Flux-flow resistivity and vortex viscosity of high-T_c films near T_c

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The flux-flow regime of high-T_c samples of different normal-state resistivities is studied in the temperature range where the sign of the Hall effect is reversed. The scaling of the vortex viscosity with normal-state resistivity is consistent with the Bardeen-Stephen theory. Estimates of the influence of possible mechanisms suggested for the sign reversal of the Hall effect are also given. [S0163-1829(97)03409-7]

I. INTRODUCTION

The dynamics of vortices in high-T_c materials shows a rich behavior that is not completely understood.[1] Part of the difficulties arise from the large variety of defects which pin the vortices. Because of the effects induced by pinning, it is a complicated task to separate intrinsic from extrinsic effects.

One of the most interesting topics in vortex dynamics is the anomalous sign reversal of the flux-flow Hall effect.[2,3] This anomaly cannot be understood within the usual Bardeen-Stephen model.[4] The alternative analysis of vortex dynamics put forward by Nozières and Vinen[5] cannot account for the phenomenon, either. The study of the origin of this anomaly has called the attention of many authors. One interesting point of view is to explore the relationship between the longitudinal ρ_{xx} and the transversal (Hall) ρ_{xy} resistivities. This has been done in the framework of scaling hypotheses of the vortex dynamics.[6,7] Another appealing explanation suggests that vortices may become charged,[8] as the chemical potential differs in the normal and in the superconducting phase. It has already been shown that this hypothesis can influence the surface properties of high-T_c samples.[9]

The unusual behavior of vortices which gives rise to the sign reversal in the Hall effect should also manifest itself in the longitudinal resistivity. Pinning effects complicate greatly the dynamics of vortices in high-T_c materials. In the present work, we study the longitudinal resistivity in the flux flow regime, where pinning becomes irrelevant. We analyze the temperature range where the sign reversal is observed in the Hall effect. In conventional superconductors, the flux flow longitudinal resistivity is given by the Bardeen-Stephen theory.[4] This theory predicts a proportionality between the normal-state resistivity and the flux flow resistivity. Thus, deviations from standard behavior can be deduced from the dependence of the flux flow resistivity on the normal one.

Note that, well below T_c, the coherence length is expected to become much shorter than the mean-free path, taking the materials to the ultraclean limit. In addition, quantization of the levels within the vortex core implies that an hydrodynamic description of the electrons within the cores is not possible.[10] Thus, we expect deviations from conventional behavior at low temperatures, although their origin can be unrelated to the anomaly in the Hall effect. We will concentrate on the temperature range where the sign reversal in the Hall effect can be observed.

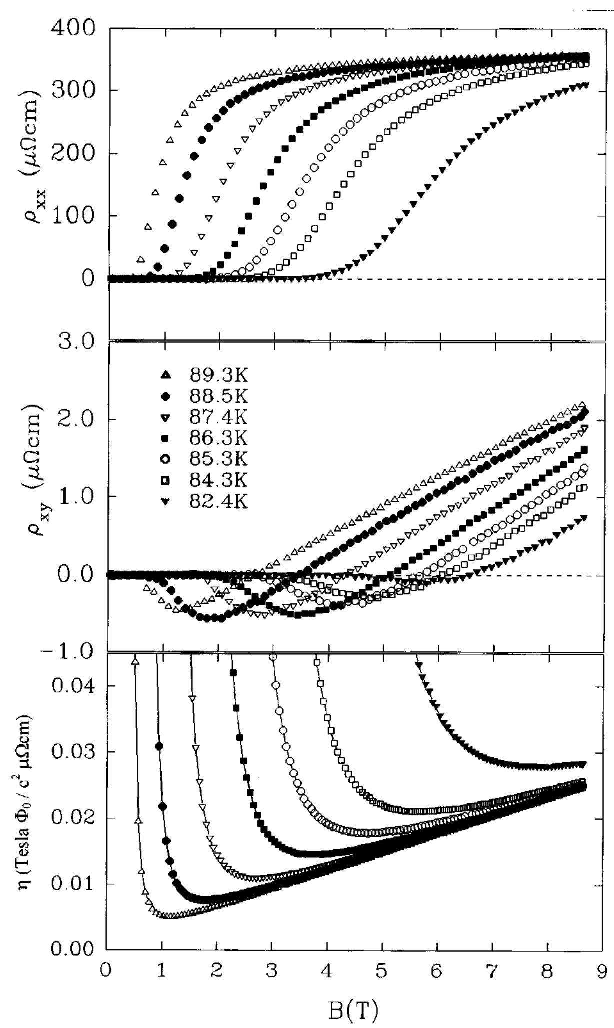

FIG. 1. Longitudinal resistivity ρ_{xx}, Hall resistivity ρ_{xy}, and vortex viscosity η, for a film with $T_c=90$ K and $\rho_N=370$ $\mu\Omega$ cm.

