AIM

Association for Information and Image Management

1100 Wayne Avenue, Suite 1100
Silver Spring, Maryland 20910
301/587-8202

Centimeter

Inches

MANUFACTURED TO AIM STANDARDS
BY APPLIED IMAGE, INC.

Conf-940121--4

Effects of interlayer coupling on the magnetic and transport properties of superconducting multilayers and high-temperature superconductors*

K.E. Gray, J.D. Hettinger, and D.H. Kim

*Materials Science Division
Argonne National Laboratory, Argonne, IL 60439*

The submitted manuscript has been authored by a contractor of the U.S. Government under contract No. W-31-109-ENG-38. Accordingly, the U.S. Government retains a nonexclusive, royalty-free license to publish or reproduce the published form of this contribution, or allow others to do so, for U.S. Government purposes.

INVITED: SPIE Conference Proceedings on Superconducting Superlattices and Multilayers, Los Angeles, CA, January 22-28, 1994, edited by Ivan Bozovic

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

*Work supported by the U.S. Department of Energy, Division of Basic Energy Sciences-Materials Sciences under contract #W-31-109-ENG-38 and the National Science Foundation Office of Science and Technology Centers under contract DMR #91-20000.

MASTER

Effects of interlayer coupling on the magnetic and transport properties of superconducting multilayers and high-temperature superconductors*

K.E. Gray, J.D. Hettinger and D.H. Kim

Materials Sciences Division, Argonne National Laboratory, Argonne, Illinois, 60439, USA

ABSTRACT

The effect of interlayer coupling on the transport properties and dissipation in a magnetic field is reviewed for superconducting multilayers including highly-anisotropic high-temperature superconductors (HTS). For the applied field *parallel* to the superconducting layers the absence of any Lorentz-force dependence of the dissipation leads to an explanation other than flux motion. This is consistent with a Lorentz-force-independent, intergranular Josephson junction dissipation which dominates flux motion because of the strong, intrinsic pinning of the insulating regions between layers. However, in superconducting multilayers the dissipation is seen to cross over from phase slips at Josephson junctions to depinning of vortices from the external field at high fields and temperatures. For fields perpendicular to the superconducting layers the much greater resistive broadening in HTS is due to dissipation by thermally-activated flux motion, consistent with a lack of intrinsic pinning. We show experimental evidence that the associated flux motion occurs as a result of a crossover from three dimensional (3D) vortex *lines* to 2D independent *pancake-like* vortices, residing in the Cu-O layers. This 3D to 2D crossover occurs after $k_B T$ exceeds the Josephson coupling energy.

1. INTRODUCTION

The purpose of this paper is to review the effects of interlayer coupling on the transport properties and dissipation in a magnetic field of superconducting multilayers including the highly-anisotropic high-temperature superconductors (HTS). Experimental results from both HTS and low-temperature superconductors (LTS) will be used to illustrate suggested models for the dissipation mechanisms. An important example in HTS is the unusual field-induced broadening of resistivity transitions, $\rho(T,H)$, which places limitations on the value of the materials for moderate to high-field applications and has led to theories of thermally-activated magnetic flux creep¹, flux-lattice melting² or a vortex-glass transition³.

For fields parallel to the superconducting layers, e.g., the Cu-O, ab-planes in HTS, (H||ab), the broadening is smaller, and the absence of any measurable Lorentz-force dependence⁴⁻⁶ of the dissipation leads to an explanation other than motion of vortices from the external field⁶, e.g., Josephson coupling. Lorentz-force independence is also found⁶ in the low-temperature critical current density, J_c , and the current-voltage characteristics, $I(V)$, for c-axis-oriented films of $Tl_2Ba_2CaCu_2O_x$, and strong intrinsic pinning has been demonstrated in J_c measurements on multilayers⁷ of the LTS superconductor NbN with insulating AlN. However, in these LTS multilayers the dissipation is seen to cross over from phase slips at intergranular Josephson junctions in the NbN layers at low fields to depinning of external-field vortices at high fields and temperatures: these phase slips do not require the existence of an external field.

