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Flux creep and the nature of a flux bundle in high-T_c thin films

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We describe measurements of the temperature and magnetic field dependence of the resistance and critical current in laser-deposited thin films of the high-temperature superconductor $YBa_2Cu_3O_{7-x}$. We identify two regimes: At low current densities and magnetic fields the sample is Ohmic. We attribute dissipation in this regime to activated flux motion of individual flux quanta in a vortex glass. In the high current and magnetic field regime we show that collective vortex motion is occurring, and we propose a model for the structure of a flux bundle.

Recently, considerable attention has focused on the mechanism of dissipation in the high-temperature oxide superconductors.[1] It has become clear that because of the extreme type-II nature of the materials, and the large thermal energies in the vicinity of the transition, thermally activated flux flow is extremely important,[2] and it has been proposed that the broadening of the resistive transition can be explained quantitatively by heavily damped flux flow.[3] Evidence for flux flow has recently been seen in *I-V* measurements on $YBa_2Cu_3O_{7-x}$ thin films and bulk material[4] and Bi-Sr-Ca-Cu-O crystals.[5] However, the nature of the moving flux bundles is still uncertain. For example, no flux-lattice images have been obtained at high temperatures,[6] and even at lower temperatures there is some evidence that the flux lattice is not static, but in fact is undergoing configurational changes which destroy short-range order.[7]

It is the purpose of this article to address in more detail the nature of the flux-pinning process in thin films of the high-temperature superconductor $YBa_2Cu_3O_{7-x}$ (Y-Ba-Cu-O) produced by pulsed laser deposition.[8] We will show that the resistive transition in a field, $R(B,T)$, may largely be explained within existing models of flux flow using the site-pinning energy proposed by Yeshurun and Malozemoff (Y-M) (Ref. 1) and expanded on by Tinkham.[3] However, we will show that dissipation is probably due to the motion of *individual* flux quanta in a vortex glass. In contrast, the critical current density $J_c(B,T)$ measured using a voltage criterion shows quite different behavior. In particular, we will demonstrate that flux motion may occur in *bundles* whose size is a strong function of temperature but *independent* of magnetic field, and we will propose a model for the structure of these bundles.

The analysis of the flux-flow measurements described in this article based on the fundamental equation of flux creep proposed by Anderson and Kim,[9,10]

$$v = v_0 \exp\left(-\frac{U_0}{k_B T}\right) \sinh\left(\frac{JBv_d l}{k_B T}\right). \qquad (1)$$

This describes voltage v developed due to thermally activated flow of flux bundles with a pinning energy U_0 acted on by a Lorentz force JBv_d, where v_d is the volume of the moving flux bundle, l is the average distance a flux bundle moves per jump, and the prefactor v_0 is related to the attempt frequency. It is straightforward to identify two regimes of behavior in Eq. (1): For low current densities or magnetic fields, we may linearize $\sinh(x) \approx x$, resulting in Ohmic behavior with approximately activated resistance.[11] In the second regime, observed at high magnetic fields and/or high current densities, the *I-V* characteristic is dominated by the sinh term, and the voltage rises exponentially with current.[4] We will explore both regimes in this paper.

The samples we have studied were produced superconducting *in situ* by pulsed laser deposition.[8] The films are epitaxial, with zero-point transition temperatures normally exceeding 88 K and critical current densities at 77 K larger than $10^6\ \mathrm{A\,cm^{-2}}$. The films are structurally single crystal and free from grain boundaries;[12] this is important since we require that the dissipation should be controlled by pinning rather than intergranular weak links which was a characteristic of early films and bulk material.[13,14] However, the films grown *in situ* have many atomic scale defects. The films are fabricated into constrictions of typical dimensions 10–20 μm by lithography and wet etching to allow high current densities to be employed without undue sample heating. Contacts are made using evaporated Ag followed by annealing in oxygen. The contact resistances are low ($<1\ \Omega$), but we restrict measuring currents to <20 mA. All magnetotransport measurements described have the magnetic field orthogonal to the current direction. The orientation quoted is the angle between the copper-oxide planes and the magnetic field.

The resistive transitions of two contrasting Y-Ba-Cu-O samples in a perpendicular magnetic field are shown in Fig. 1. The first sample has a normal-state resistivity at 100 K of 37 $\mu\Omega$ cm, a resistance ratio $R(300\ \mathrm{K})/R(100\ \mathrm{K}) = 2.86$, and a sharp transition with $R/R_N = 0.01$ at 89.5 K. However, the sample shown in Fig. 1(b), is not so good using these criteria, with a zero-point resistive transition in the absence of a magnetic field at 82 K. The parallel field configuration (not shown here) is far less severely broadened in a magnetic field (i.e., less than 3 K at 8 T for sample 1).