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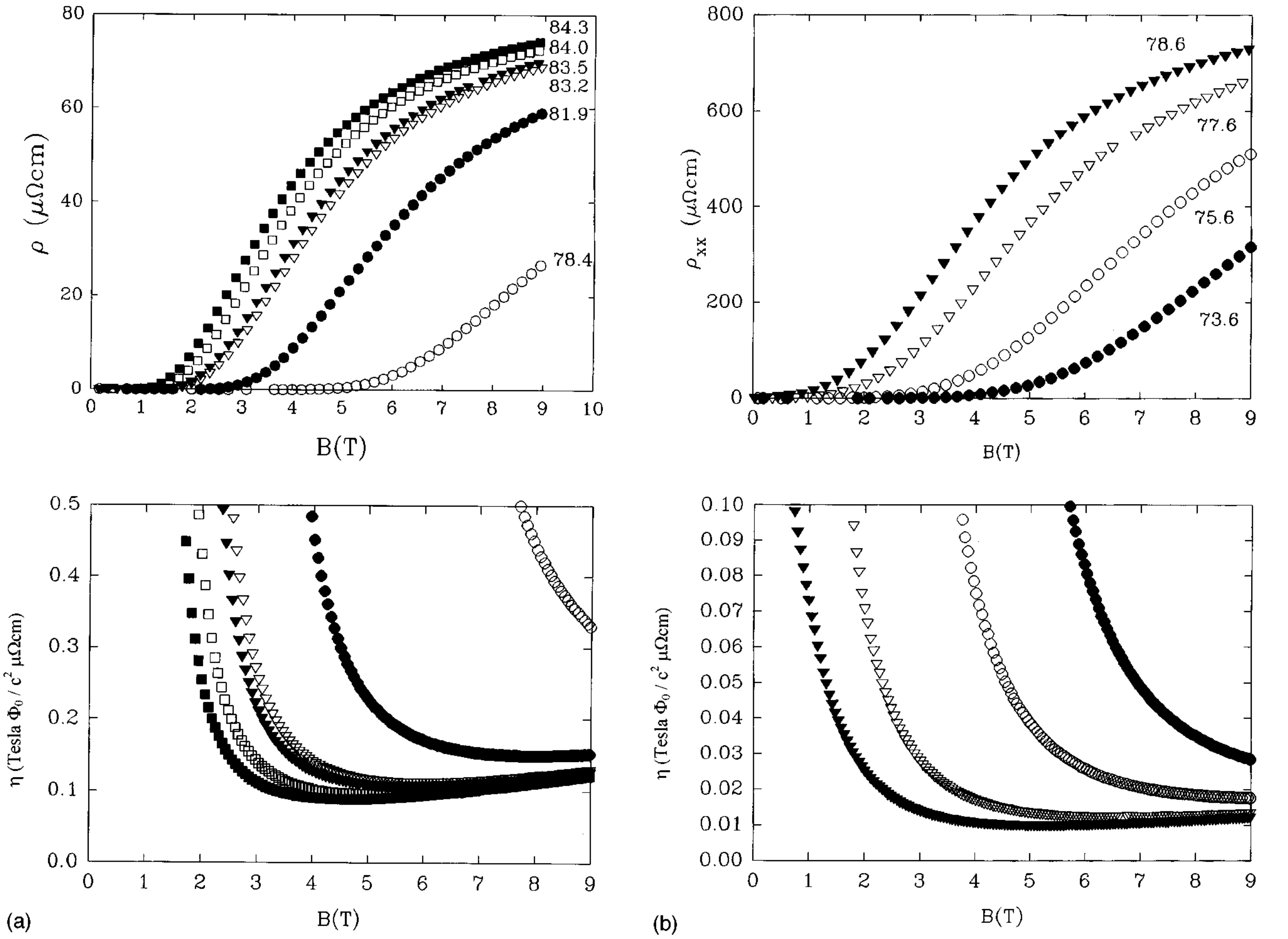

FIG. 2. Longitudinal resistivities (upper curve), and vortex viscosities (lower curves), for samples with normal resistivities $\rho_N = 76$ $\mu\Omega$ cm and $T_c = 88$ K (a), and $\rho_N = 800$ $\mu\Omega$ cm and $T_c = 80.5$ K (b).

