Flux Creep and High-Field Critical Currents in Epitaxial Thin Films of $YBa_2Cu_3O_7$

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The critical current density $J_c(B,T)$ is measured to 15 T for c-axis-perpendicular expitaxial thin films of $YBa_2Cu_3O_7$ with $J_c(0,77\ \mathrm{K})$ of order 10^6 A/cm^2. Even in the least-favorable perpendicular orientation, critical currents can exceed 5×10^5 A/cm^2 at 20 K in 15-T fields. Thermally activated flux motion (flux creep) is prominently observed, and can in large part explain the magnitude and temperature and field dependence of the high-field critical currents.

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The current-carrying capacity of the high-T_c superconductor $YBa_2Cu_3O_7$ in a magnetic field is of both scientific and technological interest. Analysis of recent thin-film critical-current measurements[1-3] has not considered the potentially important role of thermally activated flux motion (flux creep).[4-7] In this Letter, we report new measurements on high-quality c-axis-perpendicular epitaxial $YBa_2Cu_3O_7$ films that have critical current density $J_c > 10^6$ A/cm^2 at 77 K.[8,9] We find that thermally activated flux motion plays a prominent role, and can explain in large part the magnitude and temperature and field dependence of the critical currents. Even for the least-favorable (perpendicular) field orientation, we find the high-field performance of $YBa_2Cu_3O_7$ at 20 K to be superior to that of Nb_3Sn at 4.2 K.

A standard expression[10] for the dissipative v associated with thermally activated flux creep is

$$v = v_0 \exp(-U_0/k_BT)\sinh(JBV_dL/k_BT)\,, \qquad (1)$$

where v_0 is the prefactor related to an attempt frequency, U_0 is the activation energy, J is the transport current density, V_d is the volume over which the driving force on the moving flux bundle is determined, and L is the typical distance moved. This can be recast as

$$v = v_0 \exp(-U_0/k_BT)\sinh[(U_0/k_BT)J/J_0]\,, \qquad (2)$$

where $J_0 = U_0/BV_dL$ is the current density that drives the activation energy to zero by "tipping the potential well." Tinham[5] suggests that lattice shear involving the slippage of a vortex past its neighbors sets the energy scale for the activation of creep. He considers vortex lattice rearrangements (square versus triangular) and argues that

$$U_0 \simeq 0.02(B_c^2/2\mu_0)V_a\,, \qquad (3)$$

where $B_c(T) \simeq B_c(0)(1-t^2)$ is the thermodynamic critical field and V_a is the volume of the minimum region triggering the activated event. (The reduced temperature t is T/T_c.) Tinkham further suggests that at large fields V_a consists of N_1 vortex areas $a_0^3 = \Phi_0/B = h/2eB$ times N_2 longitudinal coherence lengths $\xi(T) = \xi(0)(1-t)^{1/2}$, or

$$V_a \simeq N_1N_2\xi\Phi_0/B\,. \qquad (4)$$

Thus U_0 varies as $1/B$, as postulated by Yeshurun and Malozemoff.[4] Given any fixed voltage criterion $v_c \ll v_0$, Eq. (2) defines a J_c given by

$$J_c = J_0(k_BT/U_0)\sinh^{-1}\{\exp[U_0/k_BT - \ln(v_0/v_c)]\}\,. \qquad (5)$$

In what follows, we will explicitly demonstrate the activated nature of the dissipation in our thin films, and systematically compare the measured critical currents with (2)-(5).

Our films[8,9] are fabricated by coevaporation of Y, Cu, and BaF_2 in an O_2 atmosphere on room-temperature $SrTiO_3$ substrates, followed by an 800 °C anneal in wet oxygen and a 550 °C anneal in dry oxygen. They typically have a very high degree of epitaxial order, consisting almost entirely (> 90%) of crystals with the c-axis perpendicular to the substrate (copper oxide planes parallel to the substrate).[9] In order to measure large values of J_c, we pass current through a scratched microbridge $\simeq 100$ μm wide. If the current flow is macroscopically inhomogeneous, this will cause us to underestimate J_c. Low-resistance four-terminal contacts were made with indium pressed on 1-mm gold dots evaporated directly on the film surface. Films with thicknesses of 50, 100, and 200 nm have been studied from T_c to 20 K in magnetic fields up to 15 T. Thus critical currents of order 100 mA correspond to densities of order 10^6 A/cm^2.

