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Josephson tunneling for magnetic fields perpendicular to discrete Nb junctions: Implications for the irreversibility crossover in high-temperature superconductors

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It was recently suggested that coupling between the superconducting Cu-O bilayer or trilayer units in the *highly anisotropic*, high-temperature superconductors (HAHTS) can be described by *incoherent* Josephson tunneling. This paper confirms the assumed inverse dependence of the coupling energy on field, H, through measurements of the zero-bias resistance of discrete, high-quality, thin-film Nb Josephson junctions. This $1/H$ dependence was used to explain the flux-motion-induced broadening of resistive transitions in HAHTS for equivalent geometry with H parallel to the c axis. The Nb junction results also unambiguously confirm another recent suggestion: that field-induced dissipation can occur in Josephson junctions without the motion of vortices from an externally applied field, which in this case are pinned in the Nb electrodes.

Recently,[1] the effects of anisotropy and fluctuations on the field-induced broadening of resistivity transitions, $\rho(T,H)$, and the critical current densities, $J_c(T,H)$, have been studied in high-temperature superconductors (HTS). For *highly anisotropic* HTS in an applied field, H, oriented parallel to the c axis ($H \parallel c$) one explanation considers fluctuations which affect the Josephson coupling across the interlayer junctions.[2] Fluctuations of the relative phase across these junctions occur if kT exceeds $E_{cj}(T,H)$, the Josephson coupling energy[3,4] between adjacent Cu-O bilayer or trilayer units (multilayers). For H perpendicular to the interlayer Josephson junctions, this results in a crossover from three-dimensional (3D) vortex lines to 2D pancake vortices in the isolated Cu-O multilayers, and the experimental fits to this model[2] indicated that the effective junction area in E_{cj} is Φ_0/H.

Josephson fluctuations have also been used recently to explain the broadened $\rho(T,H)$ in *granular* NbN films[5] and $J_c(T,H)$ in *granular* multilayers[6] of NbN with AlN. It was suggested that motion of the external flux could be suppressed by the relatively strong pinning, e.g., in the insulating AlN layers where a distinct crossover in the field dependence of $J_c(H)$ was observed between depinning of the *external* flux and dephasing of intergranular Josephson junctions.[6]

The results presented in this paper confirm both of these ideas by direct measurements on discrete, high-quality Josephson junctions made with Nb films. For fields perpendicular to the film plane (up to 0.03 T), broadened resistive transitions which are very similar to those found in HTS materials are observed in these junctions. For $H \geq 0.005$ T, the experimental activation energy *quantitatively* equaled the theoretical[3] $E_{cj}(T,H)$ with an effective junction area of Φ_0/H. These measurements used a current density of 0.1 A/cm^2, for which the resistive transitions of the Nb electrodes were very sharp, indicating that the external flux was completely pinned in the electrodes. Thus the field-induced dissipation was not caused by motion of the external flux, but by self-field, Josephson vortex cores which are independent of, and perpendicular to, the applied field.

The Nb films were sputter deposited in the system equipped with load-lock chamber. The substrate was a 2-in. silicon wafer, onto which a 200-nm layer of SiO_2 was formed by a thermal process. The background pressure of the system was maintained at $\sim 10^{-9}$ torr. During film deposition, the substrates were kept in contact with a water-cooled copper plate. The residual resistivity ratios of such Nb films were normally 8–9, with $T_c = 9.2$ K. The Nb junctions used in this paper were fabricated[7] on such films using an AlO_x barrier, which was formed by exposing a 7.5-nm-thick Al layer to 150 mtorr of pure oxygen gas for 30 min, which resulted in $J_c \sim 1500$ A/cm^2. Junction areas were defined by the selective niobium etching process (SNEP), which uses reactive ion etching and anodization of Nb. The figure of merit of these junctions, V_m, was normally about 50 mV, measured at $V = 2$ mV and $T = 4.2$ K. The magnetic-field penetration depth of the Nb films, λ_{Nb}, was 62 nm, as measured from diffraction patterns of Josephson critical current, while the Josephson penetration depth, λ_J, defined by

$$\lambda_J = \left[\frac{\hbar}{2e\mu_0 J_c (2\lambda_{Nb} + d_i)} \right]^{1/2}, \qquad (1)$$

was about 12 μm using the above parameter values and a reasonable insulator thickness, d_i.

