collected. Prange[14] calculated the fluctuation contribution to the magnetization in the Gaussian approximation and predicted that

$$M_{\rm fl}(H,T)/H^{1/2} = m\,[(T/T_c - 1)/H^{1/2\nu}]\ , \qquad (1)$$

where $\nu = \frac{1}{2}$ and $m(x)$ diverges at $x = -x_c$, defining the H_{c2} line. This form was found by Gollub *et al.*[15] to hold

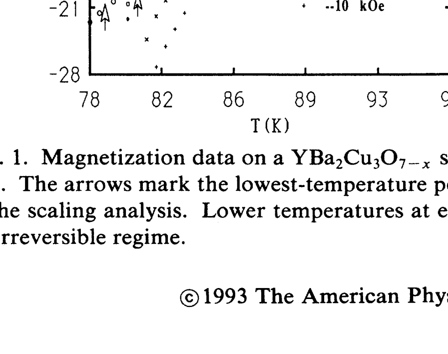

FIG. 1. Magnetization data on a $YBa_2Cu_3O_{7-x}$ single-crystal sample. The arrows mark the lowest-temperature points included in the scaling analysis. Lower temperatures at each field are in the irreversible regime.

for conventional superconductors, so long[16] as $H \leq 0.03H_{c2}(0)$. Our data lie within the range of validity, yet do not obey Eq. (1) for $\nu = \frac{1}{2}$.

For three-dimensional systems, Eq. (1) is consistent with the more general requirement that the fluctuation Gibbs free energy have the scaling form[17,10]

$$F_{\mathrm{fl}} = \xi^{-d}\mathcal{F}(H\xi^2/\Phi_0) , \tag{2}$$

where $\xi = \xi_0|t|^{-\nu}$, $t = T/T_c - 1$, Φ_0 is the flux quantum, and d is the dimensionality of the system. Other scaling forms have been proposed,[6] which go beyond the Gaussian approximation but which are valid only when $|T - T_{c2}(H)| \ll T_c H/H_{c2}(0) \simeq 2$ K for $YBa_2Cu_3O_{7-x}$ at 50 kOe, too small a range to be useful here.

Although the true critical behavior of the superconducting transition is not known, it is reasonable to expect[10] a regime where it resembles the λ transition of ^{4}He, i.e., the 3D *XY* model. Accurate estimates of the exponents for this model, obtained by renormalization-group methods,[18] are in accord with experiment[19] and we use them here: $\nu = 0.669$ and heat-capacity exponent $\alpha = 2 - d\nu = -0.007$. While such a small value of α is indistinguishable from a logarithm, it will prove useful below to retain a power-law form for the heat capacity. Shown in Fig. 2 are the data of Fig. 1 collapsed according to Eq. (1) with $1/2\nu = 0.747$. The arrows in Fig. 1 correspond to the value $x = t/H^{0.747} = -4 \times 10^{-5}$. Data with $x < -4 \times 10^{-5}$ would fall on the same curve, but are in the irreversible regime. As we will show below, the ohmic conductivity tends to diverge along a melting line located at $x = -4.2 \times 10^{-5}$ $\mathrm{Oe}^{-0.747}$, so that melting and irreversibility lines are closely correlated.

To treat the heat-capacity data in the same way, we retain the small, negative value of α at the cost of adding a nonscaling, but nonetheless critical, contribution to Eq. (2) that sets the cusp value at T_c. We then remove the nonscaling contributions by treating the difference $C(H=0,T) - C(H,T)$. This difficulty was avoided in an earlier analysis[8] through the use of a logarithmic singularity and a finite-size scaling ansatz. It is straightforward to show that

$$[C(H=0,T) - C(H,T)]H^{\alpha/2\nu} = c(x) , \tag{3}$$

with $x = t/H^{0.747}$, as above. The heat-capacity data of Ref. 8, scaled according to Eq. (3), are plotted in Fig. 3. The collapsing of the data is excellent, except in the rounding region $T_c \pm 0.2$ K. The logarithmic amplitude used in our finite-size[8] analysis can be used to obtain the cusp value at T_c. In the absence of rounding, that analysis predicts a cusp approximately ten times larger than the observed peak which, if used in place of the measured curve, would permit all the in-field data to scale. It is not only the scaling of the data in the peak (which is dominated by the low-field data) that is important, but also the "corners" near $x = -2 \times 10^{-5}$ and $x = 1 \times 10^{-5}$, where the in-field data rejoin the zero-field heat capacity.

