Observation of separate vortex-melting and vortex-glass transitions in defect-enhanced $YBa_2Cu_3O_7$ single crystals

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We have observed two *distinct* phase transitions in a magnetic field, separating the normal and superconducting phases in single-crystal $YBa_2Cu_3O_7$ samples with enhanced disorder, either naturally occurring or deliberately introduced by radiation damage using 3-MeV protons or 1-GeV Au ions. At low currents, the *I-V* data are linear on both sides of the upper transition, whereas the lower transition separates Ohmic behavior from nonlinear (i.e., superconducting) behavior. We interpret the upper transition as the remnant of the clean system's first-order vortex-lattice melting and the lower transition as the transition into a true superconducting phase, a vortex glass.

The resistive transition of high-T_c superconductors in the presence of a magnetic field continues to be an active experimental and theoretical arena. Previous *E-J* (electric field versus current density) transport data[1–3] on samples ranging from clean (crystals) to dirty (high J_c films), observed only one transition separating a high-temperature normal metal (Ohmic *E-J* characteristics) from a low-temperature "superconducting" regime with downward curving $\ln(E)$ vs $\ln(J)$ isotherms, indicative of zero or immeasurably small linear resistance.

This paper reports evidence that for some samples with *intermediate* disorder, two separate (in-temperature) phase transitions occur between the high-temperature normal metal and the low-temperature superconducting phase, and presents some theoretical ideas about why this may occur. In our picture, the higher temperature of the two transitions is the remnant, in the presence of pinning, of the first-order vortex-lattice melting transition that would occur in the absence of pinning. For sufficiently weak pinning, we argue theoretically that a thermodynamically sharp first-order transition should continue to occur at essentially the same temperature. This is consistent with our experimental observation that the temperature (which we will refer to as T_M) of this melting transition does not change when the disorder in the sample is increased by irradiation.

As argued by Larkin,[4] arbitrarily weak pinning destroys the translational correlations of the vortex lattice on length scales longer than the Larkin pinning length L_P; thus the phase just below T_M in the presence of weak disorder *cannot* be a vortex lattice. What, then, is this phase? There are two possibilities: one is a vortex glass, with infinitely long-ranged correlations in the "Edwards-Anderson" gauge invariant order parameter and strictly zero linear resistance. The other, which might occur if these long-ranged correlations are absent, will have a nonzero linear resistance, just as in the vortex fluid above T_M. Since a vortex-glass phase is expected at still lower temperature (below T_G), in this case there is an intervening sluggish vortex-fluid regime *between* T_M and T_G. Due to the (short-range) lattice correlations in this regime, the linear resistance can be much smaller than in the vortex-fluid phase above T_M. We thus refer to this regime as a "vortex slush." The experimental data reported in this paper shows evidence for this intervening vortex slush.

It should be emphasized, though, that strictly speaking there is no symmetry difference between the vortex-slush phase and the usual vortex-liquid phase that occurs above the melting transition. The differences between the two phases is only quantitative, so although the resistivity of the slush phase would be lower than the liquid, it should still be nonzero. The first-order phase transition at T_M that separates the two is, in this sense, analogous to the liquid-vapor coexistence line, and like that line can end in a critical point, beyond which it is possible to pass continuously between the two phases without a phase transition.

A transition between the vortex-slush regime and a vortex-glass phase, on the other hand, is a second-order transition with a change of symmetry; this change of symmetry is signaled experimentally by downward curving $\ln(E)$ vs $\ln(J)$ isotherms, implying a strictly zero linear resistivity and hence true superconductivity. Because the symmetry of the vortex slush is the same as that of the vortex liquid, the universality class of the vortex-slush to vortex-glass transition is expected to be the same as that of the vortex-liquid to vortex-glass transition previously studied.[5,6]

We propose that the reason we observe two transitions in these samples while previous work has only shown one transition is intimately connected with the *intermediate* strength of the disorder in these samples. If the disorder is too strong (as in films[2] or highly twinned crystals[3]), the first-order vortex-liquid to vortex-slush transition is not present at all (i.e., the system passes continuously from the liquid to slush without a phase transition, just as a gas can become a liquid without a phase transition by moving above the liquid-gas critical point); while if the disorder is too weak, the Larkin pinning length L_P in the vortex slush becomes so long that the resistance in the vortex slush, which should vanish exponentially with increasing L_P, becomes unobservably small. In this case, the vortex-slush–vortex-glass transition, whose signature is the vanishing of the linear resistivity, cannot be seen since the resistance already appears experimentally to have vanished at T_M. Indeed, in our observations, as the magnetic-field strength is changed (which is in some ways equivalent to changing the disorder strength) we pass from regimes in which two transitions can be seen to others in which only one is observable. We will now present the experimental evidence, and then the theoretical arguments, supporting the above scenario.