Resistances were obtained using a standard four-terminal, ac lock-in technique. All measurements were taken in a gas-flow cryostat after cooling the sample from above T_c in the applied field perpendicular to the thin-film electrodes. The measurements described below show that $H_{c2}(T)$ is much larger than that of pure Nb: Consequently, $H_{c1}(T)$ will be much lower. From this, H_{c1} *at zero temperature* can be estimated from the dirty limit formula to be ~ 0.023 T. This fact plus the favorable width-to-thickness ratio (>50) of the films, explains why small fields can remain in the samples after field cooling. In order to achieve transitions to a Josephson supercurrent at low temperatures, it was necessary to cool the samples quickly (and even faster as the field increased). Possibly this was necessary to keep the flux cores frozen in their initial aligned positions in each electrode, thereby preventing any parallel component in the barrier due to misalignment. Note that in a parallel field, the junction resistance periodically returned towards the quasiparticle resistance of the junction (which is greater than the normal-state value) when an integral number of flux quanta filled the junction. This manifestation of the usual Fraunhofer pattern measured in $I_c(H)$ is seen here as a function of temperature, because of the temperature dependence of $\lambda_{Nb}(T)$. For each perpendicular-field value, the resistive transitions were measured repeatedly: For analysis, we used those with the steepest slope which were reproducible and showed no evidence of vortex cores from the external field being parallel to the junction in the temperature dependence.

The resistive transitions of one junction and its electrode are shown in Fig. 1 for a field of 0.03 T. These measurements use a current density of 0.1 A/cm^2, and the extremely sharp resistive transition in the electrode indicates that *the external flux is completely pinned at all temperatures,* $T<0.98T_c$. The initial drop in apparent junction resistance corresponds to the incomplete four-terminal cancellation of both electrode resistances when they are not superconducting. Figure 2 shows the complete field dependence of this junction's resistance. Note that $H_{c2}(T)$ can be extracted from the initial drop at $T_c(H)$, and is given by 0.27 T/K near T_c.

In the theory[4] of Josephson junctions for zero field, a finite junction resistance appears even below T_c due to thermal fluctuations and it is given [for $E_{cj}(T,H)\gg k_BT$] by $R\propto\exp[-E_{cj}(T,H)/k_BT]$ where k_B is the Boltzman constant. The Josephson coupling energy, E_{cj} is proportional to I_c, the critical current in the absence of thermal fluctuations, which is proportional to the product of the superconducting order parameters on each side of the junction, ψ_a and ψ_b, divided by the normal-state resistance, R_N. Near T_c, $\psi_a\psi_b\propto(1-t)$, where $t=T/T_c$, so for convenience we define $U_0(H)=E_{cj}(T,H)/(1-t)k_B$ to obtain $R\propto\exp(-U_0(1-t)/T)$. This expression is used to fit the resistance data of Fig. 2 and determine U_0, which is plotted in Fig. 3 together with the data for two other junctions. For $H\geq 0.005$ T, U_0 agrees *quantitatively* with theory[3] if an effective junction area of Φ_0/H is used for the R_N going into E_{cj}.

In zero field, the low-temperature I_c, in the absence of

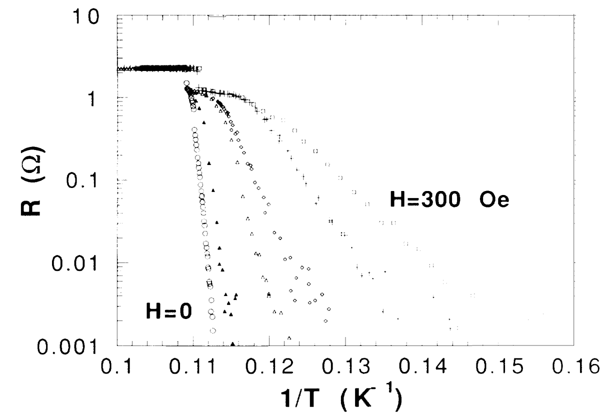

FIG. 2. The field dependence of the resistive broadening of Nb/Nb thin-film Josephson tunnel junctions looks very similar to high-temperature superconductors. Magnetic fields: 300 Oe (squares); 200 Oe (pluses); 100 Oe (diamonds); 50 Oe (open triangles); 20 Oe (solid triangles); zero applied field (circles).

