Anisotropy of the magnetothermoelectric power of high-T_c superconducting single crystals in the mixed state: Role of the coherence length

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The thermoelectric power S and the resistivity ρ of $Bi_2Sr_2CaCu_2O_x$ and $YBa_2Cu_3O_{7-\delta}$ single crystals in the mixed state have been measured in fields $H \perp ab$ and $H \parallel ab$. Abrikosov vortices and Josephson vortices respond in qualitatively different ways to a temperature gradient.

The motion of vortices in a type-II superconductor in its mixed state, in a direction perpendicular to the magnetic field $\vec{H}$ and to the temperature gradient $\vec{\nabla} T$, gives rise to an electric field $\vec{E}$ along $\vec{\nabla} T$. The coefficient of proportionality between $\vec{E}$ and $\vec{\nabla} T$ is thermoelectric power S. In the case of T_c superconductors in fields $H \perp ab$, the motion of vortices along $\vec{H} \times \vec{\nabla} T$ is caused primarily by the thermal emf of normal excitations, S_n (Refs. 1 and 2). The model of Ref. 1 is based on the idea that oppositely directed currents $\vec{j}_s$ and $\vec{j}_n$ are flowing at each point in the superconductor in which there is a temperature gradient. In a previous paper,[2] we pointed out that the imposition of boundary conditions at the n–s boundary gives rise to a normal current $\vec{j}_n \simeq -(S_n/\rho_n)\vec{\nabla} T$ in the core of a vortex and to a countercurrent of superconducting pairs, $\vec{j}_s \simeq -\vec{j}_n$, outside the core, if there is a temperature gradient. The interaction of $\vec{j}_s$ with the supercurrent of a vortex drives the vortex into motion. Accordingly, this model assumes that the vortices move because of factors other than those assumed in Ref. 1. Under these conditions, the magnetothermoelectric power is related to the magnetoresistance by

$$\Delta S \simeq (S_n/\rho_n)\Delta\rho. \qquad (1)$$

One might expect that in the orientation $H \parallel ab$ the electric field caused by the motion of the vortices would be determined to a large extent by the relation between the coherence length along the c axis, i.e., ξ_c, and the distance between the Cu–O planes, d. At small values of ξ_c, the order parameter is strongly modulated along the c axis, reaching a maximum near Cu–O planes and a minimum between them. Consequently, it is favorable from the energy standpoint for vortices to assume positions halfway between the planes; the result is a sort of inherent pinning.[3] Another important circumstance is that the very nature of a vortex depends on the relation between ξ_c and d. In the case $\xi_c < d/\sqrt{2}$, for example, the interaction between layers is a Josephson interaction[4] (as we demonstrated for $Bi_2Sr_2CaCu_2O_x$ single crystals in some recent experiments by Kleiner *et al.*[5]), and the distribution of the superconducting current around the center of a vortex is qualitatively different from that in the case of an Abrikosov vortex (Ref. 6, for example).

For $Bi_2Sr_2CaCu_2O_x$ (Bi2212), with a coherent length $\xi_{ab} \simeq 32$ Å in the ab plane,[7] with $d = 12$ Å, and with an anisotropy parameter[8] $\Gamma = (\xi_{ab}/\xi_c)^2 \simeq 3000$; we have $\xi_c(0)/d \simeq 0.05$. For $YBa_2Cu_3O_{7-\delta}$ (Y123), with $\xi_{ab} \simeq 20$ Å (Ref. 7), $d = 8.3$ Å, and $\Gamma \simeq 26$ (Ref. 8), we have $\xi_c(0)/d \simeq 0.5$. Consequently, at temperatures which are not very close to T_c ($T_c - T > 0.4$ K), we have $\xi_c < d/\sqrt{2}$ for Bi2212. In other words, the vortices are Josephson vortices. For Y123, in contrast, we have $\xi_c > d/\sqrt{2}$ at $T > 80$ K, and the vortices are Abrikoson vortices. Our purpose in the present study was to learn how a heat flux affects the motion of vortices in these two cases.

The dimensions of the samples were about $2 \times 1 \times 0.2$ mm in the case of Bi2212 and $1.3 \times 1.3 \times \times 0.02$ mm in the case of Y123. The Bi2212 single crystals were grown by the technique described in Ref. 9. The procedure for measuring the kinetic coefficients was similar to that described in Ref. 2. We selected a configuration of contacts such that the contribution to the voltage which was odd in the magnetic field was zero.