The crystals were grown using a flux-melt technique[7] and had initial T_c's near 93.5 K. The samples were selected and sometimes cut with a razor blade to have barlike geometries. The length between voltage contacts was 1–2 mm and the width 0.2–1 mm. The thickness was selected to be 10–20 μm so that the irradiation damage would be uniform through the thickness of the sample. Although the crystals are nominally identical, there is a wide variation in the critical current measured by magnetization at 77 K.[8] The reason for this variation is unknown but we assume it is due to variations in the naturally occurring pinning density. This paper presents data on three such crystals. The first crystal showed *E-J* behavior similar to that reported previously.[1] At all fields it showed sudden downward curvature in the ln(E) vs ln(J) isotherms and the linear resistivity at low current dropped below the sensitivity limit within ~ 0.5 K of the onset of nonlinearity. This crystal was then irradiated with 3-MeV protons to a dose of 10^{16}cm^2. We have shown that proton irradiation greatly increases the critical current at low temperature[9] and have argued that the relevant damage for this increased pinning is random-point defects (vacancy-interstitial pairs).[10] The second crystal was not irradiated but showed a much larger critical current by transport (voltage criteria 10^{-5} V/m) than crystal 1. The third crystal also showed "clean" crystal behavior at all fields and was irradiated with 1-GeV/Au ions at the TASCC accelerator at Chalk River, Ontario to a dose of 1.15×10^{11} ions/cm^2 which corresponds to 1-ion track per vortex at a field of 2.4 T. It is expected that this damage is similar to that reported for 580-MeV Sn ions[11] and consists of columnar defects, ~ 50 Å in diameter, extending through the entire thickness of the sample parallel to the c axis (and therefore in these data parallel to the applied magnetic field). Since the energy-loss rate for this irradiation was twice that of the Sn irradiation even more continuous columnar damage is expected.

Heating is the bane of *E-J* measurements; in crystals the source is contact resistance because it is larger than the resistance of the sample. Low-resistance contacts are therefore critical for accurate *E-J* measurements over an extended current range. The crystals were first etched with 1% Br in methanol[12] and ~ 1000 Å of gold was immediately sputtered onto one of the *a-b* faces. Gold leads were attached with silver conducting epoxy (Epotek H-31) and cured at 120 °C in air. The gold not covered by the epoxy was removed with another 1% Br etch. The crystals were then heated in 1 atm. O_2 at 400 °C for 1 h and slowly cooled to room temperature. This technique consistently produced a contact resistance of less than 0.2 Ω. The data was taken in a quasi-dc manner using a Keithley 220 current source and a Keithley 182 Voltmeter. At each of 50–100 logarithmically spaced current values, the voltage was measured for both directions of the current, switching the current direction at ~ 0.3 Hz. At the lowest current 50 pairs of readings were averaged. The standard deviation of the voltage data is $\sim 2\times 10^{-9}$ V. Contact heating was evident in our data above ~ 0.02 Å, reaching a maximum of 0.02–0.08 K at the maximum current of 0.1 Å. All of the *E-J* data reported here is taken at 0.1 K intervals, and since in some cases the resistivity changes by an order of magnitude in 0.1 K, the effect of the contact heating was removed to eliminate the annoying upturn at the very highest currents. This makes it easier to concentrate on the interesting middle current range but in no way effects the results. The magnitude of the heating was established by measuring the resistance in the normal state and establishing a δT versus current function for each set of ρ-J curves. Logarithmic interpolation was then used to construct a set of curves at constant temperature. The data were taken decreasing the current, with each ρ-J curve requiring about 1 h of data taking. We looked for, but did not observe, any change in curves taken warming and cooling. We also looked for but did not observe any hysteresis in the response to a current step down to the inductive settling time of the system, ~ 5 μsec.

Figure 1 presents ρ-J data ($\rho\equiv E/J$) at 0.1-K intervals of the proton-irradiated crystal in 0.1-, 0.2-, 0.5-, and 1.0-T magnetic field applied parallel to the c axis with the current in the *a-b* plane. The data are striking. After irradiation, at low field, there are two "special" temperatures: the upper temperature which we will call T_M, separates two regimes of linear *E-J* (constant ρ as a function of J) behavior at low current. The linear resistivity drops by about an order of magnitude in a few tenths of a degree but then stops dropping as fast. Only at a lower temperature, which we call T_G, does downward curvature set in and the linear resistivity (experimentally) vanishes. Figure 2 shows the linear resistivity, on a logarithmic scale, extracted from the ρ-J isotherms, plotted as a function of temperature. Again the "two-step" nature of the resistive transition is clearly evident. As in the previous experiments, because of geometrical constraints and limited voltage sensitivity, it is not possible to prove that T_G marks the onset of true superconducting behavior. However, magnetization measurements of similar irradiated crystals show a bulk critical current approaching that of high-quality films[9] and flux-creep measurements at only slightly lower temperature are consistent with the