For fields perpendicular to the superconducting layers, e.g., parallel to the c-axis for HTS (H||c), the lack of intrinsic pinning implies that the much greater broadening is due to thermally-activated flux motion. For the *highly-anisotropic* HTS, experimental evidence indicates that the flux motion associated with this broadening occurs⁸ as a result of a dimensional crossover from three dimensional (3D) vortex *lines* to 2D independent *pancake-like* vortices⁹, which reside in well-coupled adjacent Cu-O bi- or tri-multilayers. This decoupling of vortices renders the sparse pinning sites in individual Cu-O multilayers to be much less effective. The 3D to 2D crossover is shown⁸ to occur when $k_B T$ exceeds the Josephson coupling energy of these Cu-O multilayers. For the less anisotropic $YBa_2Cu_3O_7$, this dimensional crossover does not occur much *below* H_{c2} , presumably because the conducting Cu-O chains short circuit the Josephson

Fig. 1. Schematic of LTS multilayer transport measurement in a parallel field.

interlayer coupling. The resulting 3D dissipation may be explained by thermally-activated magnetic flux creep¹ or a vortex-glass transition³. These results explain why YBa₂Cu₃O₇ has the best properties in a field and also shows that strong interlayer coupling is a key to finding good alternatives.

2. FIELD PARALLEL TO LAYERS

2.1. Crossover in NbN/AlN multilayers

Studies related to the intrinsic pinning of the Cu-O layers for fields parallel to them have been done in sputtered artificial multilayers of the LTS superconductor NbN with insulating AlN. The fabrication and characterization is discussed elsewhere⁷. The multilayer structure and measurement geometry are shown in Fig. 1, which also indicates that the magnetic vortices, which are shown projected onto the front side, lie in the strongly pinning AlN insulating layers (thickness of 2 nm). For the various NbN layer thicknesses (shown in Fig. 2) in which T_c varied between 10 and 15 K, the field-dependent J_c at a temperature of 4.2 K is plotted against $1-b$, where $b \equiv H/H_{c2}$. The high-field behavior is always a good fit to the $(1-b)^2$ dependence of vortex depinning from an ideal planar vacuum interface¹⁰, but there is a sharp crossover at a reduced field which varies with NbN thickness. Further experiments, with the field rotated parallel to the transport current (see Fig. 1), tested the Lorentz-force dependence. As opposed to the high-field behavior, J_c was completely independent of the macroscopic Lorentz force at fields below the crossover. Thus, the low-field J_c was unlikely to be due to motion of external field vortices, since it should then continue to extrapolate with the high field $(1-b)^2$ behavior also found¹⁰ in thick films of NbN.

2.2. Intergranular Josephson junction model

It was recognized that these results could be explained in terms of Josephson coupling between the columnar grains in NbN sputtered films (10-100 nm diameter and surrounded¹¹ by thin insulating boundaries). In these junctions, dissipation occurs by phase slips which have¹² a weaker field dependence than $(1-b)^2$ and no macroscopic Lorentz-force dependence to satisfy the low-field results shown in Fig. 2. In zero field, the Josephson critical current, I_{cj} , depends on the product of the superconducting order parameters, ψ_a and ψ_b , on each side of the junction¹³. Applied fields will be quite uniform for $\mu_0 H > 0.2$ T, because then the penetration length, λ , is $\gg$ the flux line spacing, but depairing by the field will reduce ψ : Abrikosov's solution¹⁴ of the vortex lattice determines the *spatial average* of ψ from $\langle \psi^2 \rangle = \psi_0^2 (1-b)$, where ψ_0 is the zero-field value. Thus to first order, $I_{cj} \propto \psi_a \psi_b \propto (1-b)$, which is a

Fig.2. Normalized critical current density for various multilayers of NbN/AlN vs. reduced field.

weaker field dependence than the $(1-b)^2$ for pinning. Finally, phase slippage occurs in zero field and does not require flux motion of external-field vortices, so it is independent of the macroscopic Lorentz force.

