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Flux creep and critical-current anisotropy in $Bi_2Sr_2CaCu_2O_{8+\delta}$

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We have studied magnetic relaxation, temperature dependence of the critical currents, magnetic irreversibility, and lower critical fields in single crystals of the high-T_c superconductor $Bi_2Sr_2CaCu_2O_{8+\delta}$ (idealized composition) through magnetization measurements. The results are analyzed within the framework of the critical state and a thermally activated flux-creep model. The results indicate three different critical currents and flux-creep rates depending on the relative orientation of the current, the field, and the Cu-O planes. The flux-creep rate for flux lines moving perpendicular to the planes, was found to be unobservably low. Similarly, the critical-current density perpendicular to the planes was very low. However the critical-current density in the plane and the flux-pinning well depth U_0, for fluxon motion parallel to the planes, are similar to the values found in Y-Ba-Cu-O. Implications of insulating layers between the Cu-O planes are discussed.

Recent studies[1-3] of the time-dependent decay of zero-field-cooled (ZFC) magnetization and magnetic irreversibility in Y-Ba-Cu-O have shown the importance of classical flux-pinning ideas[4-6] in understanding the magnetic behavior of high-T_c superconductors.

In this study of flux-creep-related phenomena in the high-T_c superconductor $Bi_2Sr_2CaCu_2O_{8+\delta}$ (idealized composition), we measure magnetic relaxation, critical-current densities, magnetic irreversibility, and lower critical fields for two cases: field parallel to the Cu-O planes (**H** ∥) and field perpendicular to the Cu-O planes (**H** ⊥).

The pinning well depth U_0, one of the parameters which characterizes the flux-creep process, was found from the magnetic relaxation as well as the temperature dependence of the critical-current density J_c. The values obtained by the two methods are in agreement with each other. Depending on whether the fluxons are moving parallel or perpendicular to the planes, one expects quite different U_0's. For the **H** ⊥ case, fluxons always move parallel to the planes. For **H** ∥, fluxons can move both perpendicular as well as parallel to the planes. Interestingly, U_0 turns out to be the same for both field directions within experimental error. This suggests that flux motion perpendicular to the planes is negligible. This suppression of flux motion perpendicular to the planes may be the result of strong pinning by insulating layers between the Cu-O planes.

The Bean critical-state model[7] is used to determine J_c from the volume magnetization M. For the **H** ⊥ case, J_c is isotropic in the plane normal to the field, as are the sample dimensions—allowing the use of the customary isotropic-cylinder approximation: $J_c = 30M_{res}/R$ (where R is the radius of the cylinder, M_{res} is the residual magnetization which remains after the field has been raised to a high value and brought back to zero, and J_c is in A/cm^2). For the H ∥ case the cross section of the sample is a long rectangle and the current flow is both parallel as well as perpendicular to the planes and therefore not isotropic. We show that for our crystals it is appropriate to approximate the sample as an infinite slab, but one in which the large faces are neglected and only the ends are considered.

The single crystals, whose preparation has been described previously,[8] were platelets (with c axis along the small dimension) with average dimensions $0.8\times0.8\times0.07$ mm^3. The samples had a T_c of 92 K, and showed essentially 100% shielding and a large Meissner fraction (70% for **H** ⊥ Cu-O planes and 63% for **H** ∥ planes, with $H=2.5$ Oe). Measurements were made using a superconducting quantum interference device (SQUID) susceptometer (SHE model 905). Time-dependent magnetization measurements were made by stepping up the field from zero to the desired value. The magnetic irreversibility temperature T_{irr}, was determined from the point of departure between the magnetization curve obtained by cooling a sample in the field (FC), and the curve obtained by warming up the sample for which the field had been switched on after it had been cooled in zero field (ZFC).

We begin our analysis with a discussion of the relation between J_c and M. The Bean-model approach assumes that the flux-density gradient is normal to the surface and has a magnitude of $4\pi J_c/c$. This allows determination of the flux-density profile within the sample, from which the magnetization can be found. For the **H** ⊥ case, we assume the sample to be a cylinder whose radius is equal to half the transverse dimension: $R \simeq w/2 \simeq h/2$ (see inset of Fig. 2 for definitions of w and h). The field for complete flux penetration is $H^* = 4\pi J_c R/c$. For $H_{c1} \ll H < H^*$, M is given by[7]

$$4\pi M = -H + \frac{H^2}{H^*} - \frac{H^3}{3H^{*2}}\,. \tag{1}$$

For $H > H^*$, M is given by

$$4\pi M = -\frac{H^*}{3}\,, \tag{2}$$

which is also the correct expression for the magnitude of the residual magnetization $M_{\rm res}$, which remains in the sample after the field has been increased above $2H^*$ and then brought back to zero. The critical-current density plotted in Fig. 1 was calculated from $M_{\rm res}$ in zero field.