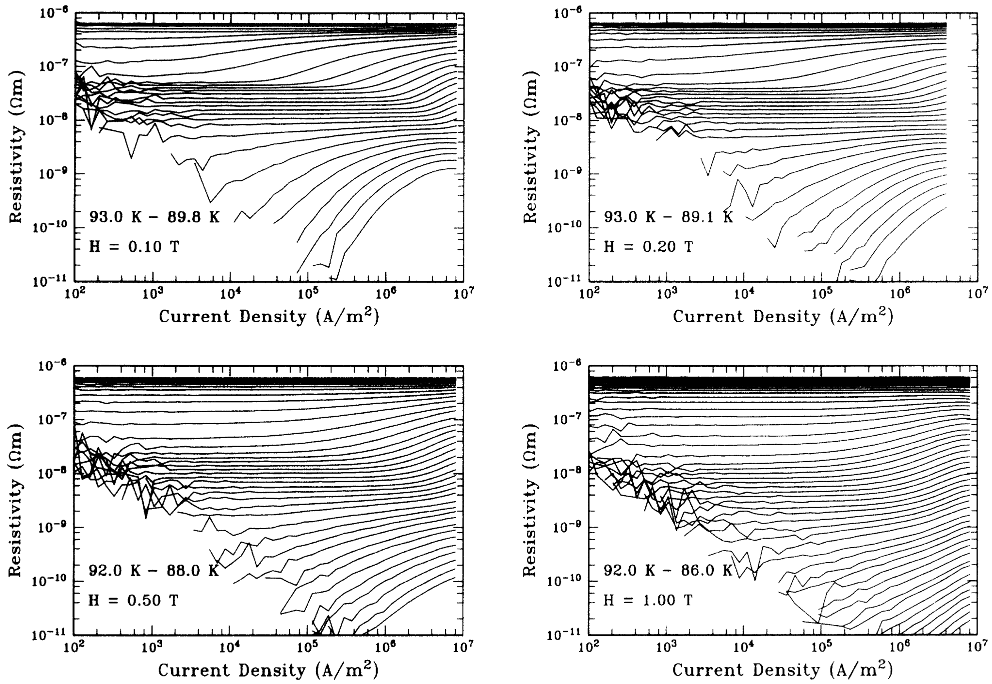

FIG. 1. Resistivity vs current density isotherms for a $YBa_2Cu_3O_{7-x}$ crystal after irradiation with 3-MeV protons at a dose of $10^{16}/cm^2$. The curves are 0.1 K apart. Data at magnetic fields of 0.1, 0.2, 0.5, and 1.0 T parallel to the c axis.

exponentially dropping E-J curves.[13]

Similar ρ-J curves are shown in Fig. 3 for an unirradiated crystal in magnetic fields of 3, 5, 7, and 9 T. Below 3 T only a single sudden change from linear to downward curving behavior is seen. Again we see two special temperatures, well separated, at 5 and 7 T. Above 7 T, the upper transition is still evident, however the lower transition appears to have moved to a lower value of resistivity and is unobservable in this experiment.

We turn now to a possible explanation of these observations. If there were no pins, we would expect a first-order vortex-lattice-melting transition at a temperature T_M below the mean-field transition T_{C2}.[14–17] Above T_M the vortex-positional correlation length is of order one intervortex spacing, indicative of the vortex-liquid regime where the thermal energy dominates. At T_M an Abrikosov lattice forms with long-range translational order (i.e., infinite vortex positional correlation length). Larkin showed that disorder (pinning) destroys the long-range crystalline correlations of the vortex lattice.[4] However, short-range crystalline order remains over a length scale L_P determined by the strength of the disorder and the elastic interaction between vortex lines.[18] We argue below that for weak disorder there may exist a temperature window between the melting temperature, T_M and the vortex-glass phase-transition temperature, T_G, where the vortex system has short-range translational order but no long-range phase order; this is the vortex-slush regime. Furthermore, for sufficiently weak disorder we expect this regime to be entered through a first-order transition at T_M, just as in the absence of pinning. As we argue below, the vortex slush has nonzero linear resistivity at low current J, because the barrier to flux-line motion does not diverge as $J \rightarrow 0$. However, as shown explicitly in Eq. (5), the *linear resistivity* can be many orders of magnitude lower than the resistivity in the vortex liquid above T_M because of the vortex-translational correlations over a distance L_P.

Whether or not the vortex-slush regime exists at weak pinning is determined by the behavior of the vortex-glass phase-transition temperature, T_G, in this limit. An indication that T_G *may* vanish in the weak-pinning limit is given by noting that the phase transition at T_G is where off-diagonal long-range order is destroyed by thermal fluctuations and that in the absence of pinning, off-diagonal long-range order is destroyed by the thermal fluctuations of the unpinned vortex lattice for arbitrarily

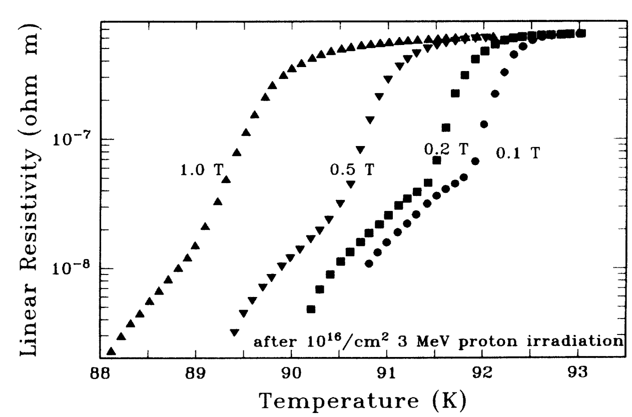

FIG. 2. The linear resistivity, extracted from the ρ vs J isotherms of Fig. 1 at low current.