The nonohmic resistivity of $YBa_2Cu_3O_{7-x}$ has been studied extensively,[12,20] providing solid evidence for a melting line $H_m(T)$ along which the ohmic resistance vanishes. Away from the melting line, the fluctuation contribution to the ohmic conductivity is predicted[21] to have the scaling form

$$\sigma_{\mathrm{fl}} = \xi^{2+z-d} S_{\pm}(H\xi^2) , \tag{4}$$

where z is the dynamical exponent and $S_{\pm}(y)$ are scaling functions that hold above (+) and below (−) T_c. It is more convenient to use the temperaturelike scaling variable x as in Eqs. (2) and (3), and take[10] $z = 2$ and $d = 3$, so that

$$\sigma_{\mathrm{fl}} H^{1/2} = s(t/H^{1/2\nu}) . \tag{5}$$

In this form, there is only one branch of the scaling function. For $T > T_c$, the conductivity must remain finite as $H \to 0$, so that $s(x \to +\infty) = x^{-\nu}$. At T_c the conductivity is finite except at $H = 0$, and $s(0) = \text{const}$. When weak pinning gives rise to a vortex glass, the ohmic conductivity will diverge along a line given by $x = -x_m$. The behavior close to that line is governed[10] by the glass coherence length $\xi_g(T) \propto (x + x_m)^{-\nu_g}$, which leads to the prediction[12]

$$\sigma_{\mathrm{fl}} H^{1/2} \propto (x/x_m + 1)^{-\nu_g(z_g - 1)} , \tag{6}$$

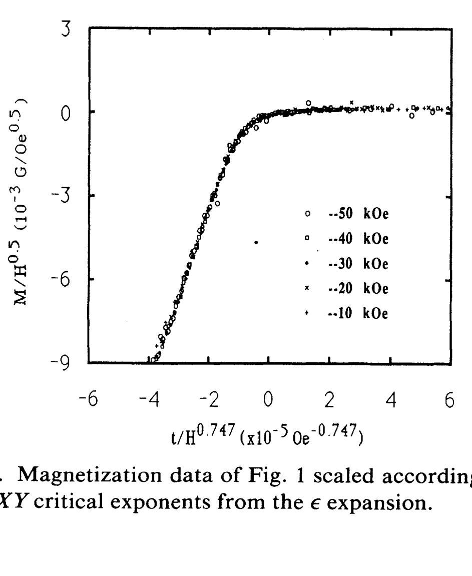

FIG. 2. Magnetization data of Fig. 1 scaled according to Eq. (1) using *XY* critical exponents from the ϵ expansion.

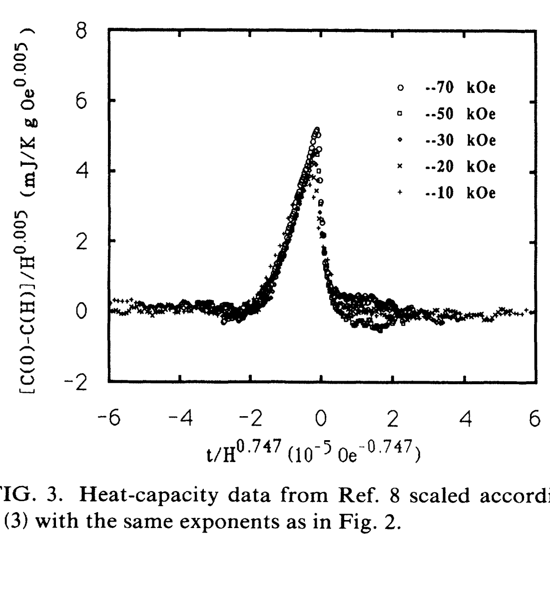

FIG. 3. Heat-capacity data from Ref. 8 scaled according to Eq. (3) with the same exponents as in Fig. 2.

where z_g and ν_g are exponents governing the glass transition. The melting line, then, is given by $H_m(T)=[(1-T/T_c)/x_m]^{2\nu}$.

The crystal used for the transport measurement was grown at Leeds by methods similar to those described previously.[22] The transition width in zero field is 0.4 K. Resistivity measurements were carried out by standard four-terminal methods using a current density of 4×10^4 A/m^2 to avoid self heating. The resistivity at 100 K is approximately 75 $\mu\Omega$ cm. Details will be published separately.[23]

To calculate the fluctuation conductivity we subtract the normal-state conductivity $\sigma_n \propto A/T$ from the measured conductivity, and scale the data according Eq. (5), with the results shown in Fig. 4 for $10 \le H \le 40$ kOe. The lowest resistivity data have been corrected for a systematic error associated with an instrumental time constant. The dashed line is $s(x) \propto x^{-0.669}$, the expected asymptotic behavior at small fields. The solid line represents the behavior predicted by Eq. (6) with $z_g(\nu_g-1)=6$ and $x_m=4.15\times10^{-5}$ Oe$^{-0.747}\equiv(695\text{ kOe})^{-0.747}$.