The data of Fig. 2 is interpreted as a competition between depinning of the external field vortices and Josephson phase slippage with the experimental J_c measuring the weaker. For b less than the crossover field, b^* , the Josephson critical current between grains of the film is smaller than the depinning critical current, which is presumed to continue to increase as $(1-b)^2$ as in the single NbN films. For LLH, the Lorentz force is a maximum and there will be a crossover to depinning (at b^*), if $J_{cj} < J_{c0}$, because of the weaker field dependence of J_{cj} . Such a crossover is not expected for IIIH, and J_c is significantly larger (~ 3 times) than for LLH. For IIIH, J_c (not shown) does not appear to follow the ideal Josephson behavior up to H_{c2} , but rather a mixed behavior which can possibly be understood as due to some residual Lorentz force due to inhomogeneous intergranular current paths.

2.3. Measurements in $Tl_2Ba_2CaCu_2O_x$

Resistive transitions for various H parallel to the Cu-O layers of epitaxial films of $Tl_2Ba_2CaCu_2O_x$ are shown in Fig. 3 for both the maximum and minimum Lorentz force orientations (HLL and HLL), where J is the current density, equal to 10 A/cm^2 . This independence of the dissipation on macroscopic Lorentz force is also found in $R(T)$, the low-temperature J_c and $I(V)$ for polycrystalline $Tl_2Ba_2CaCu_2O_x$ films⁵. As in the NbN/AlN multilayers, a dissipation mechanism other than vortex motion is suggested here⁶: again it is reasonable due to the significance of the intrinsic vortex pinning, expected by the wide insulator spacing, in this case between Cu-O bilayers. Misalignment of the external field or crystal axes can be ruled out as an explanation, because an *average* value of 3.8° would be needed to match data for HLLc. Information on the nature of the transitions is obtained by fitting to:

$$\rho_{ab}(T, H) = \rho_0 \exp\left(-\frac{U(T, H)}{T}\right), \quad (1)$$

Fig. 3. Resistive transitions in a c-axis oriented, epitaxial thin film of $\text{Tl}_2\text{Ba}_2\text{CaCu}_2\text{O}_x$ for fields parallel to the layers. Data for J parallel and perpendicular to the field are virtually indistinguishable.

where $U(T, H)$ is the activation energy. The straight lines fit the data of Fig. 3 give $U \propto (1-t)/B^{0.55}$, where $t \equiv T/T_c$. Following the results in NbN/AlN multilayers the $\text{Tl}_2\text{Ba}_2\text{CaCu}_2\text{O}_x$ data can be interpreted as due to a Josephson junction, for which U is given by E_{cj}/k_B , which in zero field is¹³:

$$E_{cj} = \frac{\hbar}{2e} I_{cj}, \quad (1)$$

$$I_{cj}(T) = \frac{\pi \Delta(T)}{2e R_N} \tanh\left(\frac{\Delta(T)}{2k_B T}\right), \quad (2)$$

where R_N is the normal-state junction resistance (e.g., above T_c) and $\Delta(T)$ is the energy gap. The resulting $U \propto (1-t)$ near to T_c has been confirmed by measurements in granular NbN films¹². While the temperature dependence for the $\text{Tl}_2\text{Ba}_2\text{CaCu}_2\text{O}_x$ agrees, the field dependence is hard to understand within a simple Josephson junction model (see Section 3.1 below), so the exact nature of the residual non-vortex motion dissipation in $\text{Tl}_2\text{Ba}_2\text{CaCu}_2\text{O}_x$ for $H \parallel ab$ is still unclear.

3. FIELD PERPENDICULAR TO LAYERS

In this section, we concentrate on the intrinsic multilayers of Cu-O bilayers found in HTS. For $H \parallel c$, it is reasonable that, due to the lack of intrinsic pinning, the much greater broadening of the resistive transition is due to thermally-activated flux motion¹⁻³. This irreversibility behavior can be quantified experimentally by a characteristic field, H^* , such that $\rho_{ab}(T, H^*)/\rho_n = 10^{-5}$, where ρ_n is the normal-state, ab-plane resistivity at T_c . Such H^* are plotted in Fig. 4 against $1-t$ for many samples of various HTS, including some from the literature¹⁵⁻¹⁷, showing that H^* is fairly independent of sample (e.g., epitaxial or polycrystalline films, single crystals, minor amounts of second-phase material, etc.), at least if our uniform criteria is used. An important step for understanding the case of $H \parallel c$ was the recognition that for