In the next section, we describe the experimental setup. Then, we present the results. Comparison to related work is also made. In the next section, we analyze the experimental consequences of the charged vortices model. Finally, we draw the main conclusions from our results.

II. THE EXPERIMENTS

In the present work, we study the flux flow regime, near T_c, of thin films of the 1:2:3 family, with different normal-state resistivities. All of them display, in the same temperature range, a sign reversal in the Hall resistivity. Our results are consistent with those reported in Ref. 11, where two samples with normal state resistivities ~ 100 $\mu\Omega$ cm were studied. We will use also data from Ref. 11 in order to enlarge the range of normal state resistivities covered in our study.

Thin films of $EuBa_2Cu_3O_7$ have been grown on (100) $SrTiO_3$ substrates by dc magnetron sputtering, following standard procedures.[12] Samples are produced with the so-called c-axis texture (CuO_2 planes being parallel to the substrate). Stoichiometric targets were used, the substrate target geometry was on-axis, the substrate temperature was, approximately, 800 °C during deposition, the sputtering atmosphere was 85% Ar and 15% O up to a total pressure of 300 Mtorr. The annealing and cooling steps were standard, as reported in the literature.

We chose three samples with different values of T_c and normal state resistivities. Normal-state resistivities are taken 10 K above the onset. Critical temperatures, with zero resistivity, varied between 80 K and 90 K. The samples were patterned into regular bars (of width 500 μm and length 5 mm). and the transverse (Hall) and longitudinal resistivities were measured. Magnotransport effects were taken by a standard dc technique, using a commercial 90 kOe magnet and temperature controller (Lake-Shore DRC 91C). The Hall voltage was obtained from the antisymmetric part of the transverse voltage under magnetic field reversal.

III. THE RESULTS

Typical measurements are shown in Fig. 1. These results are similar to previous work[7] on the sign reversal of the Hall effect near T_c. We extract the vortex viscosity η from the field independent part of the curve. Note that, at low fields, pinning effects are important, while, at high fields, superconductivity is destroyed.

Results for the other two samples are given in Fig. 2. Our range of fields are comparable to the ones used in Ref. 11, and our results are consistent with the experiments reported

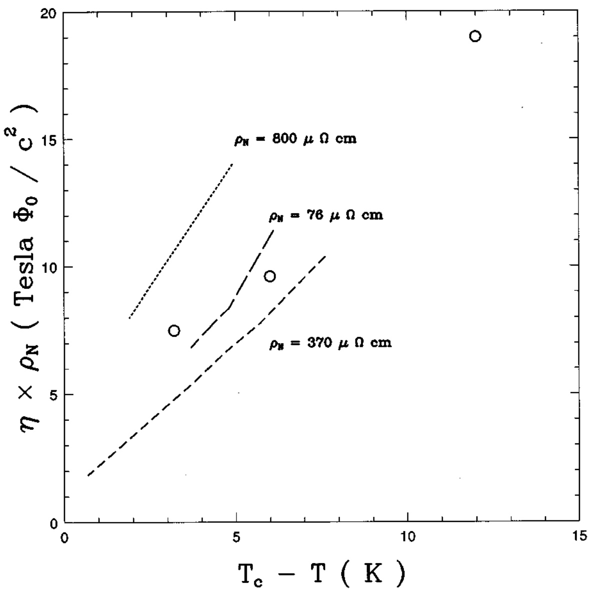

FIG. 3. Values of $\eta\rho_N$ as a function of $T_c - T$ for the three samples described in the text. Circles are data taken from Ref. 11.

there. Hence, we will use the measurements from Ref. 11 as a fourth case with different normal-state resistivity. Our results for three different samples are summarized in Fig. 3. For comparison, results from Ref. 11 are shown as circles (taken from sample II, with $T_c = 88.5$ K and $\rho_N \approx 100$ $\mu\Omega$ cm).