Thermally activated behavior is prominently observed. Figure 1 shows the activated tail of the resistive transition of a 100-nm film in a 12-T perpendicular field. (We specify field orientations with respect to the substrate

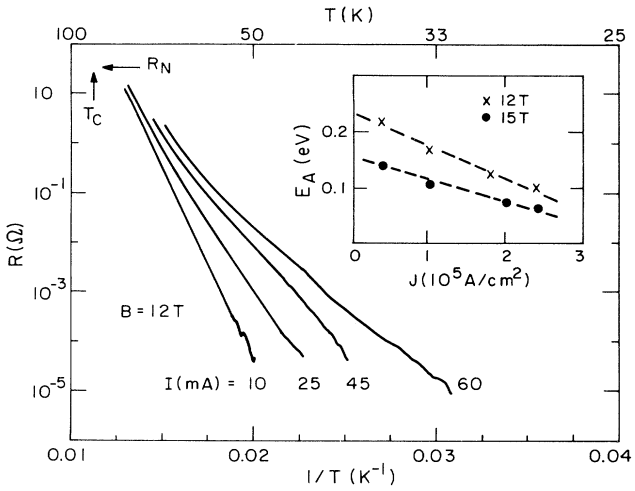

FIG. 1. Activated tail of resistive transitions at 12 T for various currents through a thin-film microbridge. Inset: Increasing current density reduces the measured activation energy E_A as described in the text.

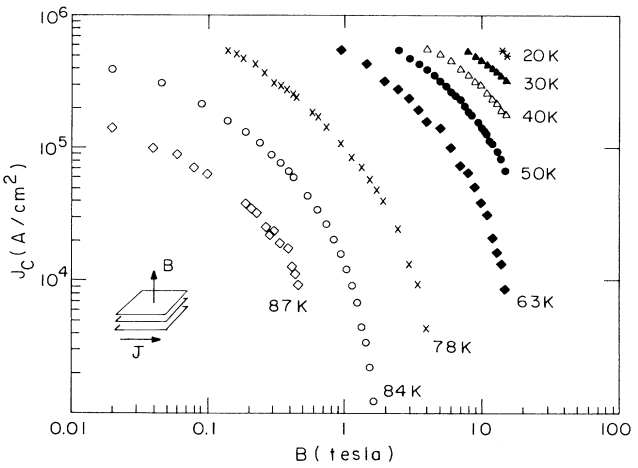

FIG. 2. Dependence of the critical current density on perpendicular magnetic field for thin-film microbridges at various temperatures.

and hence copper oxide planes; the field is always perpendicular to the direction of the current through the bridge.) In this range, Eq. (2) predicts roughly constant slopes corresponding to activation energies $E_A = U_0[1 - J/J_0]$ that decrease with increasing current density, as confirmed in the inset. Nonlinear curve fitting to a wider range of low-current data including the temperature dependence $U_0 = U_{00}(1-t^2)^2(1-t)^{-1/2}$ gives $U_{00}B = 2.2$, 2.2, 1.8, 2.1, and 2.0 eVT for fields of 15, 12, 4, 2.75, and 1.5 T, respectively. This demonstrates the distinctive $1/B$ dependence of the activation energy and corresponds to a plausible value of $N_1N_2\xi(0) \simeq 3$ nm [assuming Eq. (4) with $B_c(0) = 2.7$ T (Ref. 11)]. Thus nucleation of vortex motion apparently is initiated at a small number of vortices over a small fraction of their length. The $1/B$ dependence suggests the importance of interactions between adjacent vortices in preventing their motion.

The intercept $J_0 \simeq 4\times10^5$ A/cm^2 in the Fig. 1 inset corresponds to $V_dL \simeq 5\times10^{-31}$ m^4. Under the assumptions above, this exceeds V_a by a length $L_0 \simeq 1$ μm. V_dL determines the energy gain which biases the hopping rate in the driven direction. We suggest that it is likely to be determined by a small number of vortices moving one lattice spacing and relaxing the stress of a column of vortices behind them for approximately a penetration depth λ, which is a typical vortex interaction distance. Thus we expect the macroscopic Lorentz force $\mathbf{J}\times\mathbf{B}$ on a volume $V_d = a_0(N_3\lambda)(N_4\xi)$ to act for a distance $L = a_0$. This product differs from V_a by L_0 of a few penetration depths. Since $\lambda = 150$–200 nm for our $YBa_2Cu_3O_7$ films,[12] this interpretation is consistent with our experimental result.