Fig. 4. The characteristic fields for thermal activation for various HTS fall in four groups. From left, $\text{YBa}_2\text{Cu}_3\text{O}_7$; $\text{TiBa}_2\text{CaCu}_2\text{O}_x$ and $\text{TiBa}_2\text{Ca}_2\text{Cu}_3\text{O}_x$; $\text{Ti}_2\text{Ba}_2\text{CaCu}_2\text{O}_x$; and $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x$. See text.

sufficiently weak coupling between superconducting layers, implied by a high degree of anisotropy, that the usual concept of three dimensional (3D) vortex lines should be replaced with one of 2D pancake-like⁹ vortices confined to individual Cu-O bilayers. In addition, this weak coupling allows for *thermally-activated* decoupling of pancake vortices in adjacent layers, which greatly reduces the effectiveness of pinning (e.g., point defects in superconducting layers), when compared to extended, 3D vortex *lines*.

3.1. Nature of the interlayer vortex coupling

The nature and strength of pancake vortex coupling between adjacent layers is crucially important to the vortex dimensionality and dynamics. It is necessary to consider the relative importance of various types of coupling between layers: electromagnetic, i.e., due magnetic fields or currents associated with the vortices; and Josephson, i.e., due to the superconducting pairs' wave function.

Electromagnetic coupling was first studied in the ingenious experiments of Giaever¹⁸, in which a perpendicular field was applied to two electrically insulated, superimposed, superconducting films with vortices threading both films. The Lorentz force due to a current in one film causes flux motion, and vortices in the secondary film are dragged along until the maximum electromagnetic coupling force, F_{cm} , is exceeded, and then a secondary voltage is observed due to the Faraday effect. The theoretical analysis of Clem¹⁹ has been adequately verified near T_c in small fields by Ekin and coworkers²⁰ and at significantly higher fields and lower temperatures by Gray²¹. In particular, the very strong, $F_{cm} \sim \lambda^{-4}$, penetration-length, λ , dependence was verified.

Clem²² has recently extended his analysis of electromagnetic coupling to the case $B \gg B_{c1}$ (where B_{c1} is the lower critical field and the vortex spacing is smaller than the penetration length *in the layers*, $\lambda_{||}$), and found the energy of a pancake vortex, misaligned at a distance r (with s the layer repeat distance), to be:

$$U_{em}(r) \sim \{(\phi_0^2/4\pi^2)(1+\gamma) s/16\lambda_{||}^4\} r^2 \equiv e_{cm} r^2, \quad (3)$$

where $\gamma \sim 0.577215$. Note the strong temperature dependence of $U_{em} \sim (1-t)^2$ near T_c , where $t \equiv T/T_c$. In finite fields, the flux-lattice spacing, $r_B \sim \sqrt{\Phi_0/B}$, defines the spatial periodicity of the fields and currents, so that U_{em} can never be larger than $U_{em}(r_B)$. For the highly anisotropic Bi- and Tl-based cuprates, $s \sim 1.5$ nm and $\lambda_{||} \sim 230$ nm, so the characteristic temperature for loss of alignment is given by $T_B \equiv U_{em}(r_B)/k_B \equiv Q/B$ with $Q \sim 8.4$ K T. Thus for $B > 1$ T, one finds $T_B < T_c$ in $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$.