For the **H** $\parallel$ case, flux penetration occurs as indicated by the shaded portion of the inset of Fig. 2. As will be justified by our data analysis, we assume from the outset that $1/(tJ_{cy})$ and $1/(wJ_{cx})$ are quite different [calling the larger one $1/(sJ_c)$] so that flux penetration from the ends is not comparable to that from the sides, allowing one of them to be neglected. A straightforward extension of the Bean model shows that the field for complete penetration is given by $H^*=2\pi sJ_c/c$. For partial penetration, the magnetization can be easily calculated by noting that (A) there is complete flux exclusion from the inner unshaded rectangle (i.e., $4\pi M=-H$) and (B) the average magnetization in the shaded portion is half that for complete exclusion (i.e., $4\pi M=-H/2$). Thus for $H_{c1}\ll H<H^*$

$$4\pi M\simeq -H+\frac{H^2}{2H^*}=-H+\frac{H^2c}{4\pi sJ_c}\,. \tag{3}$$

Similarly, for $H>H^*$

$$4\pi M\simeq =\frac{H^*}{2}=-\frac{\pi sJ_c}{c}\,, \tag{4}$$

which also gives the magnitude of the residual magnetization as for the **H** $\perp$ case. Once again we use the value of $M_{\rm res}$ in zero field to calculate J_c (plotted in Fig. 2). We measured the magnetization of samples with different w/t ratios ($\sim$8–23) and found that Eqs. (3) and (4) give approximately the same value of J_c in each case if w is substituted for s, showing that $1/(tJ_{cy})$ is much smaller than $1/(wJ_{cx})$. The fact that flux penetration from the faces (i.e., along x) was negligible was also demonstrated in another way. Equation (4) implies that $M\propto s$. If flux penetration is from the large faces (x direction) then $s=t$, and breaking a crystal into two pieces without altering the thickness should not affect M. Therefore, the two pieces measured together will have the same moment m, as the original crystal. However, if flux penetration is from the edges (y direction), then breaking the crystal in the manner described above would reduce the combined moment of the pieces to 50% of the original moment (if the two pieces were equal and regular, and all other assumptions were satisfied). What we found was that the combined moment of the two pieces was 65% of the original moment. This convincingly demonstrates that $J_{cy}\gg(w/t)J_{cx}$ and allows us to put a lower bound on J_{cy} of 2×10^6 A/cm^2 at 7 K. This influence of flux-creep anisotropy on the interpretation of specimen magnetization in terms of various J_c's for the **H** $\parallel$ case (which can lead to errors as large as an order of magnitude) has not been fully addressed in the literature.

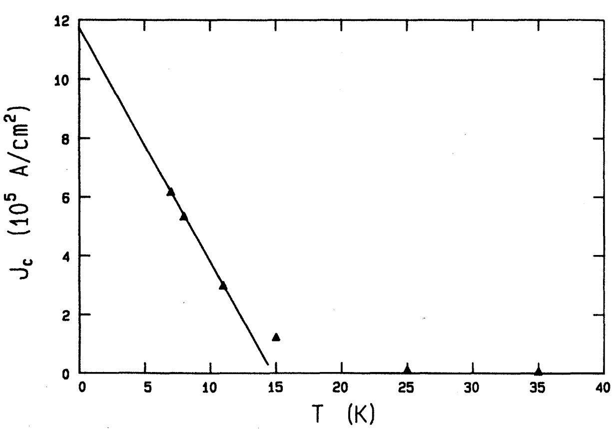

FIG. 1. Temperature dependence of J_c for **H** $\perp$ Cu-O planes.

