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Lorentz-force independence of resistance tails in magnetic fields near T_c for the low-temperature superconductor granular NbN: A Josephson-junction model

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The broadened resistive transitions in a granular NbN film have been measured for field orientations *both perpendicular and parallel to the current.* These are essentially the same over a span of 6 orders of magnitude below the normal-state resistance (just above T_c) demonstrating the absence of the macroscopic Lorentz force. We show that a model of Josephson coupling between grains, rather than flux motion, can *quantitatively* explain the results with no adjustable parameters. Broadening of the resistive transitions of high-temperature superconductors (HTS) has also been shown to be independent of the macroscopic Lorentz force. This research demonstrates unambiguously that resistive losses near T_c and the absence of the macroscopic Lorentz force on them are not restricted to exotic HTS.

The broadened resistive transitions in high-temperature superconductors (HTS) in a magnetic field are dramatic in[1] Bi-Sr-Ca-Cu-O and[2] Tl-Ba-Ca-Cu-O materials but are also found in[3] $YBa_2Cu_3O_7$. This behavior has led authors to suggest thermally activated magnetic-flux creep,[1,4] flux-line melting,[5] or another unknown loss mechanism.[1] Recent studies[6-8] of resistive transitions in HTS report essentially no difference for the current *perpendicular or parallel to the field*, indicating that the *macroscopic* Lorentz force does not play a role in this loss mechanism, which questions explanations based on flux motion. In order to better understand the origin of this unexpected result, we have studied the low-temperature superconductor, NbN. The broadened resistive transitions for field orientations *both perpendicular and parallel to the current* are essentially the same over a span of 6 orders of magnitude below the normal-state resistance (just above T_c) for the granular NbN film. We further expand on a previously suggested model[9] of Josephson coupling between the grains in NbN, rather than flux motion, to *quantitatively* explain the results with no adjustable parameters.

Granular NbN films are prepared on polished sapphire substrates using magnetron reactive sputtering.[10] The substrate temperature was maintained at $\approx 300\,^\circ C$ during deposition, while the system base pressure prior to deposition was typically $\leq 5\times10^{-7}$ torr. The sputtering conditions are typical for films made under conditions optimized for the highest J_c in a 20-T field. These optimized conditions were determined by finding the Ar pressure (at a fixed nominal N_2 partial pressure) which gave the highest J_c at 20 T, then varying the N_2 partial pressure to maximize J_c. Samples were etched into 1-mm-wide strips using standard photolithographic techniques. Thickness was measured to be 1.12 ± 0.05 μm using a surface profilometer. Transverse section TEM studies[11] of similar films have indicated a conical, columnar grain structure. Grain diameters are 10 nm at the substrate and increase linearly with thickness to 100 nm at 2 μm thickness.

Samples were mounted in a variable temperature insert of a 6.5-T superconducting solenoid. Resistance measurements (ac) were taken with a standard lock-in technique as a heater surrounding the sample chamber increased the temperature slowly through the transition. Strictly identical thermal conditions (starting temperature and heating rate) were used for all runs. Standard magnetoresistance corrections were used for the carbon glass thermometer, but these do not affect the relative measurements between current parallel and perpendicular to the field.

The data are shown in Fig. 1 for various fields, which together with the current are parallel to the film plane, but both field orientations with respect to the current are displayed. It is apparent that for all fields up to 5 T, the differences for both field orientations with respect to the current are minor compared to the overall effect. The broadening is therefore not dependent on the *macroscopic* Lorentz force which is zero for the field parallel to the current. Although one can consider that the current and/or flux-line direction may deviate from the macro-

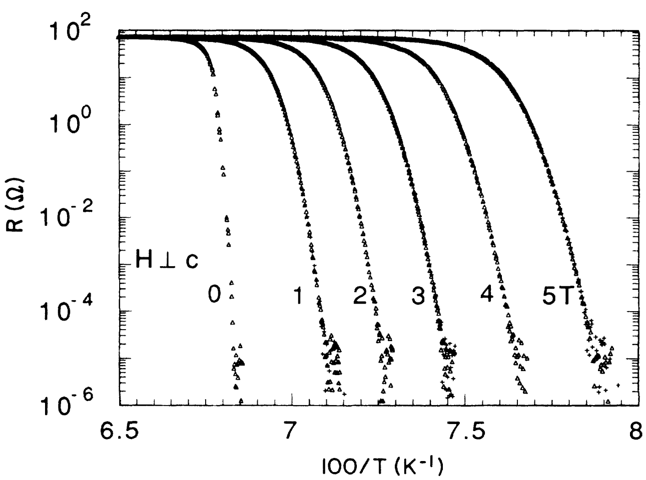

FIG. 1. Resistive transitions for a granular NbN film for various fields, which together with the current I, are parallel to the film plane: Δ, H⊥I; and +, H∥I. It is apparent that for fields up to 5 T, the differences for both field orientations with respect to the current are minor compared to the overall effect.