The results of Fig. 4 are consistent with a power-law divergence of the linear conductivity along the line $H_m(T)=(695\text{ kOe})(1-T/T_c)^{1.34}$, reasonably close to the previous values.[12,20] Note in Fig. 3, however, that there is no obvious feature in the scaled heat capacity at the point $x=-4.1\times10^{-5}$. It is also evident in Fig. 4 that the data fail to scale in a narrow region above $x=-2\times10^{-5}$, close to the "corner" in the heat-capacity curve seen in Fig. 3. The conductivity is larger in higher fields than would be expected for perfect collapsing. This excess is the unexplained "knee" in the resistivity curves which is associated with pinning. The coincidence of the two features suggests some sort of a crossover line, perhaps the residue of the mean-field H_{c2} line.

To address the applicability of the critical-point analysis, we turn to Ginzburg's famous 1960 analysis,[24] in which he introduced a characteristic reduced temperature ζ_T that depends inversely on the sixth power of the zero-temperature coherence length. When $T_c\zeta_T$ is experimentally unresolvable, the singularities regularly associated with a critical point are unobservable. For conventional superconductors, $T_c\zeta_T$ is of order 10^{-14} K, clearly beyond experimental resolution, while for Y-Ba-Cu-O we have[9] a readily resolvable $T_c\zeta_T\simeq1$ K. There has been a tendency, in discussing fluctuations in high-T_c superconductors, to take the converse of the Ginzburg criterion and argue[25] that critical fluctuations cannot be observed except in a temperature range much *narrower* than $T_c\zeta_T$. At $T_c\zeta_T$, to the contrary, higher-order fluctuation contributions exceed the Gaussian term considered by Ginsburg. Such considerations[10] lead to a revised estimate of $T_c(1\pm25\zeta_T)$ for the temperature range *outside of which* Gaussian corrections suffice. Consequently, for Y-Ba-Cu-O, the entire experimental range is inadequately described by Gaussian perturbations.

Finally, we must comment on the use of perturbation expansions about the mean-field Ginzburg-Landau solution. Prange[14] was among the first to use the Landau expansion to treat this problem, as did some later researchers.[16] The Landau expansion results in a cascade of phase transitions, the highest transition temperature T_{LLL} that of the lowest Landau level and the others lower by multiples of $\Delta T=H/|dH_{c2}/dT|$. So long as $|T-T_{\mathrm{LLL}}|\ll\Delta T$, it is generally argued that the remaining transitions can be ignored. However, in the Prange calculation, the sum over all levels brings in a leading H^{-1} factor that exactly cancels the Landau degeneracy factor, guaranteeing the scaling properties of Eq. (2). When the sum is truncated, the degeneracy factor enters as a distinct thermodynamic field H' that scales as $H'\xi^0$, rather than as $H\xi^2$. In the language of the renormalization group, the truncation process introduces a marginal operator into the problem that was not present in the original functional. As a consequence, the exponents that appear are no longer related to the usual exponents β, ν, etc.

In summary, we have demonstrated that the magnetization, heat capacity, and fluctuation conductivity of $YBa_2Cu_3O_{7-x}$ single crystals support a non-mean-field picture of the transition region of a high-temperature superconductor. In this intermediate critical regime, in the field-temperature region above the vortex melting line, the data are consistent with the critical properties of the 3D XY model, the model that describes the λ transition of ^{4}He. We have observed identical scaling in the conductivity of epitaxial Y-Ba-Cu-O films, in the magnetization of $YBa_2Cu_4O_8$, and in other materials, results of which will be reported separately.

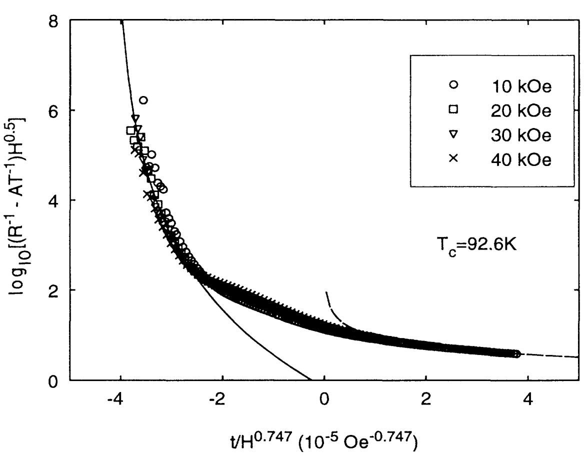

FIG. 4. Scaled conductance data. The fluctuation contribution is obtained from the measured conductance by subtracting the normal background term $(2900\text{ K}/\Omega)/T$ from the data. The dashed line is the expected asymptotic behavior $\approx x^{-\nu}$; the solid line is the divergence predicted for a vortex glass transition, $\approx(x+x_m)^{-\nu_g(z_g-1)}$, with $x_m=4.15\times10^{-5}$ Oe$^{-0.747}$ and $\nu_g(z_g-1)=6\pm0.5$.

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