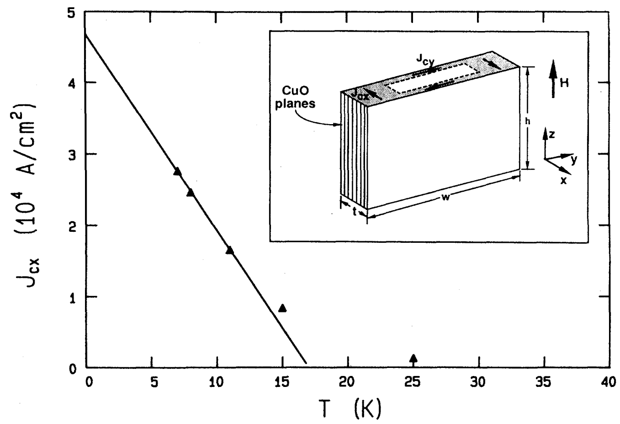

FIG. 2. Temperature dependence of J_{cx} for **H** $\parallel$ Cu-O planes. Inset: Definition of relevant directions for current flow and flux penetration. Shaded areas represent the area of flux penetration.

We analyze the temperature dependence of J_c in the framework of a thermally activated flux-creep model.[4-6] Thermal activation increases the flux-creep rate thereby reducing the observed critical-current density. This effective J_c is given by[2,6]

$$J_c=J_{c0}\left[1-\frac{k_BT}{U_0}\ln\left(\frac{t}{t_0}\right)\right], \tag{5}$$

where J_{c0} is the value of J_c in the absence of thermal activation of flux lines, U_0 is the depth of the potential wells which trap the flux vortices, t is a characteristic time for an experiment (the time since the field was changed and the critical state established), and t_0 is a constant which depends on parameters such as the fluxon oscillation frequency, the average hopping distance of the fluxons, etc.[6] The value of the logarithm for typical parameters[6] is 30, which roughly corresponds to the value of t_0 found in Y-Ba-Cu-O by Malozemoff *et al.*[9] from the frequency dependence of ac susceptibility. We will use this value for $\ln(t/t_0)$ noting that for the logarithm to change by 50%, its argument is required to change by 10^7. Since our calculations are correct only within a factor of 2 or 3, the exact choice has no significant effect on the results.

Equation (5) has two consequences. The temperature dependence is much stronger than the time dependence. Therefore, to a first approximation J_c is time independent and linear in T at low temperatures (where J_{c0} is independent of T), as is evident from Figs. (1) and (2). Thus, extrapolation of the data to $T=0$ K gives J_{c0} for each field direction: $J_{c0x}=4.7\times10^4$ A/cm^2 (**H** ∥) and $J_{c0}=1.2\times10^6$ A/cm^2 (**H** ⊥). We return to a discussion of these current densities later. From Eq. (5) and the observed low-temperature slopes in J_c vs T, we find the following U_0's: $U_0=0.038$ eV (**H** ⊥) and $U_0=0.044$ eV (**H** ∥). The former is for flux lines perpendicular to the planes, and the latter is for flux lines parallel to the planes. The observation that the two U_0's are about the same, supports the conclusion made earlier that flux creep is parallel to the planes, in the **H** ∥ case as well as the **H** ⊥ case.

The second consequence of Eq. (5) is that at a fixed temperature the magnetization in a specimen will decay logarithmically with time. This behavior can be seen in Fig. 3 which shows data for the **H** ⊥ case at 10 K. Inserting the temperature-dependent form of J_c into Eqs. (1) and (3) gives the logarithmic decay rates for the two field directions:

$$\frac{d(4\pi M)}{d\ln t}=\begin{cases}(H^2/H^{*2}-2H^3/3H^{*3})(4\pi J_{c0}R/c)k_BT/U_0, & \mathbf{H}\perp, \quad (6)\\ (H^2cJ_{c0x}/4\pi wJ_{cx}^2)k_BT/U_0, & \mathbf{H}\parallel. \quad (7)\end{cases}$$

This allows an independent determination of U_0 for each direction—a determination which does not involve the unknown parameter t_0. The decay rates $dM/d\ln(t)$ measured at several different temperatures and fields were qualitatively well described by Eqs. (6) and (7). At the lowest temperatures the decay rates were found to be proportional to temperature, as would be expected from the above equations since the implicit temperature dependence of the other terms becomes negligible. The U_0's obtained from these low-temperature decay rates are $U_0=0.012$ eV (**H** ⊥) and $U_0=0.048$ eV (**H** ∥). For the **H** ⊥ case, the values differ by a factor of 3. In view of the approximations made in the analysis, the two values are not inconsistent. For the **H** ∥ case, the agreement is within 10%, perhaps reflecting the fact that an infinite slab represents this situation more accurately than a cylinder does the other. The fair agreement between the U_0's obtained by the two methods shows that the value of 30 chosen for $\ln(t/t_0)$ is reasonable. The U_0's in this material are similar to $U_0^\perp$ in Y-Ba-Cu-O (Ref. 2).