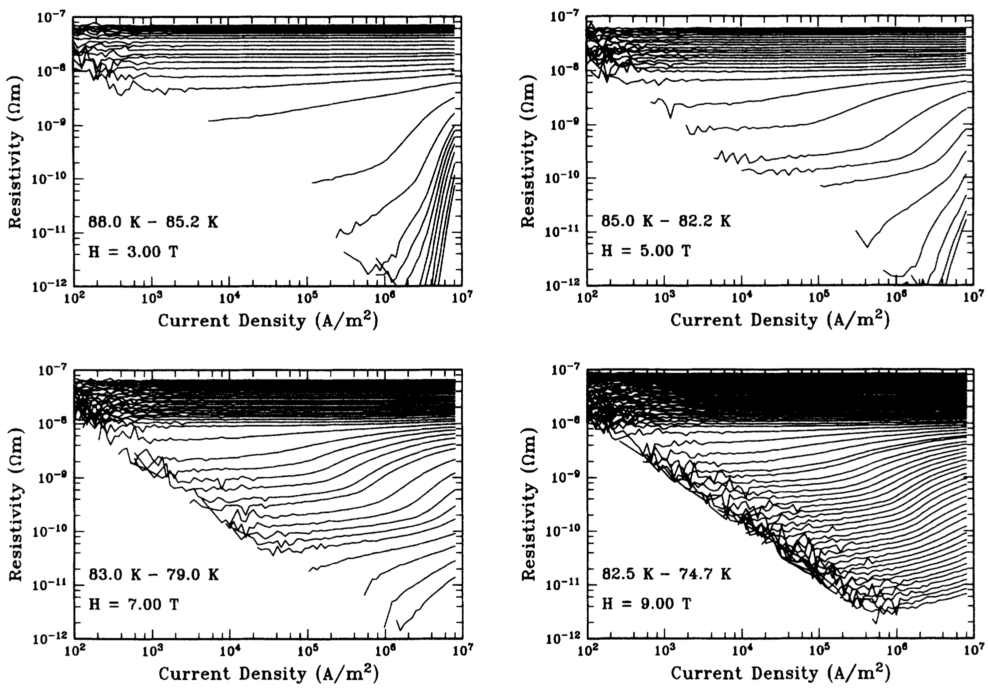

FIG. 3. Resistivity vs current density isotherms for a $YBa_2Cu_3O_{7-x}$ crystal at applied fields of 3, 5, 7, and 9 T. The curves are 0.1 K apart.

low temperature.[19] However, the limit of weak pinning is sufficiently subtle that this indication could well be misleading. Indeed, we have not been able to reliably estimate T_G in this limit of weak pinning. An honest calculation of T_G, as discussed below, would involve a calculation of the energy of topological defects in the vortex-glass state. In the absence of such a calculation, we can try to estimate T_G using an oversimplified spin-wave theory, which we will outline below. The vortex-glass-transition temperature, T_G^{SW}, predicted by the spin-wave theory does, indeed, vanish in the limit of weak pinning; this at least suggests that the actual T_G also vanishes in this limit. However this argument is less than compelling, for reasons that will be elaborated on below.

The spin-wave argument proceeds as follows: In the local, isotropic elastic approximation, the thermal energy to melt the vortex lattice, $k_B T_M$, can be approximated as the elastic energy to strain a vortex line by a_0, the separation between lines, in a length a_0,

$$k_B T_M \simeq K_{\mathrm{el}} a_0^3 \ . \tag{1}$$

Here the elastic properties of the vortex lattice have been approximated by a single elastic constant K_{el}. Now we will estimate the vortex-glass-transition temperature using the (potentially erroneous) assumption that the thermal energy necessary to melt the vortex glass can be estimated as the energy necessary to change the relative phase between adjacent Larkin domains by π, thereby apparently destroying the phase relationship of the domains. The relative phase of a domain is determined by the position of each vortex line in the domain so a large phase change can be achieved by a small movement of each line. The number of lines is $(L_p/a_0)^2$. Thus a displacement of each line by $\delta \simeq a_0^2/L_p$ will change the phase of the domain by 1. Each vortex line is at local minimum of the combined elastic and pinning energy so again we approximate the strain energy with a single elastic constant. The strain energy in a Larkin domain, $K_{\mathrm{el}}(\delta/L_p)^2 L_p^3$, gives us the spin-wave estimate for T_G. Combining this with Eq. (1), we get

$$k_B T_{\mathrm{vg}}^{\mathrm{sw}} \simeq k T_M \frac{a_0}{L_p} \ . \tag{2}$$

Thus we see that for sufficiently weak disorder (L_P sufficiently large), this naive argument suggests that $T_G \ll T_M$ and a window exists between the melting temperature and the vortex-glass phase transition.

The potential flaw in the above logic is that we have used a spin-wave theory for T_G; that is, we have assumed that the energy required to disorder the system can be obtained from the energy needed to produce a smooth, continuous, spin wave of sufficient amplitude. In fact, to disorder the system one must introduce topological defects, whose energy may be much higher. The most dramatic example of this is the simple Gaussian model, $H = K(\phi_i - \phi_j)^2$ where the field ϕ runs from $-\infty$ to ∞. A naive argument like that presented above would suggest that this model disorders at a temperature $T_C \simeq K/k_B$; whereas, in fact, it can trivially be shown that this model remains ordered at all temperatures. Thus, these simple spin-wave arguments can sometimes greatly underestimate T_c, and so must be used with caution.

A more honest calculation of T_G would begin as follows: The weakly pinned vortex lattice at equilibrium has local strains of order a_0/L_P. These strains accumulate and the system thus presumably has a network of dislocation lines present with typical spacing L_D. Naively, L_D is of order L_P, since this is the distance over which total distortions of order a_0 occur, but L_D could be significantly larger than L_P; how to systematically address this question is at present unclear. Now let us consider the typical free-energy cost, $\Delta F(l)$, to move a single vortex line segment of length of order l a distance of order l somewhere in a given region of size of order β. If, for some l, this free energy is small enough (below some threshold that is presumably of order $k_B T$), then these vortices will be able to move at equilibrium, destroying any possible vortex-glass long-range order and producing a vortex slush (or liquid) with nonzero linear resistivity. For weak pinning, the smallest length scale on which one might readily move vortices is L_D. Consider a dislocation loop of size L_D that is present in the weakly pinned lattice. By removing a vortex line along one-half of this loop and adding it to the other half, the dislocation locally "climbs" by one lattice unit. Since the pinning is weak, the average free-energy cost per unit line length of this climb is small and it will vanish as the pinning vanishes. However, L_D diverges in this limit and we are interested in the total free-energy cost over the whole loop, $\Delta F(L_D)$, since the vortex-glass-transition temperature T_G is expected to be of order $\Delta F(L_D)/k_B$. How this behaves in the weak-pinning limit is, at present, unclear since it requires understanding not only what determines L_D but also the energetics of climbing these weakly pinned dislocation lines by small amounts.