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scopic average and identify flux flow dissipation mechanisms for each field orientation, it is hard to imagine *any* flux flow scenario, including flux cutting, which would give identical dissipation for both orientations. In addition, in spite of the columnar grains, the anisotropy of granular NbN, with respect to the columnar grain axes, is much smaller than for the HTS (e.g., $H_{c2\parallel c} \sim 1.25 H_{c2\perp c}$). Therefore, we look to explanations other than flux flow for these losses.

Recently, studies[12] of the critical current density in multilayers of NbN and AlN showed a crossover from flux pinning to Josephson coupling as the field strength was varied. In the Josephson coupling regime, the critical current density was independent of the macroscopic Lorentz force. In the following, we present a model of Josephson coupling between grains which can *quantitatively* explain the above results with no adjustable parameters.

We begin with the theory of Ambegaokar and Halperin[13] for the dissipation (resistance) in a Josephson junction due to *phase* fluctuations of the order parameter, $\psi e^{i\phi}$. When $k_B T$ becomes comparable with the Josephson coupling energy $\hbar I_{cj}/e$, the resistance $R(T)$ in the limit of low bias current, $I \ll I_{cj}$, is given by[13]

$$R(T)/R_N = [I_0(\gamma/2)]^{-2}, \qquad (1)$$

where I_0 is the modified Bessel function, R_N is the normal-state resistance of the junction, and $\gamma = \hbar I_{cj}/e k_B T$. The Josephson critical current I_{cj} also depends[14] on the energy gap $\Delta(T)$, so

$$I_{cj}(T) = \frac{\pi\Delta(T)}{2eR_N}\tanh\left[\frac{\Delta(T)}{2k_BT}\right], \qquad (2)$$

or the corresponding low-temperature limit

$$I_{cj}(0) \sim \pi\Delta(0)/2eR_N . \qquad (3)$$

Since the experimental data are very near T_c, we expand Eq. (2) for $\Delta(T) \ll k_B T$ and substitute the appropriate BCS approximation, $\Delta(T) = 3.2 k_B T_c \sqrt{1-t}$, where $t \equiv T/T_c$, so

$$\gamma = 2.56\pi \frac{R_0}{R_N}\frac{1-t}{t^2}, \qquad (4)$$

where $R_0 \equiv \hbar/e^2 = 4114\ \Omega$. In order to verify this model, it is necessary to demonstrate that independently measured values of R_N, $I_{cj}(0)$, and $\Delta(0)$ are consistent with Eqs. (1), (3), and (4).

Individual Josephson junctions between the columnar grains of the NbN film cannot be directly measured, but average values of the above parameters can be determined. For example, R_N corresponds to the average resistance between grains, $\langle R_N \rangle = \rho_N/d = 9.6\ \Omega$, where ρ_N is the normal-state resistivity of 1.08 mΩ cm and d is the film thickness of 1.12 μm. The value of γ in zero field is determined by the best match of the data of Fig. 1 at the lowest temperatures. Nearer T_c, fluctuations in the *magnitude* of the order parameter ψ can account for the rounding of the experimental data for all field values. Putting γ in Eq. (4) gives $R_N = 19.7\ \Omega$, which is in relatively good agreement considering the granularity of the films. The difference may reflect the distribution width of actual R_N, in that the resistive tail is more likely to select the largest values. The independent evaluation of $I_{cj}(0)$ requires the measured critical current density $J_{cj}(0)$ and the area A of the junctions between grains. The latter has been determined from extensive transmission electron microscopy studies on similar films.[11] For the present case, the average contact area of the conical, columnar grains is estimated to be $\sim 0.04\ \mu m^2$. Using zero-field values of $J_{cj}(0)$ at 4.2 K and the R_N obtained from γ, then Eq. (3) implies that $\Delta(0) = 2.3$ meV. Tunneling data,[15] scaled to our experimental T_c of 14.7 K, gives $\Delta(0) \sim 2.5$ meV. The very reasonable agreement for two independent parameters (R_N and Δ) with the Josephson model is strong evidence for its validity to describe conduction in these granular NbN films.