Tinkham has proposed a model for the dynamics of flux motion in high-T_c superconductors.[3] He considers heavily damped flux motion in a system with a site-pinning energy

$$U_0 = \gamma_{00}(1-t^2)^2(1-t)^{-1/2}/B$$

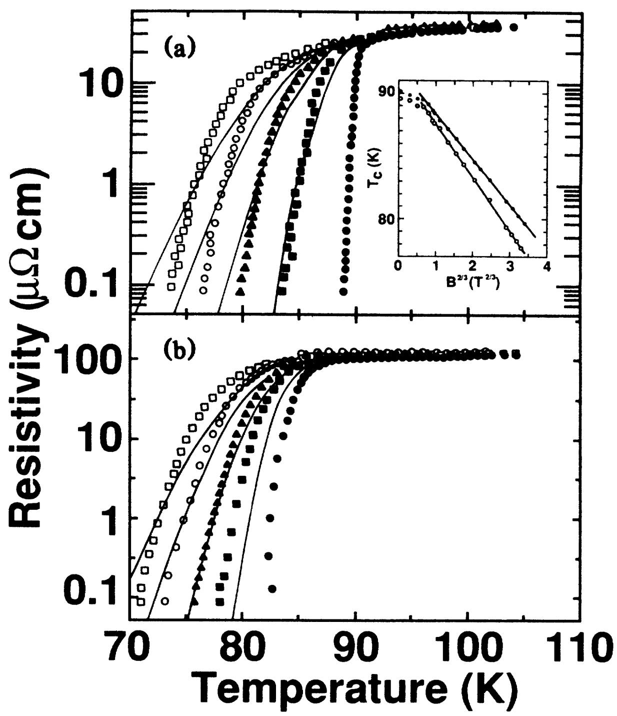

FIG. 1. The resistive transitions of two contrasting samples in magnetic fields of $B=0$, 2, 4, 6, and 8 T (from the right). Also shown by the solid lines are Tinkhams model for $R(T)$. Inset: the transition temperature measured at $0.1R_N$ (closed circles) and $0.03R_N$ (open circles) as a function of magnetic field.

and arrives at a broadened resistive transition given by

$$\frac{R}{R_N}\approx\left[I_0\left(\frac{2U_{00}}{k_BT_c}\right)\right]^{-2}, \qquad (2)$$

where I_0 is a modified Bessel function. We have attempted to fit this expression to our experimentally observed resistive transitions shown in Fig. 1. The fitting parameters are the normal-state resistivity R_N [we have chosen $R_N(T)$ to have a linear temperature dependence as extrapolated from above T_c], the transition temperature (which is chosen as the 10% point of the zero-field resistive transition), and the parameter γ_{00} which quantifies the zero-temperature site-pinning energy (which we vary to obtain a good fit to the $R=0.1R_N$ point on the characteristics). The results of the fit are indicated by the solid lines, where $\gamma_{00}=1.36$ and 1.95 eV T in (a) and (b), respectively. The corresponding unfluctuated critical current densities are $J_{c0}(0)=2\times10^6$ and 3×10^6 A cm^{-2}, respectively;[15] comparable to recent measurements.[4] Clearly, although the fits contain the essential aspects of the broadening, they do not fit the data very well. However, a more resistive test of the form of the activation function (rather than the dynamics of flux flow) is obtained from the expression for the broadening $\Delta T_c\approx B^{2/3}$ implied by the Y-M form of the activation energy. We show the shift in transition temperature measured at $0.1R_N$ and $0.03R_N$ (inset to Fig. 1). Excellent agreement to the $\frac{2}{3}$ power law is seen at magnetic fields above 1 T; however, at low magnetic fields, the dependence is faster. The conclusion we draw is that the Y-M form of the activation energy is accurately obeyed (above 1 T), but the model of the dynamics of flux flow is incomplete.