If, for temperature just below T_M, this free-energy cost, $\Delta F(L_D)$, to move a vortex line on the scale of L_D is smaller than $k_B T_M$, then the vortex-slush regime should be present below T_M. Only at a lower vortex-glass-transition temperature, $T_G \sim \Delta F(L_D)/k_B$, will the resistance drop to zero.

The linear resistance in the (assumed) vortex-slush regime will be determined by the free-energy barriers F_B which must be surmounted to move a length L_D segment of vortex line from one side of a dislocation loop to the other. This free-energy barrier may, in principle, be *much* bigger than the free-energy cost of the excitation itself, i.e., $\Delta F(L_D)$. Indeed, since the barrier involves moving a vortex line through a region of nominally perfect vortex lattice of size L_D, one might expect $F_B(L_D) \simeq \varepsilon L_D$, where $\varepsilon \simeq \Phi_0^2/8\pi^2\lambda^2$ is the line tension of an interstitial vortex line in a perfect lattice. [Within factors of $\ln(\kappa)$, this line tension is comparable to that of a single vortex line inserted into an otherwise vortex-less system.] Since the vortex-lattice melting temperature is of order the energy of a length a_0 segment of vortex line, $k_B T_M \simeq \varepsilon a_0$, the dimensionless free-energy barrier can be very large for weak pinning

$$\frac{F_B(L_D)}{k_B T_M} \sim \frac{L_D}{a_0} \ . \tag{3}$$

The linear resistance in the vortex slush arises from thermal excitation over these barriers and can be expressed as

$$R_{\text{slush}} = R_{\text{liquid}} e^{-F_B(L_D)/k_B T} \tag{4}$$

or

$$R_{\text{slush}}(T_M^-) = R_{\text{liquid}} e^{-cL_D/a_0} \ , \tag{5}$$

where the result (3) for F_B implies that $c = O(1)$. A constant c of order 1 is also reasonable on physical grounds since in highly disordered samples with $L_D \simeq L_P \simeq a_0$, we do not expect an appreciable drop in resistance at T_M.

This picture provides a good qualitative explanation for our data. If the defect density is low, i.e., weak pinning, L_D/a_0 will be large and the resistance in the slush phase will be below the resolution of our experiment. The resistance will appear to vanish suddenly because, as will be discussed below, the first-order vortex-lattice-melting transition remains sharp under moderate disorder even though the vortex lattice itself is destroyed. For the proton-irradiated sample, (Fig. 1), the defect density is very high[10] and so presumably L_D is small and there is only a modest drop in the linear resistivity at T_M. This picture also explains the absence of a second linear regime in $YBa_xCu_{1-x}O$ films. There the vortex-glass scaling indicates that the films are always in the limit where $L_P/a_0 \simeq 1$,[2] and therefore no change at the melting temperature would be expected.

For a given sample with fixed disorder, L_P will presumably have a magnetic-field dependence. There are a number of competing effects which may make L_P/a_0 nonmonotonic in field. However, the exact dependence of L_P/a_0 on field along the melting curve is difficult to calculate for a number of reasons, not the least of which being the need for a quantitatively correct prediction for the melting locus $T_M(H)$, which does not appear to be currently available. Nonetheless, the following arguments, hopefully, give some idea of the field dependence of L_P/a_0: As the field increases along the melting line, the decreasing effectiveness of thermal smearing[6] will tend to decrease L_P. The thermally induced vibrations of the vortex lines will significantly decrease the pinning interaction whenever the amplitude of vibration is greater than the size of the pinning defect.[20] For point defects, from either proton irradiation or chemical impurities, or our columnar damage tracks, this is $\sim\xi$, the coherence length. In a Lindemann formulation, the thermal vibration amplitude is $\sim 0.1a_0$ at melting.[16] This implies that, along the melting curve, for $B \lesssim 0.5$ T the vibration amplitude will be greater than ξ and the pinning will be significantly weakened. However, increasing the field decreases the lattice constant, and therefore decreases the absolute vibration at melting and therefore increases the effectiveness of the pinning sites. For this reason we expect an eventual decrease in the size of the Larkin domains, L_P, which form just below the melting temperature, as the field is increased. This is manifested in our data (see Fig. 3) by a decrease in the drop in linear resistance upon entering the slush phase. At 1 T, only a sudden vanishing of the resistivity is observed, however, as the field is increased a second linear regime develops out of the noise floor and is clearly present in the ρ vs J curves at 5 T and above. The observation that the magnitude of the drop in resistivity at melting decreases as the field increases is also consistent with this picture.