An explanation of the field dependence requires extending the zero-field Josephson model, and we start with the simple interpretation used in Ref. 12. It was argued that since the magnetic-field penetration length is considerably greater than the grain diameters and the flux-line spacing for $H > 1$ T, the induction will be quite uniform. The consequent depairing will reduce ψ uniformly and, because of the short electron mean free path, isotropically, except for the normal cores of the vortices. For b close to 1, where $b \equiv H/H_{c2}$, Abrikosov's solution[16] of the vortex lattice determines the *spatial average* of ψ as $\langle \psi^2 \rangle = \psi_\infty^2(1-b)$.

Experimentally, we wish to connect $J_{cj}(0,H)$ to $\gamma(H)$ [or $\Omega(H) \equiv t^2\gamma(H)/(1-t)$ to remove the t dependence], as was done above for zero field. The measured values of $J_{cj}(0,H)$ and $\Omega(H)$, e.g., from the low-temperature portions of the data in Fig. 1, for fields applied both parallel and perpendicular to the columnar grains, are displayed in Fig. 2. The data all decrease linearly with H, although for $\mathbf{H} \perp c$, they do not extrapolate to the same zero-field point. The latter observation must result from the anisotropy of the columnar grains, but is not easily explained. The universal linear decrease with H is anticipated from Abrikosov's solution[16] $\langle \psi^2 \rangle = \psi_\infty^2(1-b)$ and Eq. (2), for which $I_{cj} \propto \Delta^2 \propto \psi^2$. The ratios of coefficients of the linear

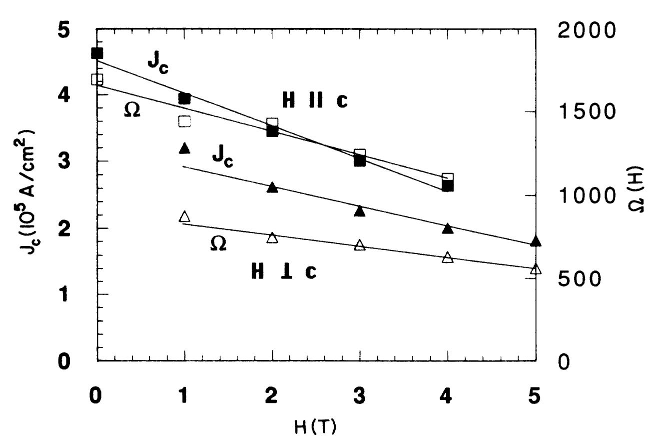

FIG. 2. The measured values of $J_{cj}(0,H)$ (solid symbols) and $\Omega(H)$ (open symbols), e.g., from the low-temperature portions of the data in Fig. 1, for $\mathbf{H}\parallel c$ (squares) and $\mathbf{H}\perp c$ (triangles).

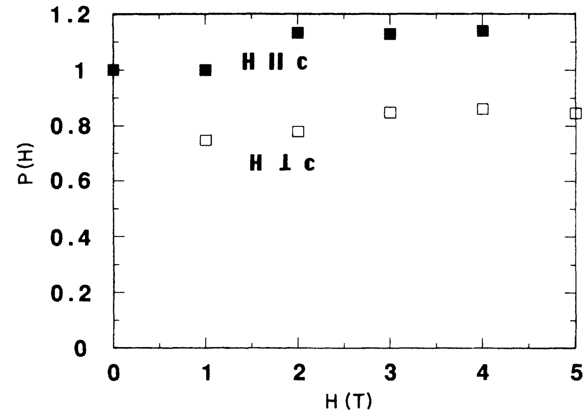

FIG. 3. Agreement of the field- and orientation-dependent data with the Josephson-junction model as represented by $P(H)$ from Eq. (5).