The results are consistent with the Bardeen-Stephen theory of flux flow viscosity. The values of $\eta\rho_N$ should scale with $H_{c_2}\Phi_0/c^2$, where H_{c_2} is the upper critical field, and Φ_0 is the flux quantum. As shown in Fig. 3, the experimental data are consistent with a linear dependence of H_{c_2} on $T - T_c$, with $dH_{c_2}/dT \sim 2T/K$. Note that no adjustements have been made on the available experimental data. The small variations in dH_{c_s}/dT suggest that the superfluid density changes little from sample to sample, as expected.

IV. INFLUENCE OF THE VORTEX CHARGE

The hypothesis that vortex cores may acquire charge, leading to the reversal of the sign in the Hall effect, is very appealing. It provides a simple explanation of the phenomenon, and it is supported by reasonable numerical estimates. In the following we analyze the expected effect of this hypothesis on the longitudinal resistivity.

In order to compute the contribution of the core charge to the viscosity of a single vortex, we consider first an isolated pancake vortex, localized in a single CuO_2 plane. This point particle, as it moves under the influence of a voltage, creates excitations in the medium, and dissipates energy. Assuming that the leading cause of disspation is the creation of electron-hole pairs, we can write the energy loss per unit time as

$$\frac{\partial E}{\partial t} = \eta|\vec{v}|^2 = \int d^D q V_q^2 \mathrm{Im}\chi(\vec{q},\vec{q}\vec{v}), \tag{1}$$

where $\chi(\vec{q},\omega)$ is the polarizability due to electron-hole pairs of the medium, and V_q is the coupling between the core charge and the electron-hole pairs. Neglecting for the moment the influence of the superconducting gap, we know that, in a metal, Im $\chi(\vec{q},\omega) \propto \omega/\epsilon_F^2$, where ϵ_F is the width of the conduction band of the metal. From Eq. (1), we can infer the value of the viscosity, η. In general, for short range potentials, Eq. (1) leads to $\eta \sim (\hbar V^2)/d(\epsilon_F^2 a^2)$, where V is the potential induced by the ''impurity'' on the metal, and a is its range, and d is the separation between planes. This expression gives the viscosity per unit length of the vortex. Alternatively, we can replace V/ϵ_F by δ, the phase shift induced by the potential on the electrons at the Fermi level. For the charged vortex considered here, δ should scale with the charge of the vortex, in dimensionless units. Hence, $V/\epsilon_F \sim \delta \sim Q/e$. The smallness of Q justifies, *a posteriori*, the use of second order perturbation theory in the present analysis of the dissipation. The range of the potential goes like the size of the core, that is, the coherence length, ξ. Finally, the vortex viscosity per unit length is

$$\eta_Q \sim \frac{\hbar Q^2}{e^2 \xi^2 d}. \tag{2}$$

The standard theory of the stopping power of charges in metals[13–15] gives a larger value for the vortex viscosity per unit length, $\eta_Q \sim 0.1Q^2/d$ (in atomic units), for typical metallic densities. The main reason for this difference lies in the size of the potential due to the core charge, which is taken to be of the order of the inverse Fermi-Thomas wave vector, k_{FT}^{-1} in the second case. A complete elucidation of this question requires a detailed knowledge of the screening processes near the vortex core.[9] The value of η_Q is to be compared to the Bardeen-Stephen contribution

$$\eta_{\mathrm{BS}} \sim \frac{\Phi_0 B_{c_2}}{\rho_N c^2}, \tag{3}$$

where Φ_O is the quantum unit of magnetic flux, B_{c_2} is the upper critical field, ρ_n is the normal-state resistivity, and c is the velocity of light.

We can write $\Phi_0 B_{c_s} \sim B_{c_2}^2 \xi^2$ as $\Delta F \xi^2$, where ΔF is the condensation energy per coherence length to the cube, so that $\Phi_0 B_{c_2} \sim \Delta^2/(\epsilon_F d)$, where Δ is the superconducting gap. Using this last expression, and the value of η_Q given in (2), we obtain

$$\frac{\eta_Q}{\eta_{BS}} \sim \frac{\rho_N}{[(\Delta^2 e^2 \xi^2)/(\hbar Q^2 \epsilon_F c^2)]}. \tag{4}$$

We estimate the denominator in Eq. (4) assuming that $Q = 10^{-3}e, \epsilon_F = 1$ eV, $\Delta = 0.05$ eV, and $\xi = 50$ Å. Then $\rho_Q = (\Delta^2 e^2 \xi^2)/(\hbar Q^2 \epsilon_F c^2) \approx 10^3$ $\mu\Omega$ cm. If we use the standard expression for the stopping power of a charge moving within a metal, we obtain $\rho_Q \sim 10$ $\mu\Omega$ cm. The relative importance of the vortex charge in the flux flow dissipation can be inferred by comparing the value above to the normal-state resistivity of the material under consideration. The effect of the vortex charge will be important if $\rho_Q < \rho_N$. The two estimates given above can be considered as an upper and a lower bound, so that $10\ \mu\Omega\ \mathrm{cm} < \rho_Q < 10^3\ \mu\Omega\ \mathrm{cm}$. The samples studied here have normal-state resistivities within this range.