As shown by Yeshurun and Malozemoff,[1] the temperature-dependent form for J_c [Eq. (5)] leads to the existence of a "quasi-de Almeida–Thouless" (AT), or irreversibility line ($1-T_{irr}/T_c\propto H^{2/3}$). Since in this model a nonvanishing J_c gives rise to magnetic hysteresis, the temperature T_{irr} above which the magnetization is reversible can be identified with the temperature above which J_c is negligible, i.e., when $kT_{irr}\ln(t/t_0)/U_0\simeq1$. Based on this, the authors of Ref. 1 have predicted a scaling relation (valid near T_c) between T_{irr} and the applied field H:

$$1-t\simeq\left[\frac{8\pi f^2k_BT_c\ln(t/t_0)}{2.56H_c^2(0)\Phi_0\xi_0}\right]^{2/3}H^{2/3}, \tag{8}$$

where $t=T_{irr}/T_c$ is the reduced irreversibility temperature, ξ_0 is the coherence length in the direction of the applied field, $H_c(0)$ is the thermodynamic critical field at zero temperature, and f describes the value of the flux lattice parameter a_0 at which a crossover to collective pinning effects occurs, $a_0=f\xi$.[1] We assume that f has the same value as in Y-Ba-Cu-O ($f\approx6$).[1] From the values

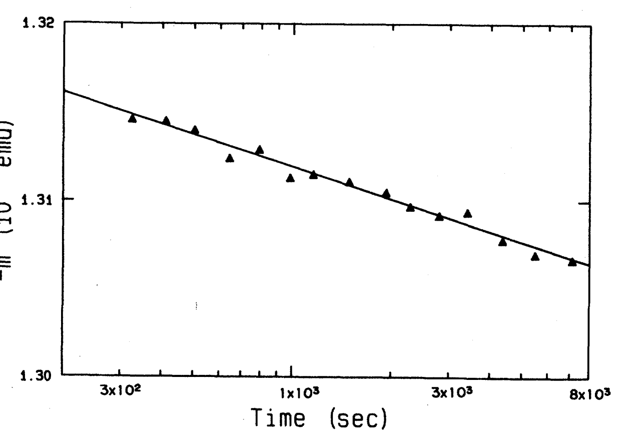

FIG. 3. Time logarithmic decay of the magnetic moment m, at $T=10$ K for **H** ⊥ Cu-O planes. The field was stepped up to 120 Oe after zero-field cooling the specimen.

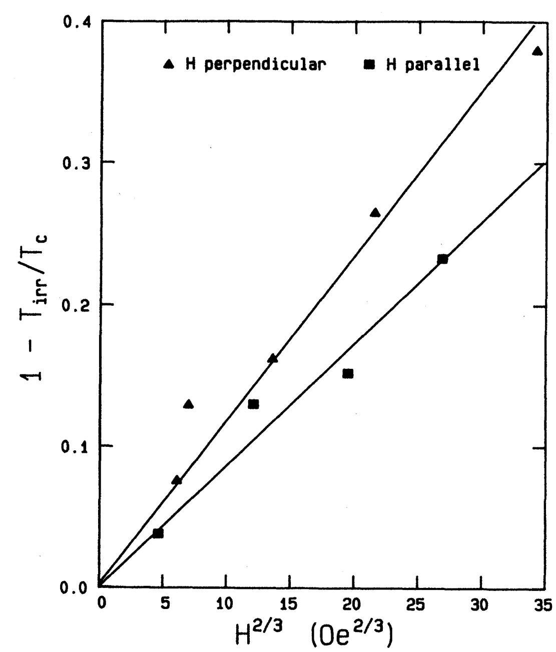

FIG. 4. Irreversibility lines for **H** ⊥ and **H** ∥ to the Cu-O planes.