Josephson coupling has its origin in the energy costs associated with any differences in the *phase* of the superconducting wave function across a Josephson junction. Zero phase difference occurs when vortices are aligned across a junction, but when they are misaligned, there is necessarily a spatially dependent phase difference across the junction. Clem²³ has calculated, U_j , the energy for a pancake vortex misaligned at a small distance r to be:

$$U_j = e_j r^2, \quad \text{for} \quad r \ll r_c = s \frac{\lambda_c}{\lambda_{ab}} = s \sqrt{\frac{m_c}{m_{ab}}}, \quad (4)$$

$$e_j = \frac{\hbar I_{cj}(T)}{eA} \quad (5)$$

where $I_{cj}(T)$ is the unfluctuated Josephson critical current (Eq. 2), R_N is the normal-state junction resistance, A is the junction area and $\Delta(T)$ is the superconducting energy gap, while for $r \gg r_c$, $U_j = e_j r_c r$. For $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$, the effective-mass anisotropy of $\sim 10^5$ means the crossover occurs for $r_c \sim 450$ nm, which corresponds to the average flux spacing at 0.01 T. Thus Eq. 4 is valid for most experimental situations, and similar to electromagnetic coupling, the maximum value of U_j in a finite field, U_{jmax} , is

$$U_{jmax} \sim U_j(r_B) \sim e_j \Phi_0/B, \quad (6)$$

in this case because of the spatially periodic superconducting wave function in an Abrikosov lattice. The temperature dependence of U_j is given by Eq. 2; near T_c the dependence is $(1-t)$, and therefore weaker than the $(1-t)^2$ dependence for electromagnetic coupling. From standard tunneling theory²⁴ and experiment²⁵, e_j should be proportional to $\exp(d_i/d_0)$, where d_i is the insulator thickness and $d_0 = \hbar/\sqrt{8m\Phi}$ accounts for the tunneling barrier height, Φ . It is important to thus recognize that U_{jmax} is intrinsic, albeit field-dependent, with a sample-dependent, zero-field cutoff given by $e_j A$.

There is a useful analogy between superconducting multilayers, e.g., in the highly-anisotropic HTS cuprates, a *pair* of Cu-O bilayers separated by a region of weaker superconductivity, and *discrete* Josephson tunneling junctions, comprised usually of two metallic superconducting films separated by an insulating oxide of one of them. Discrete Josephson junctions in a perpendicular magnetic field, have not received the attention that the parallel field case has. However, recent experiments on small-area, thin-film Nb tunnel junctions²⁶, in which e_j can be separately measured, have shown that the appropriate U_j is in excellent *quantitative* agreement with Eq. 6.

The experimental evidence which implies that Josephson coupling dominates in the cuprates is, briefly: (1) that the coupling energy⁸ for a series of cuprates depends exponentially on the insulator spacing (between adjacent Cu-O bilayers) with a reasonable tunneling barrier height of ~ 0.8 eV; and (2) that the 'irreversibility behavior' defined below by a crossover field, H^* , shows the $1-t$ dependence expected for Josephson coupling and not the $(1-t)^2$ dependence predicted for electromagnetic coupling²². The analogy between discrete Josephson junctions and layered cuprates is further supported by microwave emission experiments on $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ and $\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$ crystals²⁷, in which it is observed that each pair of bi- or tri-layers of Cu-O forms a separate working Josephson contact.

Fig. 5. Characteristic fields for thermally-activated 3D to 2D crossover, flux depinning and the combined effect of both to produce H^* . Also shown is H_{c2} .

3.2. Josephson interlayer coupling model for HTS

The dimensional crossover occurs when the coupling energy between adjacent Cu-O multilayers, $U_{jmax}(H,T)$, is $\sim k_B T$. For Josephson tunneling through an insulating barrier, U_{jmax} given by Eq. 6 must include a $(1-b)$ term to account for the spatially-averaged order parameter over the flux-line lattice. The total Josephson coupling of a vortex to *both* neighboring Cu-O layers is $2U_{jmax}$. The resulting 3D to 2D dimensional crossover is shown as a thin solid line in Fig. 5 for $[\pi \Phi_0 / e^2 \rho_{nc} s B_{c2}(0)] [\Delta(0) / k_B T_c] = 0.05$, where ρ_{nc} is the normal-state c-axis resistivity. However, even in totally decoupled layers, the 2D pancake vortices must overcome some finite pinning energy. The energy associated with depinning is written as that fraction, α , of the loss in superconducting condensation energy in a vortex core, which is compensated by a particular pinning site¹⁰:

$$U_p = \alpha \pi \xi_{ab}^2 d_s \frac{B_c^2}{2\mu_0} (1-b)^2. \quad (7)$$

where B_c is the thermodynamic critical field. An example of thermally-activated depinning is also shown in Fig. 5 as a dashed line for $\alpha \Phi_0 d_s B_c(0)^2 / 4\mu_0 k_B T_c B_{c2}(0) = 0.5$. Vortices must overcome *both* energy barriers, so a characteristic field for dissipation, H^* , can be defined theoretically by solving $k_B T = U_p(H,T) + 2U_{jmax}(H,T)$, and this is shown in Fig. 5 as the heavy solid line.