We turn now to the second part of the paper which describes the nonlinear transport properties of the Y-Ba-Cu-O thin films. At high magnetic fields and/or high current densities, Eq. (1) reduces to

$$v=\frac{v_0}{2}\exp\left[-\frac{U_0}{k_BT}+\frac{JBv_dl}{k_BT}\right]\approx\frac{v_0}{2}\exp\left[\frac{JBv_dl}{k_BT}\right], \qquad (3)$$

implying approximately exponential current-voltage characteristics, which we observe but are not reported here. We will be most concerned with the form of the critical current density $J_c(B,T)$. The operational definition for J_c we will use is the curent density at which a small, measurable voltage v_x is developed along the sample. Using this definition, Eq. (3) may be recast as

$$J_c\approx\left[\frac{k_BT}{v_dl}\right]\frac{1}{B}\ln\left[\frac{2v_x}{v_0}\right]. \qquad (4)$$

Therefore, the characteristic of flux flow in this regime is a reciprocal relationship between the critical current density defined using a voltage criterion, and the magnetic field B *if the bundle size and jump distance is independent of magnetic field*[16] (we take no account of the dynamics of *bundle* motion).

The sample described in Fig. 1(a) shows such behavior in the critical current characteristics (Fig. 2), where the voltage criterion is 3 μV developed along the length of the constriction (≈30 μm). Clearly, below 75 K, and in magnetic fields ranging from 2 to 9 T, we find $J_c\propto1/B$.[9] We can immediately draw some interesting conclusions from this behavior. First, dissipation is probably occurring through the motion of flux bundles, and *not* individual flux quanta which would show a different power-law characteristic.[17,18] This is an important conclusion since there has been some doubt whether a flux lattice exists in the high-temperature superconductors at elevated temperatures.[6] However, we remark that although it is clear that flux slip occurs in bundles, it does not imply that the flux structure is crystalline, merely that the motion is strongly correlated. We will return to the precise nature of a flux bundle shortly. Second, the straight-line charac-

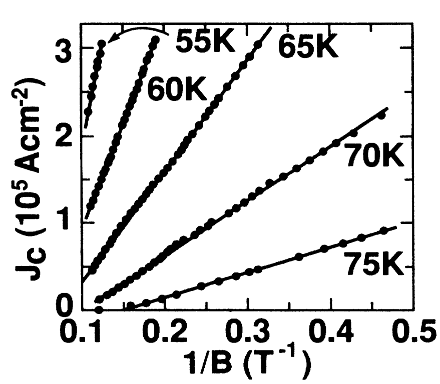

FIG. 2. The critical current density in a magnetic field.

teristic in Fig. 2 also indicates that the bundle size v_d and the jump distance L are essentially unaffected by magnetic field and are therefore not directly related to the flux-lattice spacing in this regime, which would be the case for the motion of individual fluxons (since the vortex spacing is field dependent).[4]

We turn now to the temperature dependence of the gradient G of the $1/B$ plots illustrated in Fig. 2. In Fig. 3 we have plotted G/k_BT for this sample and two others in the perpendicular field orientation. From Eq. (4) we obtain $G/k_BT = C(v_dL)^{-1}$, where $C = \ln(2v_x/v_0) \approx -10 \pm 5$.[10] Therefore, from Fig. 3 we may obtain *directly* the (bundle volume) (jump distance) product. Assuming the jump distance and bundle dimensions are all equal, we arrive at the characteristic length $(v_dL)^{1/4} \approx 50$ nm. Further, this length diverges at T_c, which is to be expected if it is related to either of the superconducting length scales $\lambda(T)$ or $\xi(T)$.

Finally, we will discuss in more detail the elementary flux structures implied by our data. The resistive transition data presented here and by other authors lends support to the activation energy

$$U_0 \approx \gamma_{00}(1-t^2)^2(1-t)^{-1/2}/B\,.$$

This has been justified on the basis of a constant *number* of vortices (independent of T and B) moving collectively.[3,4] This is rather remarkable since the vortex interaction energy is a strong function of temperature, and the vortex spacing depends on the magnetic field. However, a minor modification to this picture is that a single vortex is moving quasi-independently in the configurational energy space defined by its neighbors (some of which are presumably pinned). In this case, the energy of a vortex in an unfavorable energy configuration (i.e., the energy barrier) is related to the energy density of the vortex lattice multiplied by the volume occupied by a *single* vortex in the lattice, i.e., $U_0 = \alpha(B_c^2/2\mu_0)[(\phi_0/B)\xi]$, where α is a numerical constant. Using the activation energy $\gamma_{00}B = 1.75$ eV at 1 T, and using published values for the other parameters, we find $\alpha \approx 10^{-2}$, which is reasonable for a local configuration change.[3] However, we still expect the picture of vortices interacting predominantly with each other rather than sample defects should fail when the average fluxon spacing exceeds the pinning-site spacing, whereupon the vortex-defect interaction rather than the vortex-vortex interaction should dominate. This explains the breakdown of $\Delta T_c \propto B^{2/3}$ at low magnetic fields, and, more generally, why no universal broadening is seen. In our sample, the breakdown of the power-law behavior occurs at $B \approx 1$ T, implying a typical defect spacing of 45 nm.