The general features of Eq. (1) are also manifest directly in current-voltage characteristics. As the critical current approaches zero with increasing B, the I-v shape at low voltages changes from exponential to linear (with B-dependent slope $\ll R_N$), as predicted by (1). Measurements on a 200-nm-thick film at 15 T in the temperature range 25–50 K yield slopes $d(\ln v)/dI$ corresponding to $L_0 \equiv V_dL/V_a = 150$–$600$ nm, roughly consistent with the discussion above. However, systematic presentation of such results is beyond the scope of this Letter.

The central question that we now address is whether the observed flux-creep phenomena are actually sufficient to account for the high-field critical currents of our microbridges, or whether they only marginally affect the behavior near current levels set by other phenomena. In principle, flux creep implies the existence of nonzero dissipation under all conditions, but for practical purposes one can always define an appropriate critical level. In what follows we analyze the critical currents determined by a 1-μV criterion. Higher criteria give similar results, while substantially lower criteria are obscured by experimental noise.

Figure 2 shows $J_c(B,T)$ for a 200-nm film also in perpendicular-field orientation. Current densities of 5×10^5 A/cm^2 can be carried at 20 K up to 15 T. For comparison, typical Nb_3Sn filaments carry such current densities only up to about 8 T at 4.2 K (comparable T/T_c). The J_c then drops a further order of magnitude by 15 T.[13,14] Thus $YBa_2Cu_3O_7$ at the boiling point of liquid hydrogen can carry more current at high fields than Nb_3Sn at the boiling point of liquid helium.

By analogy with conventional superconductors, we examine the scaling behavior of the pinning force per unit volume $F_p = BJ_c$ with field and temperature.[15–17] Figure 3 shows that the data of Fig. 2 can indeed be scaled into a single curve by scaling factors $F_{p\max}(T)$ and $B^*(T)$. At low temperatures, the available data are sufficient to locate the peak, and thus provide estimates of B^* to values as large as 100 T. Figure 3 also shows the flux-

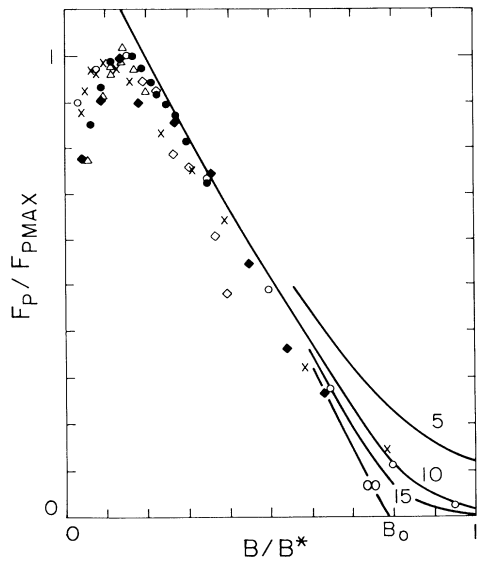

FIG. 3. Scaling of J_cB vs B for data of Fig. 2 (same symbols). The solid curves are the high-field predictions of the flux-creep theory, for various values of $\ln(v_0/v_c)$ (see text).

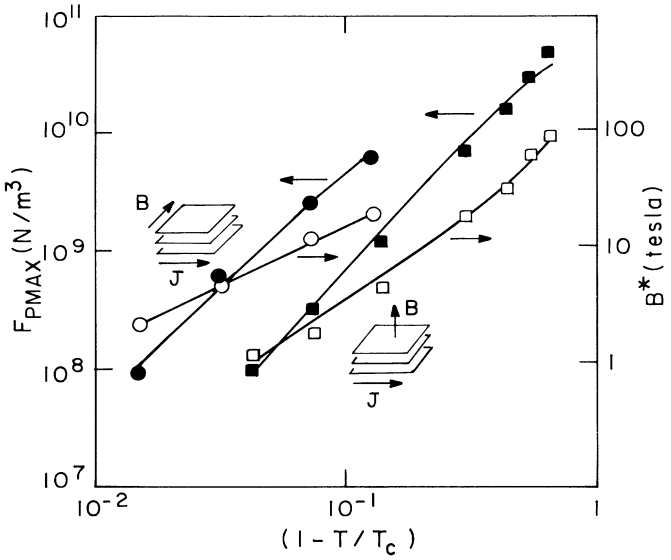

FIG. 4. Temperature dependence of the scaling parameters $F_{p\max}$ (solid symbols) and B^* (open symbols) for the perpendicular-field data of Fig. 3 (squares) and similar parallel-field data (circles). The solid (fitted) curves are Eqs. (8)–(11).