The H^* data of Fig. 4 are thus fit to H^* of the Josephson model using three parameters: H_{c2} ; ρ_c , which is related to the strength of Josephson coupling; and $\alpha B_c^2 / 2\mu_0$, which is related to the 2D pinning strength. The fits are relatively good and the parameters derived agree reasonably well with other measurements⁸. An important further prediction of coupling by Josephson tunneling is an exponential dependence of ρ_c on d_i , and the data are shown in Fig. 6 for the Tl- and Bi-based HTS materials. While these data do not fall on one straight line, we anticipate the possibility of different tunneling barrier heights or tunneling effective masses, at least between Bi-O and Tl-O insulating layers. Nevertheless, the best fit

Fig. 6. Values of ρ_c for the fits of the model to various single-crystalline or epitaxial materials indicating the exponential dependence on the Cu-O multilayer spacing d_i , expected for Josephson tunneling.

gives a reasonable value of $d_0=1$ Å and a barrier height of 0.8 eV, assuming a free electron mass. Values of $d_0=1-1.6$ Å are typically found in artificial or natural oxide tunneling barriers^{24,25}. Note that data for fully-oxygenated $\text{YBa}_2\text{Cu}_3\text{O}_7$ could not be fit to the model with a reasonable value of ρ_c : because of its must lower anisotropy, $\text{YBa}_2\text{Cu}_3\text{O}_7$ appears to be 3D over most of the accessible H-T phase diagram.

3.3. C-axis resistivity

Recent results²⁸ in $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ single crystals reveal a magnetic-field dependent peak near T_c for H||c, which has been confirmed by later studies²⁹. Initial attempts to explain the field-dependent transition at low temperatures by 1D phase slippage²⁸ or by fluctuation effects³⁰, gave no explanation of the empirical 'semiconducting' behavior of the normal-state. Consistent with the irreversibility model presented in Section 3.1, a single explanation for both the low-T, field-dependent transition and the high-T, field-independent normal-state 'semiconducting' behavior has been constructed³¹ by considering the c-axis conduction to be that of a stack of ordinary Josephson tunnel junctions in series, connecting well-coupled units of superconducting Cu-O bilayers.

A model for the conductance of a Josephson junction as the sum of that due to pair tunneling and that due to the tunneling of unpaired electrons was proposed: it shows a peak with temperature and was verified with data on discrete, Al thin-film Josephson junctions³¹. Using 2D fluctuations, calculated beyond the Gaussian approximation³², and values of Δ measured by tunneling³³⁻³⁵, good agreement was found³¹ with the measured $\rho_c(T, H)$. In Fig. 7, the data²⁸ are shown as open symbols, together with a two-parameter fit to the above model ($\gamma \equiv E_j(0)/k_B T_c$ to match the field-dependent resistance increase below the peak and to match the field-independent shape above the peak, d_c , the thickness of the isolated conducting bilayers). From the field-independent decrease at high temperatures, d_c is found to be ~0.6 nm which is less than s , confirming that the 2D fluctuation model is appropriate, but greater than the Cu-O bilayer spacing of 0.3 nm, implying that the bilayers can be treated as a single layer. As an independent check, the 2D fluctuation amplitude thus obtained predicts³¹ the absolute magnitude of the fluctuation conductivity in ab-plane transport within 20%. At low temperatures, the field-dependent increase of ρ_c is

Fig. 7. Data²⁸ on the c-axis resistivity of $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ single crystals fit to the model³¹ described in the text. Solid diamonds are break-junction tunneling results³⁵.