At very low field, as the vortex lines get further apart, $a_0 \simeq \lambda$, the elastic interaction weakens, eventually decreasing L_P/a_0 as the applied field decreases. Unfortunately, the melting locus is expected to have a certain field strength below which T_M drops with field;[14] this turn-around also occurs when $a_0 \simeq \lambda$. Hence, without a detailed theory, it is difficult to predict whether this regime of decreasing L_P/a_0 is ever accessible. If it is, as the data in Fig. 4 suggests, then for a particular strength of disorder and a particular experimental resistance sensitivity, a vortex-slush regime might be visible at both low and high field but unobservable (or absent) at intermediate-field strengths. This behavior is indeed evident in the data taken on the heavy ion-irradiated crystal shown in Fig. 4. At 500 G an Ohmic regime is visible below T_M, indicative of a vortex slush. Upon increasing the field, only a sudden vanishing of the linear resistance is observed in the range 1–3 T. However at 4 T and above, the vortex-slush regime becomes visible again.

There are several indications that the upper transition is indeed a remnant of the pure system's vortex-lattice melting. The first is the observation that it does not move upon irradiation (see Fig. 6). Moreover, vortex-lattice melting is predicted to be first order[17,21] and first-order phase transitions remain thermodynamically sharp in the presence of sufficiently weak disorder. This can be demonstrated using the following argument due to Imry and Wortis.[22] In order for the transition to be rounded by disorder, at the pure system's coexistence line, the system must break up into a mixture of vortex-liquid and vortex-slush regions. We can argue that for weak disorder this is energetically unfavorable, as follows: Consider a slushy region surrounded by liquid ones. The liquid to slush interface has a width $\sim\xi_T(T_M)$, the translational correlation length in the liquid at melting. Because the pure system's melting transition is first order, this width does not diverge at T_M, and hence the interface is sharp. Therefore, the interfacial energy cost E_{cost} of creating such an island of slush is proportional to the surface area of the slushy region; i.e., $E_{\text{cost}} \propto L_d^2$, where L_d is the size of the slushy domain ($L_d > L_P$). Competing against this is the energy gained by the random pinning, which favors creation of a slushy domain (in some places) due purely to statistical fluctuations. On statistical grounds, we expect this energy gain due to the pinning to be proportional to the *square root* of the volume and the pinning strength, V_P. Thus $E_{\text{gain}} \propto V_P\sqrt{L_d^3}$. Hence the ratio $E_{\text{gain}}/E_{\text{cost}} \propto V_P L_d^{3/2}/L_d^2 \propto V_P/\sqrt{L_d}$, which is always small for weak pinning. Hence, for sufficiently weak disorder, it costs far more energy to break the system up into domains; therefore, the system remains monodomain, and the first-order vortex liquid to slush transition remains thermodynamically sharp. Furthermore, since the physics of the transition all happens on the length scale $\xi_T(T_M)$, which is finite, and therefore less than the pinning length L_P for sufficiently weak disorder, the temperature at which the transition occurs should not be strongly altered by weak disorder.

We observe that the upper transition remains sharp except for the proton-irradiated sample where L_P is presumed small. Figure 5 shows four ρ vs J curves at 4 T

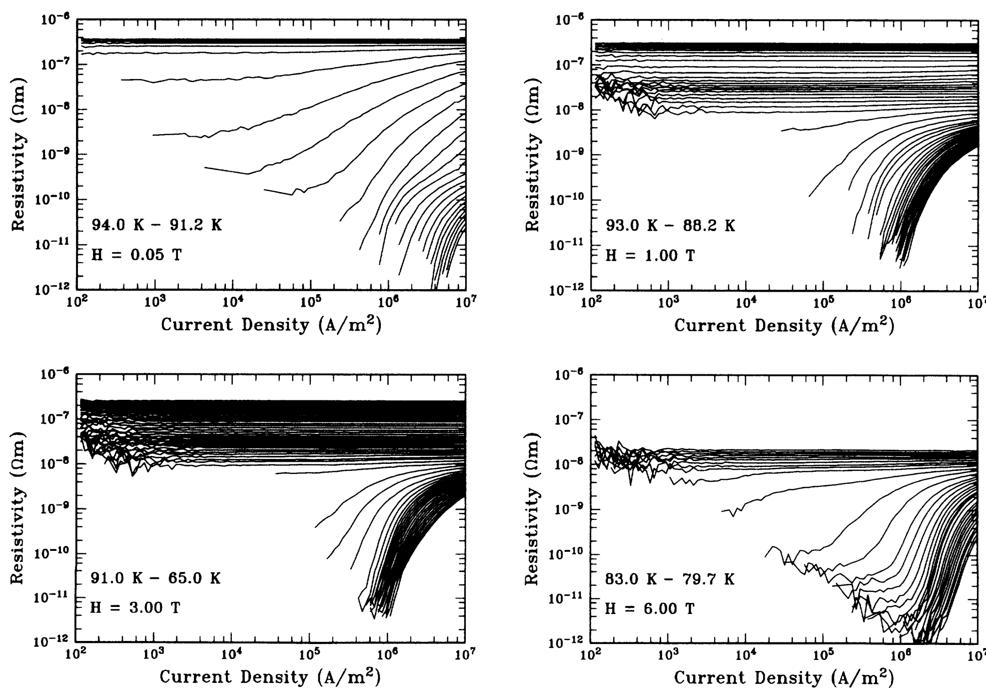

FIG. 4. Resistivity vs current density isotherms for a $YBa_2Cu_3O_{7-x}$ crystal after irradiation with 1-GeV Au ions at a dose of $B_\Phi = 2.4$ F. Data at magnetic fields of 0.05, 1.0, 3.0, and 6.0 T parallel to the c axis. The curves are 0.1 K apart.