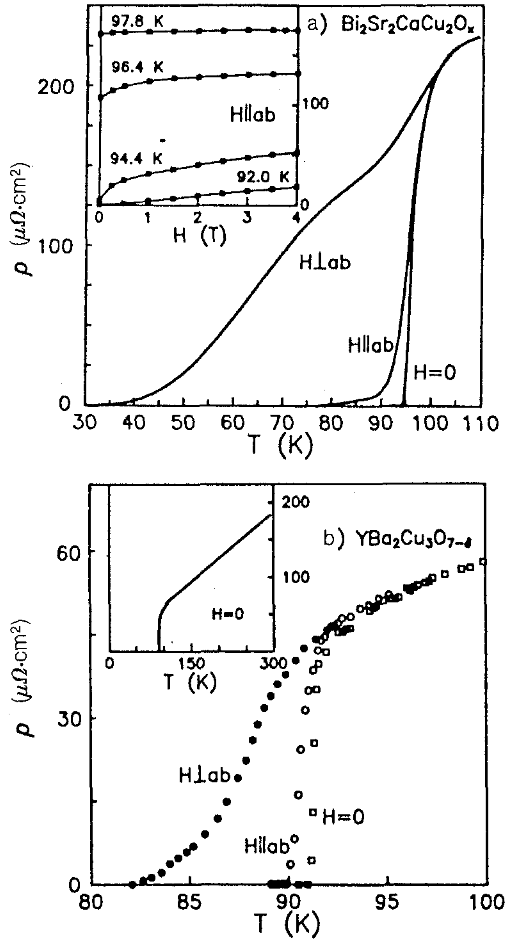

FIG. 1. a: $\rho(T)$ for Bi2212 in magnetic fields of 0 and 4 T ($H \perp ab$, $H \| ab$). The inset shows $\rho(H)$ for Bi2212 with $\Delta T = 1.5$ K and $H \| ab$ at the temperatures which are the curve labels. b: $\rho(T)$ for Y123 in magnetic fields of 0 and 4 T ($H \perp ab$, $H \| ab$). The inset shows $\rho(T)$ for Y123 in a zero magnetic field up to room temperature.

The voltage was measured with a Keithley 181 nanovoltmeter. For all the results reported here, the conditions $\vec{j}\|\vec{\nabla}T\perp\vec{H}$ and $(\vec{j}, \vec{\nabla}T)\|ab$ held. The magnetic field was imposed in the ab plane within an error of 0.5°. For Y123, the direction of the magnetic field $H\|ab$ made an angle of about 45° with twin planes. The temperature difference between the potential contacts was $\Delta T\approx 1.5$ K in the Bi2212 case and $\Delta T\approx 0.5$ K for Y123. The resistance measurements were carried out at a direct current density of 10 A/cm^2; the current direction was reversed. The thermoelectric power and the resistivity were measured as functions of the magnetic field at fixed temperatures (with a temperature difference existing). In addition, the temperature dependence of the resistivity was measured at fixed magnetic fields up to 4 T with $\Delta T=0$. The error of the measurements of the thermoelectric power was 0.02 of S_n for Bi2212 and 0.08 of S_n for Y123.

Figure 1, a and b, shows the results of measurements of the resistivity of Bi2212 and Y123 with $\Delta T=0$ in magnetic fields of 0 and 4 T (for $H\perp ab$ and $H\|ab$). The

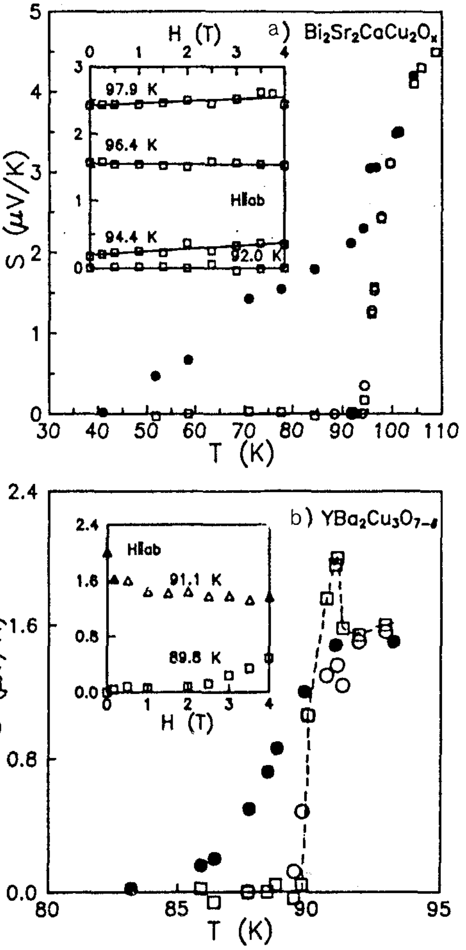

FIG. 2. a: $S(T)$ for Bi2212 in several magnetic fields. □—0 T; ●—4 T, $H\perp ab$; ○—4 T, $H\|ab$. The inset shows $S(H)$ for Bi2212 in the case $\Delta T=1.5$ K, $H\|ab$, at the temperatures which are the curve labels. b: $S(T)$ for Y123 in several magnetic fields. □—0 T; ●—4 T, $H\perp ab$; ○—4 T, $H\|ab$. The inset shows $S(H)$ for Y123 with $\Delta T=0.5$ K, $H\|ab$, at the temperatures which are the curve labels.

widths of the transitions, between the 10% and 90% levels, in a zero magnetic field are 3.5 K and 0.35 K. The midpoints of the transitions are at 97.0 K and 91.3 K for Bi2212 and Y123, respectively. The $\rho(T)$ curve for Bi2212 in its normal state is similar to that which has been reported previously,[2] while that for Y123 is shown in the inset in Fig. 1(b). The inset in Fig. 1(a) shows the resistivity versus the magnetic field $H \| ab$, measured at the same temperatures and at the ΔT values at which the thermoelectric power was measured. We see that at certain temperatures near 94 K the magnetoresistance is substantial, reaching 30% of the value of the resistance in the normal state.