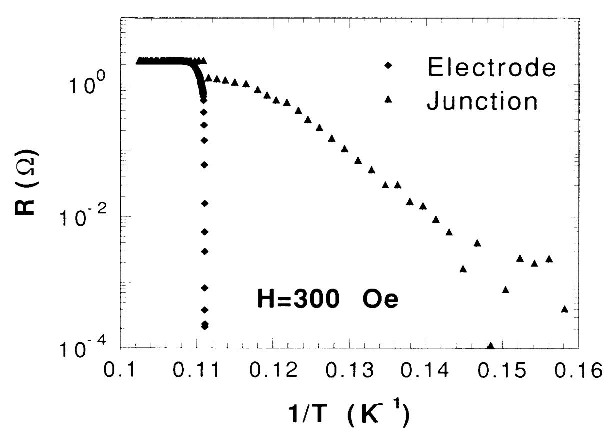

FIG. 1. The resistance of a Nb/Nb thin-film Josephson tunnel junction in a field of 300 Oe which is perpendicular to the film plane (triangles). For comparison, the resistance of one of the electrodes (diamonds) is plotted in scaled units under the same conditions. Clearly vortices from the external field are well pinned in the electrodes at temperatures just below T_c, whereas the junction dissipation extends to much lower temperatures.

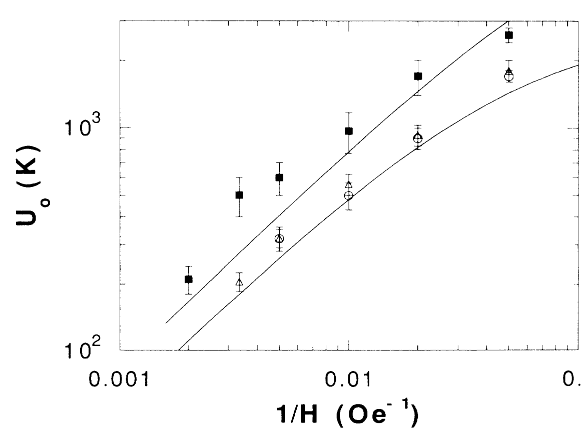

FIG. 3. The activation energy obtained from fitting data of Fig. 2 (triangles), for a Nb/Nb junction of area 8×12 μm^2, and those of two other Nb/Nb junctions, of the same area (circles) and of area 16×16 μm^2 (squares). The lines are calculated from the model described in the text: At high fields they represent an effective junction area of Φ_0/H with no adjustable parameters.

thermal fluctuations, was measured to be ~ 1.4 mA, which is about 70% of that predicted by using the measured junction resistance (e.g., from Fig. 2) in E_{cj}. However, a fit to the broadened resistive transition in zero field requires a much larger resistance in E_{cj}. This can be understood if the junction is sectioned into phase-coherent areas which are significant smaller ($A_0 \sim 1\ \mu m^2$) than the geometrical area ($\sim 96\ \mu m^2$), and may be caused by disorder in the films, which were deposited at low temperatures. Then each section, and the junction as a whole, will have the same broader transition. However, the low-temperature critical current *density* is expected to be little changed by such sectioning, so I_c, which sums the (super) conductivity of all the phase-coherent sections in parallel, should be affected very little, and the above measurements confirm this. The solid line in Fig. 3 interpolates between the low-field limit of area (A_0) and the high-field limit (Φ_0/H), using an effective area, $A = (A_0^{-1} + H/\Phi_0)^{-1}$.