fit by $\gamma \sim 4.1$ at 0.5 T and ~ 1.1 at 7 T, which shows qualitatively the field dependence of Eq. 6 and gives values of $\rho_{cN} \sim 1.4$ and $0.4 \, \Omega\text{cm}$, respectively. These are reasonably consistent with the measured²⁸ value of ρ_c well above T_c , which is $\sim 1 \, \Omega\text{cm}$. Also shown in Fig. 7 as solid diamonds are zero-bias tunneling resistance measurements in discrete break-junctions³⁵ of $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ single crystals: the excellent agreement with the measured ρ_c of Ref. 28 above the peak supports the contention that the origin of $\rho_c(H,T)$ is ordinary tunneling. Note that Josephson coupling is weak in these high- R_N break-junctions and is seen³⁵ only at much lower temperatures (with $R_0/R_N \geq 10$).

4. SUMMARY OF CONCLUSIONS

The dissipation associated with the broadened resistive transitions of the highly-anisotropic high-temperature superconductors has been studied for fields parallel and perpendicular to the superconducting layers. The absence of any Lorentz force for H parallel to the Cu-O layers is shown, by analogy with artificial multilayers of superconducting NbN with insulating AlN, to be due to the strong intrinsic pinning of the layered structure, such that the residual measured dissipation is due to a mechanism not involving vortex motion. For H perpendicular to the Cu-O layers, vortex motion causes the dissipation, but only after a crossover from 3D vortex lines to 2D vortex pancakes. This intrinsic crossover is determined by the Josephson coupling between Cu-O bi- or tri-layers and thus is unaffected by additional flux pinning. At low temperatures, pinning can be effective even in the uncoupled 2D layers, and can also increase H^* and the critical current density in this model.

5. ACKNOWLEDGEMENTS

This research is supported by the U.S. Department of Energy, Basic Energy Sciences-Materials Sciences under contract #W-31-109-ENG-38, the National Science Foundation-Office of Science and Technology Centers under contract DMR 91-20000, and was done in collaboration with R.T. Kampwirth, J.M. Murduck, D.M. McKay, K.C. Woo, S.J. Stein, J.C. Smith, D.S. Richeson, T.J. Marks, J.H. Kang, J. Talvacchio, M. Eddy, D. Hampshire, R. Herzog and H.W. Weber.

6. REFERENCES

1. P.H. Kes, J. Aarts, J. van den Berg, C.J. van der Beek and J.A. Mydosh, *Supercond. Sci. Technol.* **1**, 242 (1989).
2. J.M. Kosterlitz and D.J. Thouless *J. Phys.* **C6**, 1181 (1973); B.A. Huberman and S. Doniach, *Phys. Rev. Lett.* **43**, 950 (1979).
3. M.P.A. Fisher, *Phys. Rev. Lett.* **62**, 1415 (1989).
4. Y. Iye, S. Nakamura and T. Tamegai, *Physica C* **159**, 433 (1989); K. Kitazawa, S. Kambe, M. Naito, I. Tanaka and H. Kojima, *Japan. J. Appl. Phys.* **28**, L555 (1989); H. Iwasaki, N. Kobayashi, M. Kikuchi, T. Kajitani, Y. Syono, Y. Muto and S. Nakajima, *Physica C* **159**, 301 (1989).
5. K.C. Woo, K.E. Gray, R.T. Kampwirth, J.H. Kang, S.J. Stein, R. East and D.M. McKay, *Phys. Rev. Lett.* **63**, 1877 (1989).
6. D.H. Kim, K.E. Gray, R.T. Kampwirth and D.M. McKay, *Phys. Rev.* **B42**, 6249 (1990).
7. K.E. Gray, R.T. Kampwirth, D.J. Miller, J.M. Murduck, D. Hampshire, R. Herzog and H.W. Weber, *Physica C* **174**, 340 (1991).
8. D.H. Kim, K.E. Gray, R.T. Kampwirth, J.C. Smith, D.S. Richeson, T.J. Marks, J.H. Kang, J. Talvacchio and M. Eddy, *Physica C* **177**, 431 (1991).
9. J.R. Clem, *Phys. Rev.* **B43**, 7837 (1991).
10. K.E. Gray, R.T. Kampwirth, J.M. Murduck and D.W. Capone II, *Physica C* **152**, 445 (1988).
11. H.L. Ho, R.T. Kampwirth, K.E. Gray, D.W. Capone II and L.S. Chumbley, *Ultramicroscopy* **22**, 297 (1987).
12. D.H. Kim, K.E. Gray, R.T. Kampwirth, K.C. Woo, D.M. McKay and J. Stein, *Phys. Rev.* **B41**, 11642 (1990).
13. P.W. Anderson, Lectures at Ravello Spring School, 1963; V. Ambegaokar and A. Baratoff, *Phys. Rev. Lett.* **10**, 486 (1963).
14. A.A. Abrikosov, *Zh. Eksperim. i Teor. Fiz.* **32**, 1442 (1957); *Soviet Phys. JETP* **5**, 1174 (1957), and also A.L. Fetter and P.C. Hohenberg, in Superconductivity, edited by R.D. Parks, (Marcel Dekker, New York, 1969), p.851.
15. T.T.M. Palstra, B. Batlogg, L.F. Schneemeyer and J.V. Waszczak, *Phys. Rev. Lett.* **61**, 1662 (1988).
16. T.T.M. Palstra, B. Batlogg, R.B. van Dover, L.F. Schneemeyer and J.V. Waszczak, *Appl. Phys. Lett.* **54**, 763 (1989).
17. W.K. Kwok, (private communication).
18. I. Giaever, *Phys. Rev. Lett.* **15**, 825 (1966); *Phys. Rev. Lett.* **16**, 460 (1966).
19. J.R. Clem, *Phys. Rev.* **B9**, 898 (1974); *Phys. Rev.* **B12**, 1742 (1975).