Figure 2 shows the temperature dependence of the thermoelectric power. In the absence of a magnetic field, the thermoelectric power of Bi2212 (Fig. 2a) vanishes below the superconducting transition temperature. In a magnetic field $H \perp ab$, the $S(T)$ curve broadens, as the $\rho(T)$ curve does (Fig. 1a). The results in the configuration $H \| ab$ are completely different. As we see from the inset in Fig. 2a, the thermoelectric power of Bi2212 is essentially independent of the magnetic field, within the experimental error, up to 4T [an exceptional case is the $S(H)$ curve at $T = 94.4$ K, in which case the change in the thermoelectric power amounts ≈ 0.03 of S_n]. When the angle be-

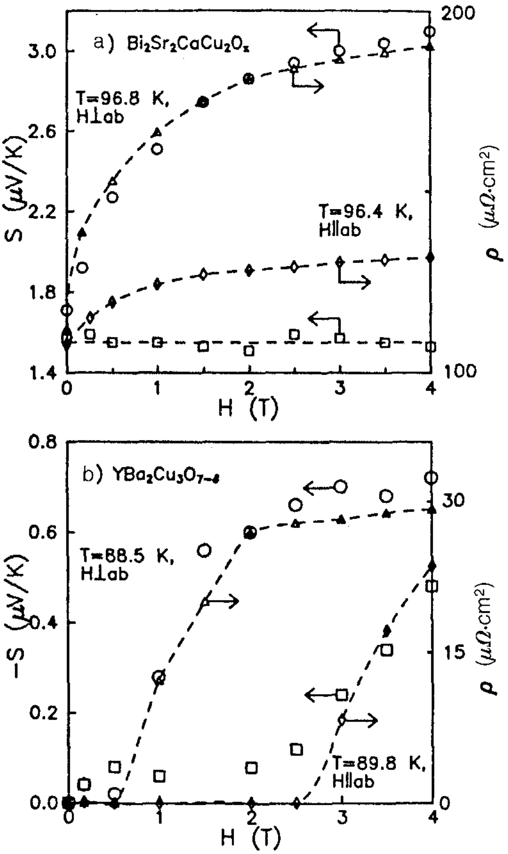

FIG. 3. a: Δ, ◇—$\rho(H)$ for $H \perp ab$ and $H \| ab$, respectively; O,□—$S(H)$ for $H \perp ab$ and $H \| ab$, respectively, for Bi2212. b; Δ, ◇—$\rho(H)$ for $H \perp ab$ and $H \| ab$, respectively; O,□—$S(H)$ for $H \perp ab$ and $H \| ab$, respectively, for Y123.

tween the magnetic field and the *ab* plane is about 9°, the change in the thermoelectric power with the magnetic field is again similar to the change in the resistivity.

A distinctive feature of the $S(T)$ curve for Y123 (Fig. 2b) is the presence of a fairly narrow peak near T_c. This peak was first discovered by Howson *et al.*[10] and was attributed to fluctuation effects.[10,11] This peak is suppressed by a magnetic field ≈ 0.2 T (in either orientation, $H \perp ab$ or $H \| ab$). The inset in Fig. 2b shows the thermoelectric power as a function of the magnetic field $H \| ab$. Near the peak the thermoelectric power decreases with increasing magnetic field, while it increases at lower temperatures.

Figure 3 shows the thermoelectric power and the resistivity versus the magnetic field for Bi2212 (Fig. 3a) and Y123 (Fig. 3b) for $H \perp ab$ and $H \| ab$. We see that the nature of the magnetic field dependence of S and ρ is the same for the two compounds in the case $H \perp ab$. It is also the same for Y1234 in the $H \| ab$. For Bi2212 with $H \| ab$, we find $\Delta S \approx 0$. The heat flux thus acts in quite different ways on Josephson vortices (Bi2212 in the case $H \| ab$) and Abrikosov vortices.

We believe that the key to reaching an understanding of this fact is the absence of a normal core and thus the absence of an *n-s* boundary in the case of a Josephson vortex. When there is a temperature gradient, the boundary conditions are responsible for the normal currents inside the core, for the motion of the vortex in the direction of $\vec{H} \times \vec{\nabla} T$, and for the energy dissipation. The small value of ξ_c and the deviation of the order parameter from zero in a small region near the Cu–O layers in the case of Bi2212 require a substantial change in our understanding of the effect, of a heat flux on a vortex structure.

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