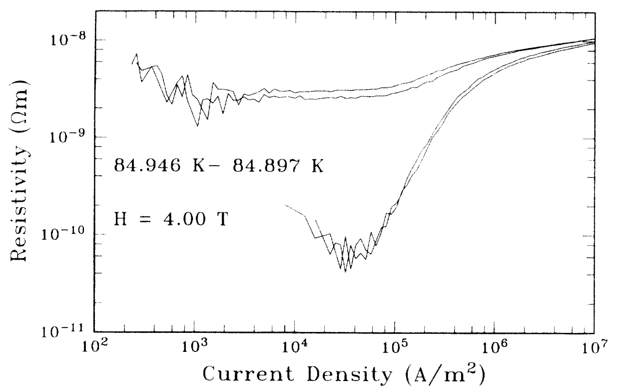

FIG. 5. Four resistivity vs J isotherms for the heavy ion-irradiated sample at an applied field of 4 T. The temperatures are 84.95, 84.93, 84.91, and 84.90 K from top to bottom.

for the heavy ion-irradiated sample. The curves are separated from each other by only 20 mK, yet between the two middle curves there is a drop in linear resistivity of almost a factor of 50. Figure 5 shows the melting temperature measured by selecting the first obviously nonlinear ρ vs J isotherm, versus magnetic field before (squares) and after (circles) irradiation. They coincide remarkably well. In fact T_M in all of the crystals with similar T_c's ($\sim$93.5 K) lies on this curve both before and after moderate doses of irradiation. This is a clear indication that this line is unaffected by the defect density and suggests that it is a vestige of the pure system's Abriskov vortex-lattice-melting line. We observe a field dependence proportional to $(T_c - T)^{3/2}$ which is shown as the line in Fig. 6. In earlier work on an untwinned single crystal of $YBa_2Cu_3O_7$, Farrell, Rice, and Ginsberg[23] present evidence for melting based on a dissipation peak in a torsional oscillator experiment. Our transport data tends to support their interpretation, and moreover enables us to distinguish between vortex-slush and vortex-glass phases below T_M in samples with intermediate disorder.

Further evidence for this vortex-melting picture comes from the s-shaped feature evident in the ρ vs J curves in data sets which exhibit a vortex-slush regime at low current [see especially in Fig. 1(a)]. We believe this feature, which connects the vortex-slush resistivity at low current to the vortex-liquid resistivity at high current can be interpreted as a *current-induced "melting"* at fixed temperature of the vortex-lattice domains in the slush phase. The basic idea is that, as the Magnus force due to the current pushes the vortex-lattice domains across the sample, the vortex lines encounter random forces as they pass over pinning sites. Because of the motion of the vortex lines, these forces are time dependent, and act (crudely speaking) like an additional random thermal noise agitating the lattice. As a result, the faster the lattice moves, the higher its "effective temperature." At sufficiently large currents, the increase in effective temperature is sufficient to raise the temperature from the true, equilibrium temperature to above the melting temperature thereby "melting" the lattice. For larger currents, the lines move as a vortex liquid, rather than vortex slush, and hence show a higher resistivity, as we observe experimentally. Clearly, the closer the temperature is to T_M, the smaller the current necessary to "melt" the crystal (since the required increase in effective temperature is smaller). This also agrees qualitatively with our observations as shown in Fig. 7, where the melting current [defined as the current where $\partial \ln(\rho)/\partial \ln(J)$ is a maximum] is plotted as a function of temperature for the two irradiated crystals. It is noteworthy that the slope of the melting current versus temperature, below T_M, is essentially independent of the strength and type of the disorder.

We also note that there is a preliminary report of a similar sudden drop in linear resistivity in untwinned $YBa_2Cu_3O_7$ crystals, accompanied by hysteresis in decreasing versus increasing temperature sweeps of up to 40 mK.[24] We looked for but did not observe hysteresis greater than our thermometer resolution of $\sim$10 mK.

In summary, we have reported evidence, in transport

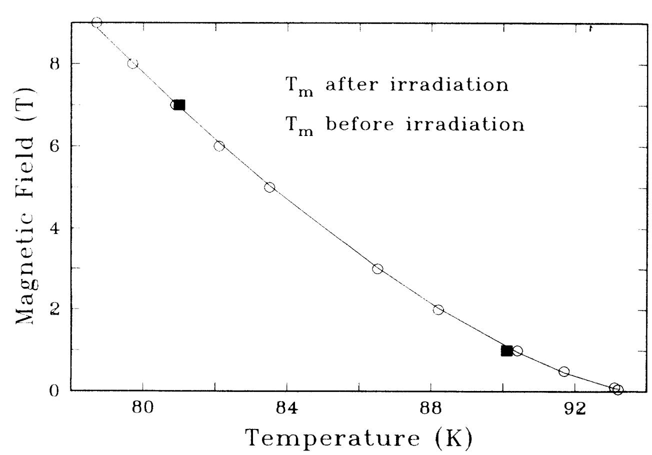

FIG. 6. The field dependence of the melting temperature for the heavy ion-irradiated crystal before (squares) and after irradiation (circles).