We now consider the possible sources of error in the derivation of the estimate of ρ_Q given earlier. We assume that the response of the material is that of a gapless metal. This is justified as far as $\Delta \ll k_B T$, that is, near T_c. At lower temperatures, electron-hole pairs cannot be excited at low energies, and the dissipation is reduced. The other main approximation made in estimating ρ_Q lies in the neglect of the temperature dependence of Q, Δ and ξ. Note, however, that the product $\Delta\xi$ which enters in ρ_Q is independent of temperature. Q goes to zero as $T \to T_c$, reducing the relative importance of η_Q near the transition temperature. On general grounds, $Q \sim \Delta^2(T) \sim T_c - T$,[8] and $\rho_Q \sim (T_c - T)^2$, which is not consistent with the results shown in Fig. 3. Outside the critical region, our estimate $Q \sim 10^{-3} e$ per plane[8] is probably too conservative.[9] In any case, the value of Q is the most uncertain parameter in ρ_Q.

V. CONCLUSIONS

In conclusion, we have analyzed the flux-flow regime of high-T_c samples in the range where the sign reversal of the Hall effect is observed. Samples with different normal state resistivities were used, in order to verify the validity of the standard Bardeen-Stephen theory of flux flow dissipation. Our results are consistent with this theory, with $dH_{c_2}/dT \approx 2T/K$. Estimates of the expected deviations associated to the charging of the vortices suggest that this effect should influence dissipation in samples with normal state resistivities similar to those of the samples studied here. We find, however, no significant deviations from the Bardeen-Stephen theory within our experimental errors.

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[1] G. Blatter, M. V. Feigel'man, V. B. Geshkenbein, A. I. Larkin, and I. M. Vinokur, Rev. Mod. Phys. **66**, 1125 (1994).
[2] S. J. Hagen, A. W. Smith, M. Rajeswari, J. L. Peng, Z. Y. Li, R. L. Greene, S. N. Mao, X. X. Xi, S. Bhattacharya, Qi Li, and C. J. Lobb, Phys. Rev. B **47**, 1064 (1993).
[3] J. M. Graybeal, J. Luo, and W. R. White, Phys. Rev. B **49**, 12 923 (1994).
[4] J. Bardeen and M. J. Stephen, Phys. Rev. **140**, 1197 (1965).
[5] P. Nozières and W. F. Vinen, Philos. Mag. **14**, 667 (1966).
[6] V. M. Vinokur, V. B. Geshkenbein, M. V. Feigel'man, and G. Blatter, Phys. Rev Lett. **71**, 1242 (1993).
[7] J. I. Martin, M. Velez, J. Colino, P. Prieto, and J. L. Vicent, Solid State Commun. **94**, 341 (1995).
[8] D. I. Khomskii and A. Freimuth, Phys. Rev. Lett. **75**, 1384 (1995).
[9] G. Blatter, M. Feigel'man, V. Geshkenbein, A. Larkin, and A. van Otterlo, Phys. Rev. Lett. **77**, 566 (1966).
[10] F. Guinea and Yu. Pogorelov, Phys. Rev. Lett. **74**, 462 (1995).
[11] M. N. Kunchur, D. K. Christen, and J. M. Phillips, Phys. Rev. Lett. **70**, 998 (1993).
[12] J. Colino, J. L. Sacedón, and J. L. Vicent, Appl. Phys. Lett. **59**, 3327 (1991).
[13] A. van Otterlo, M. Feigel'man, V. Geshkenbein, and G. Blatter, Phys. Rev. Lett. **75**, 3736 (1995).
[14] J. Lindhard and A. Winther, K. Dan. Vidensk. Selsk. Mat. Fys. Medd. **34**, 4 (1964).
[15] P. M. Echenique, R. M. Nieminen, J. C. Ashley, and R. H. Ritchie, Phys. Rev. A **33**, 897 (1986).