20. J.W. Ekin, B. Serin and J.R. Clem, *Phys. Rev.* **B9**, 912 (1974).
21. K.E. Gray, *Phys. Rev.* **B28**, 1281 (1983).
22. J.R. Clem, *Phys. Rev.* **B43**, 7837 (1991).
23. J.R. Clem, private communication.
24. See e.g., E.L. Wolf, Principles of Electron Tunneling Spectroscopy (Oxford Univ. Press, New York, 1985).
25. K.E. Gray, *Solid State Commun.* **13**, 1787 (1973).
26. D.H. Kim, K.E. Gray and J.H. Kang, *Phys. Rev.* **B45**, 7563 (1992).
27. R. Kleiner, F. Steinmeyer, G. Kunkel and P. Müller, *Phys. Rev. Lett.* **68**, 2394 (1992); R. Kleiner and P. Müller, preprint.
28. G. Briceno, M.F. Crommie and A. Zettl, *Phys. Rev. Lett.* **66**, 2164 (1991).
29. Y.I. Latyshev and A.F. Volkov, *Physica C* **182**, 47 (1991); and K. Kadowaki, J.N. Li and J.J.M. Franse, *Physica C* **170**, 298 (1990).
30. R. Ikeda, T. Ohmi and T. Tsuneto, *Phys. Rev. Lett.* **67**, 3874 (1991); *J. Phys. Soc. Jpn.* **60**, 1051 (1991).
31. K.E. Gray and D.H. Kim, *Phys. Rev. Lett.* **70**, 1693 (1993).
32. S. Ullah and A.T. Dorsey, *Phys. Rev.* **B44**, 262 (1991).
33. Q. Huang, J.F. Zasadzinski, K.E. Gray, J.Z. Liu and H. Claus, *Phys. Rev.* **B40**, 9366 (1989); Q. Huang, J.F. Zasadzinski, K.E. Gray, E.D. Bukowski and D.M. Ginsberg, *Physica C* **161**, 141 (1989).
34. J.X. Liu, J.C. Wan, A.M. Goldman, Y.C. Chang and P.Z. Jiang, *Phys. Rev. Lett.* **67**, 2195 (1991).
35. D. Mandrus, L. Forro, D. Koller and L. Mihaly, *Nature* **351**, 460 (1991).

DATE
FILMED

7/25/94

END