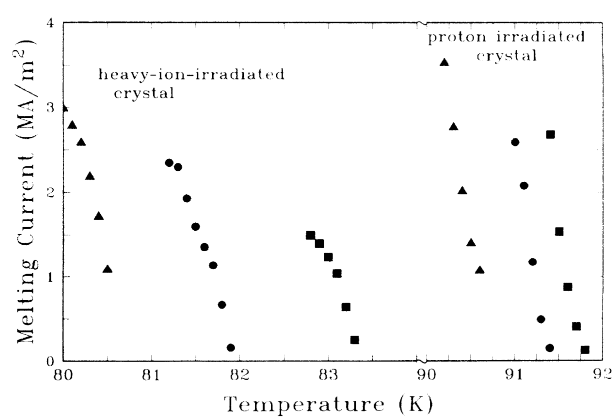

FIG. 7. The melting current in MA/m² vs temperature defined as the current where $\partial \ln(\rho)/\partial \ln(J)$ is maximal. Data for crystal 1 at applied fields of 0.1, 0.2, and 0.5 T are shown on the right as squares, circles, and triangles. Data for crystal 3 after heavy ion-irradiation is shown on the left at 5, 6, and 7 T.

behavior, for two separate transitions near the irreversibility line in $YBa_2Cu_3O_7$ single crystals. We interpret the upper transition as a remnant of the first-order vortex-lattice-melting transition, frustrated by disorder, and the lower as a vortex-glass transition into a superconducting phase. In clean samples the melting transition may be the only visible transition, if the drop in linear resistivity at the melting temperature results in resistivity level below the experimental threshold. In very disordered samples, only the vortex-glass transition would be expected because the samples are always in the $L_p/a_0 \simeq 1$ limit where the melting transition is completely rounded. However, if the disorder is adjusted properly, two well-separated transitions are visible.

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[1]T. K. Worthington, F. Holtzberg, and C. A. Feild, Cryogenics **30**, 417 (1990).

[2]R. H. Koch, V. Foglietti, W. J. Gallagher, G. Koren, A. Gupta, and M. P. A. Fisher, Phys. Rev. Lett. **63**, 1511 (1989).

[3]P. L. Gammel, L. F. Schneemeyer, and D. J. Bishop, Phys. Rev. Lett. **66**, 953 (1991).

[4]A. I. Larkin, Zh. Eksp. Teor. Fiz. **58**, 1466 (1970) [Sov. Phys. JETP **31**, 784 (1970)].

[5]M. P. A. Fisher, Phys. Rev. Lett. **62**, 1415 (1989).

[6]D. S. Fisher, M. P. A. Fisher, and D. A. Huse, Phys. Rev. B **43**, 130 (1991).

[7]D. L. Kaiser, F. Holtzberg, M. F. Chisholm, and T. K. Worthington, J. Cryst. Growth **85**, 593 (1987).

[8]L. Civale (unpublished).

[9]L. Civale, A. D. Marwick, M. W. McElfresh, T. K. Worthington, A. P. Malozemoff, F. Holtzberg, J. R. Thompson, and M. A. Kirk, Phys. Rev. Lett. **65**, 1164 (1990).

[10]L. Civale, M. W. McElfresh, A D. Marwick, T. K. Worthington, A. P. Malozemoff, F. Holtzberg, C. A. Feild, J. R. Thompson, D. K. Christen, and M. A. Kirk, in *Progress in High Temperature Superconductivity* (World Scientific, Singapore, 1991), Vol. 31, p. 25.

[11]L. Civale, A. D. Marwick, T. K. Worthington, M. A. Kirk, J. R. Thompson, L. Krusin-Elbaum, Y. Sun, J. R. Clemm, and F. Holtzberg, Phys. Rev. Lett. **67**, 648 (1991).

[12]R. P. Vasquez, B. D. Hunt, and M. C. Foote, Appl. Phys. Lett. **53**, 2692 (1988).

[13]J. R. Thompson, Y. R. Sun, D. K. Christen, H. R. Kerchner, A. P. Malozemoff, L. Civale, A. D. Marwick, T. K. Worthington, L. Krusin-Elbaum, and F. Holtzberg, in *Physics and Materials Science of High Temperature Superconductors II*, Vol. 209 of *NATO Advanced Studies Institute, Series E*, edited by R. Kossowsky, B. Raveau, D. Wohlleven, and S. K. Patatis (Kentennity, Dordrecht, 1992).

[14]D. R. Nelson, Phys. Rev. Lett. **60**, 1973 (1988).

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

[16]A. Houghton, R. A. Pelcovitz, and A. Sudbo, Phys. Rev. B **40**, 6763 (1989).

[17]R. E. Hetzel, A. Sudbo, and D. A. Huse, Phys. Rev. Lett. **69**, 518 (1992).

[18]A. I. Larkin and Y. N. Ovchinnikov, J. Low Temp. Phys. **34**, 409 (1979).

[19]M. A. Moore, Phys. Rev. B **45**, 7336 (1992).

[20]M. V. Feigel'man and V. M. Vinokur, Phys. Rev. B **41**, 8986 (1990).

[21]E. Brezin, D. R. Nelson, and A. Thaiville, Phys. Rev. B **31**, 7124 (1985).

[22]Y. Imry and M. Wortis, Phys. Rev. B **19**, 3580 (1979).

[23]D. E. Farrell, J. P. Rice, and D. M. Ginsberg, Phys. Rev. Lett. **67**, 1165 (1991).

[24]H. Safar, P. L. Gammel, D. A. Huse, D. J. Bishop, J. P. Rice, and D. M. Ginsberg, Phys. Rev. Lett. **69**, 824 (1992).