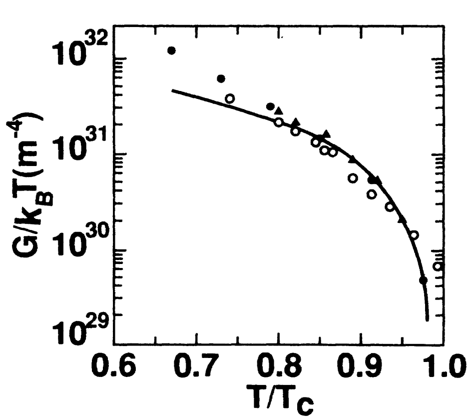

FIG. 3. Symbols: the gradient of the characteristics in Fig. 2 divided by the temperature. Line: the results of the semiempirical model of a flux bundle described in the text.

The preceding discussion implies that there is no short-range order in the fluxon structure in the vicinity of the transition, which is undergoing continual, activated configurational changes and moving distances comparable to the average vortex spacing. Clearly this provides evidence for the vortex liquid phase which is receiving considerable current attention.[19–21] However, the high-field and current nonlinear characteristics require a rather different interpretation. The distinctive $J_c \propto 1/B$ dependence implies correlated motion of a bundle of vortices whose size and jump distance v_dl are independent of magnetic field (and hence vortex spacing). Furthermore, the strong temperature dependence of v_dl argues against the flux bundle size being controlled directly by the defect distribution within the material. Therefore, we propose that the bundle size is determined predominantly by the vortex-vortex energy interaction length: the penetration depth $\lambda(T)$. Hence, a plausible explanation of the vortex dynamics in the current-driven, collective regime is that a collection of vortices, whose area is related to λ^2 and whose length is a few $\xi(T)$, is moving a distance also related to $\lambda(T)$ at each activated event. These arguments imply

$$v_dl \propto \lambda(0)^3\xi(0)(1-t^4)^{-3/2}(1-t)^{-1/2}$$

(we could also choose a vortex length determined by the penetration depth without significantly affecting the temperature dependence). We have plotted the form of this expression in Fig. 3, giving a reasonable fit to the observed temperature dependence. Furthermore, assuming the length of the bundle is $\xi(T)$ (an underestimate), and the jump distance is equal to the bundle size, we conclude a bundle size at $t = 0.9$, and 5 T of $\delta \approx 100$ nm consisting of ≈ 25 fluxons which is comparable to that observed in conventional superconductors.[10]

The picture of two differing length scales for vortex motion in the low-current and high-current regimes is not unreasonable. At low fields and currents, the flux lattice is so soft, and thermal energies are so high that individual vortex motions are to be expected. However, at high currents, when the dissipation is controlled by the driving force on the fluxons, the Lorentz force per resolved unit length on each flux quantum is the same. Therefore, in the absence of some extrinsic pinning mechanisms, collective motion is to be expected (although thermal configurational changes will still destroy short-range order) and a second length scale *must* be invoked to keep the dissipation finite.

In summary, we have delineated a number of regimes in the dissipative behavior of the high-temperature superconductor $YBa_2Cu_3O_{7-x}$. Ohmic behavior is observed at low current densities and high temperatures. This can be explained in terms of flux creep of individual flux quanta in a vortex glass or liquid. This description fails at low magnetic field when the vortex-defect interaction exceeds the vortex-vortex interaction. At high current densities and magnetic fields, a second process becomes dominant. In particular, collective, Lorentz driven motion of bundles of flux lines takes place. We propose that the bundle structure is determined primarily by the magnetic penetration depth in the material.

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[15]J_{c0} is the theoretical depairing current density, that is the critical current density in the absence of thermal fluctuations and vortices, as discussed in Ref. 3.

[16]Strictly, only the product $J_c v_d$ needs to be field independent, but it is unlikely that the separate dependences cancel out.

[17]In the case of individual fluxons, we expect the bundle size to be related either to the vortex spacing, which is clearly magnetic field dependent, or the spatial defect distribution, which is presumably temperature independent. Neither of these possibilities fits the data.

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