creep scaling behavior predicted by Eq. (5). The precise curvature at large fields depends on $\ln(v_0/v_c)$, but because $2\sinh^{-1}[\exp(x)]\simeq 1$ for large x, the J_cB curve is basically linear and extrapolates to a scaling field B_0, given by

$$B_0=\frac{U_0B/k_BT}{\ln(v_0/2v_c)}. \tag{6}$$

In the low-field direction, it extrapolates linearly toward

$$J_0B=U_0/V_dL\simeq 0.02(B_c^2/2\mu_0)V_a/V_dL. \tag{7}$$

We expect the behavior to change when the flux density is so low that vortices can be individually pinned, so that J_cB becomes proportional to B. Because the data show that this crossover occurs at small B, the observed $F_{p\max}$ should be smaller than the theoretical J_0B by only a minor factor ($\simeq 0.7$). Similarly, the fit indicates that B_0 corresponds to approximately 0.8 of our empirically defined B^*.

The temperature dependence of $F_{p\max}$ and B^* (Fig. 4) is shown for both perpendicular and parallel orientations, plotted assuming $T_c=91$ K, which corresponds to the midpoint of the zero-field transition. For the perpendicular orientation, the fitted solid lines are

$$B^*=1.3B_0=30(1-t^2)^2/t(1-t)^{1/2}\ \mathrm{T} \tag{8}$$

and

$$F_{p\max}=0.7J_0B$$
$$=6\times 10^{10}(1-t^2)^2(1-t)^{1/2}\ \mathrm{N/m^3}. \tag{9}$$

The fit to B^* agrees with Eq. (6) using the values of U_0 obtained from the activation analysis, provided that $\ln(v_0/2v_c)\simeq 12$ (a value consistent with the shape of Fig. 3). The fit to $F_{p\max}$ corresponds to Eq. (7) with $L_0\equiv V_dL/V_a\simeq(700\ \mathrm{nm})/(1-t)^{1/2}$, plausibly corresponding to a few times the (transverse) penetration depth.

Our parallel-field data scales to the same shape as Fig. 3, but the scaling parameters show a reduced temperature dependence. The fitted curves in Fig. 4 are given by

$$B^*=1.3B_0=75(1-t^2)^2/t(1-t)\ \mathrm{T} \tag{10}$$

and

$$F_{p\max}=0.7J_0B=1.2\times 10^{11}(1-t^2)^2\ \mathrm{N/m^3}. \tag{11}$$

The fit to $F_{p\max}$ implies a temperature independent $L_0\equiv V_dL/V_a\simeq 350$ nm. In this parallel-field orientation the Lorentz force is perpendicular to the film, and one would logically expect L_0 to be limited by the film thickness of 200 nm. Given the uncertainty in numerical parameters, we consider our result consistent with this idea. Within the flux-creep theory, the orientation dependence of B^* reflects the anisotropy of the longitudinal coherence length ξ, assuming that other factors such as N_1 and N_2 do not change. A careful study of the angular dependence at 77 K yields a B^* anisotropy of approximately 10, varying rapidly near parallel orientation. Because Fig. 4 shows a factor-of-2 smaller anisotropy at 78 K, our "parallel" alignment may not be exact. For this reason, we do not give significant weight to the unexplained reduced temperature dependence.

Given the uncertainty for numerical values of Eq. (3), we consider the self-consistency of our data strong evidence for the importance of thermally activated flux-

lattice motion in limiting the high-field current density of $YBa_2Cu_3O_7$. In conventional materials, these sort of data are frequently compared with theories of flux-lattice shear,[16] although controversy still abounds.[17] Because of the short coherence length and high temperature of $YBa_2Cu_3O_7$ thermally activated flux-creep processes should be expected to play a much larger role. Initially, we compared our results to the flux-lattice-shear theory of Kramer.[16] Because it is based on a static comparison between elastic distortions and Lorentz forces similar to the balance which occurs inside the exponents of the flux-creep theory, the agreement was also quite good, by order of magnitude. However, in this theory B^* must be interpreted as B_{c2}, where superconductivity completely disappears. This is not consistent with either the observed breadth of the resistive transitions,[5] or the unusual temperature dependence of B^* in the perpendicular orientation.

Finally, we have found that the high-field behavior can be explained without detailed knowledge of the microstructure of our films. Data from films of various thickness agree within factors of 3. More limited data from other high-current-density films made by a variety of techniques[1-3] are also generally consistent with our experimental results.

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