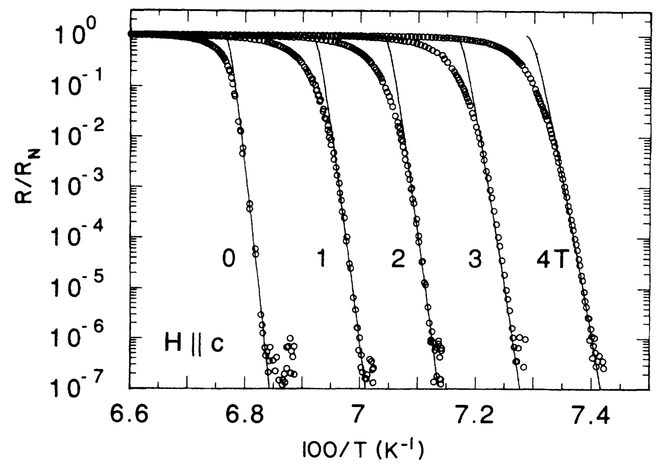

FIG. 4. The calculated $R(T)/R_N$ from Eqs. (1) and (5), (solid lines), using the appropriate values of $J_{cj}(0)/\Delta(0,0)$ for H$\|c$, together with the experimental data (open circles). Note that fluctuations in the *magnitude* of the order parameter ψ can account for the rounding of the experimental data near T_c for all field values.

decreasing terms, for H$\perp c$ and H$\| c$, are ~ 1.24 for *both* $\Omega(H)$ and $J_{cj}(0,H)$, in excellent agreement with the measured H_{c2} anisotropy of 1.25. However, using the measured $H_{c2\perp c} \sim 22.8$ T, one anticipates a field dependence of $1\text{-}0.044H(T)$, whereas the data of Fig. 2 give larger coefficients [by a factor of ~ 1.5 for $\Omega(H)$ and 2 for $J_{cj}(0,H)$]. One might explain this by (i) concentration of the cores at the Josephson junctions due to the minimum of the core potential resulting from the finite width of the insulating grain boundaries[17] and/or (ii) misalignment of the cores across the junction, resulting in a further reduced $\langle\psi^2\rangle$.

The connection of $\gamma(H)$ with $J_{cj}(0,H)$ is found by rederiving Eq. (4), using the appropriate field dependences of the spatial average of Δ in Eqs. (2) and (3)

$$\Omega(H) \equiv \frac{t^2\gamma(H)}{(1-t)} = 5.12\frac{eR_0J_{cj}(0,H)A}{\Delta(0,0)}P(H)\,, \quad (5)$$

where $P(H)$ accounts for the fact that the field dependences of $\langle\Delta\rangle$ may be different since $J_{cj}(0,H)$ and $\Omega(H)$ are measured at different temperatures. Unfortunately, $P(H)$ cannot be simply calculated [in the zero-field case, $P(0)=1$]. Using the measured values of $J_{cj}(0,H)$ and $\Omega(H)$, Eq. (5) is solved for $P(H)$, and the results plotted in Fig. 3. Overall, the agreement is excellent considering the simplicity of the model and/or experimental errors. The abrupt drop of $P(H)$ from zero field to finite H for H$\perp c$ reflects the residual effect of similar drops in $J_{cj\perp c}(0,H)$ and $\Omega_{\perp c}(H)$ shown in Fig. 2. One expects a weak-field dependence for $P(H)$ since the flux-line spacing is smaller than the grain size for $H > 1$ T, and thus the spatial profile of ψ will scale with the equivalent $H/H_{c2}(T)$ for all the tails.

Figure 4 shows the calculated $R(T)/R_N$ from Eqs. (1) and (5), using the appropriate values of $J_{cj}(0)/\Delta(0,0)$ for H$\|c$, together with the experimental data. Note that fluctuations in the *magnitude* of the order parameter ψ can account for the rounding of the experimental data near T_c for all field values.

The overall agreement of the data with the Josephson-junction model implies that it correctly describes the transport in these granular NbN films. Of greater significance is the fact that the model can explain the absence of any effect of the macroscopic Lorentz force on the resistive dissipation just below T_c. This may have implications on similar measurements in high-temperature superconductors which show the independence of resistive losses near T_c on the Lorentz force,[6–9] but it should be emphasized that those HTS measurements were made on single crystals or very large-grained thin films which do not have obvious structural defects to produce Josephson junctions as do the NbN films. This research also demonstrates unambiguously such effects, the resistive losses, and the absence of the macroscopic Lorentz force on them, are not restricted to exotic HTS.

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