A potential picture to understand the effective area being Φ_0/H has been proposed[8] to explain the c-axis resistivity in $Bi_2Sr_2CaCu_2O_x$ crystals. It borrows the concept of a coherence radius used for superconducting fluctuations in a magnetic field,[9] either above T_c or in resistive states below T_c. However, the Nb electrodes are in a zero-resistance state with the vortices pinned, and the above interpretation would miss the long-range coherence of the phase of ψ in these electrodes. A more conventional[10] approach begins with thermally activated Josephson vortex cores, which are parallel to the films in the insulating region of the junction. For finite dissipation, these must cross the entire junction area, A_0, and thus, e.g., in zero field, $E_{cj} \propto 1/R_N \propto A_0$. However, we may suggest that in a finite field, the minimum-sized Josephson core excitation is a loop of area $\sim \Phi_0/H$, since then it is possible to *connect* the Josephson cores to those of the external field which are pinned in the Nb electrodes on a lattice of that unit-cell size. It is necessary to make the unproven presumption that this reconnection makes the thermally activated loop more stable against collapse. In such a case, E_{cj} would correspond to an effective junction area, Φ_0/H. Note that the energy barrier for flux-line cutting associated with this reconnection is small: The field of the Josephson cores is weakly localized ($\sim 2\lambda_J = 24\ \mu m$) compared to the flux-line spacing ($\sim 0.4\ \mu m$ at 100 G) and the magnitude of ψ in the insulating regions is small (its value is ~ 0.007 of that in the Nb electrodes, i.e., the square root of the ratio of the critical current densities of the Josephson junction to that for depairing in Nb).

Dissipation eventually occurs when these loops further expand in size and/or merge with others to cross the entire junction area. The junction current favors both the formation and expansion of one sense of the circulation in the loops, as it does in zero field. This may be viewed as the result of a Lorentz-like force between the Josephson cores and the tunnel current. However, the tunnel current is *parallel* to the applied field, so there is no *macroscopic* Lorentz force, except in the electrodes where pinning prevents any flux motion. Further theoretical analysis is highly desirable, and it can be noted here that the pinned flux cores in the Nb may be analogous to the role of dislocations in dislocation-mediated shearing (melting) of crystals.[11]

In summary, we report the field dependence of the zero-bias resistance of discreet, high-quality, thin-film Nb Josephson junctions and find a $1/H$ dependence for the thermally activated Josephson coupling energy for H perpendicular to the film plane. This confirms the use of such a field dependence to explain[2] the flux-motion-induced broadening of resistive transitions for the HTS with H parallel to the c axis. The Nb junction results have also unambiguously confirmed another recent suggestion,[6] that field-induced dissipation can occur in Josephson junctions without the motion of vortices from an externally applied field.

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[1] D. H. Kim, K. E. Gray, R. T. Kampwirth, and D. M. McKay, Phys. Rev. B **42**, 6249 (1990).

[2] 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 (Amsterdam) **177C**, 431 (1991).

[3] P. W. Anderson (unpublished); V. Ambegaokar and A. Baratoff, Phys. Rev. Lett. **10**, 486 (1963).

[4] V. Ambegaokar and B. I. Halperin, Phys. Rev. Lett. **22**, 1364 (1969).

[5] D. H. Kim, K. E. Gray, R. T. Kampwirth, K. C. Woo, D. M. McKay, and S. J. Stein, Phys. Rev. B **41**, 11642 (1990).

[6] K. E. Gray, R. T. Kampwirth, D. J. Miller, J. M. Murduck, D. Hampshire, R. Herzog, and H. W. Weber, Physica (Amsterdam) **174C**, 340 (1991).

[7] J. H. Kang, D. L. Miller, J. X. Przybysz, and M. A. Janocko, IEEE Trans. Magn. **MAG-27**, 3117 (1991).

[8] G. Briceno, M. F. Crommie, and A. Zettl, Phys. Rev. Lett. **66**, 2164 (1991).

[9] R. Ikeda, T. Ohmi, and T. Tsuneto, J. Phys. Soc. Jpn. **58**, 1377 (1989); S. Ullah and A. T. Dorsey, Phys. Rev. B **44**, 262 (1991); M. Tinkham, *Introduction to Superconductivity* (McGraw-Hill, New York, 1975), p. 240.

[10] This approach was suggested by S. Doniach (private communication).

[11] This analogy was pointed out by V. M. Vinokur (private communication).