of H_{c2} measured by Palstra *et al.*,[10] we use the anisotropic Ginzburg-Landau theory[11,12] to find $H_c(0)=3090$ Oe and the coherence lengths $\xi_\perp=1.6$ Å and $\xi_\parallel=38$ Å (perpendicular and parallel to the planes, respectively). Thus, for the $\mathbf{H}\perp$ case we predict $1-t=(0.03\ \mathrm{Oe}^{-2/3})H^{2/3}$. The measured irreversibility line, shown in Fig. 4, is $1-t=(0.011\ \mathrm{Oe}^{-2/3})H^{2/3}$. This agreement within a factor of 3 is not bad considering the many approximations. For the $\mathbf{H}\parallel$ case, the expected behavior from Eq. (8) is $1-t=(3\times10^{-3}\ \mathrm{Oe}^{-2/3})H^{2/3}$. The measured curve (Fig. 4) is $1-t=(8.4\times10^{-3}\ \mathrm{Oe}^{-2/3})H^{2/3}$, again with agreement within a factor of 3. We also find the relation $M_{\mathrm{rem}}=M_{\mathrm{FC}}-M_{\mathrm{ZFC}}$ to hold closely, where the remanent magnetization M_{rem} is that obtained by cooling the sample in a magnetic field, reducing the field to zero, and warming the sample in zero field. The observed equality between the difference $M_{\mathrm{FC}}-M_{\mathrm{ZFC}}$ and M_{rem} supports the flux-pinning picture.[13]

Estimates of the intrinsic J_c can be obtained by measuring the lower critical field H_{c1} and using the approximation[12]

$$\frac{H_{c1}}{\lambda}\simeq\frac{dH}{dx}=\frac{4\pi J_c}{c}\,. \tag{9}$$

Our measured values of H_{c1}($H_{c1}^{\perp}=175$ Oe and $H_{c1}^{\parallel}=9$ Oe) are consistent with earlier measurements on similar crystals by Lin *et al.*[8] We ascribe differences in our values and those reported earlier to uncertainties in the demagnetization factor and in the determination of the field at which deviation from linearity occurs in the M vs H curves. This last source of error is less serious for Bi-Sr-Ca-Cu-O than it is for some of the other high-temperature superconductors because of the relatively lower flux trapping (implied by the higher Meissner fractions) in this material. We use the anisotropic Ginzburg-Landau theory,[11] following the procedure of Umezawa *et al.*,[12] to calculate the penetration depths $\lambda^{\perp}=2000$ Å and $\lambda_y^{\parallel}=70000$ Å from H_{c1} and H_{c2} data, where the subscripts refer to the direction of field penetration and the superscripts refer to the direction of the applied field. The J_c's calculated from Eq. (9) are $J_c^{\perp}=7\times10^6$ A/cm^2 and $J_{cx}^{\parallel}=1\times10^4$ A/cm^2 which, considering the uncertainty in the H_{c2}'s and H_{c1}'s, are similar to the J_{c0}'s calculated from magnetization.

A number of interesting conclusions can be drawn from these results. First of all, the value of J_c in the plane for this material is not too different from that in Y-Ba-Cu-O—both are in the range 10^6–10^7 A/cm^2. The pinning well depth U_0 for fluxon motion parallel to the planes is also about the same (0.04 and 0.02 eV, respectively). On the other hand, the value U_0 for fluxon motion perpendicular to the planes is huge (0.06 eV) in Y-Ba-Cu-O and practically immeasurable in our crystals of Bi-Sr-Ca-Cu-O—flux creep in this direction was unobservable. These observations seem to support the notion that the Cu-O planes are the basic conduction units in the Cu-O based superconductors and, especially in Bi-Sr-Ca-Cu-O, that the conducting planes are separated by insulating layers. This layered structure in Bi-Sr-Ca-Cu-O, which has been invoked recently to explain anisotropies in other properties,[14] may be the reason for the excessively large U_0: The insulating layers may serve as very effective pinning sites (the layer thickness is larger than ξ), so that flux vortices are confined within them for the $\mathbf{H}\parallel$ case. At present we are growing very thin crystals ($w/t\sim1000$) to facilitate direct measurement of this large U_0 for interplanar fluxon motion and the corresponding J_c.

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