TRANSPORT BEHAVIOR IN SUPERCONDUCTORS AT EXTREME DISSIPATION LEVELS

M. N. Kunchur[1,2], D. K. Christen[3] and B. I. Ivlev[4]

[1] *WL/POOX-2, Bldg. 450, Area B, Wright Patterson Air Force Base, OH 45433*
[2] *Dept. of Physics and Astronomy, University of South Carolina, Columbia, SC 29208*
[3] *Solid State Division, Oak Ridge National Lab, P. O. Box 2008, Oak Ridge, TN 37831*
[4] *Landau Institute for Theoretical Physics, Kosygin St. 2, 117940 Moscow, Russia.*

A number of fundamental physical phenomena unfold in the mixed state of superconductors, when subjected to enormous current and power-dissipation levels. A sufficiently large current can destroy the superconducting state itself—the so-called pair-breaking effect. At intermediate current densities, below the onset of pair-breaking, one expects to see the free viscous flow of flux vortices. In the present work a pulsed-current technique was used to explore this dissipative regime of high-T_c superconductors, verifying both free flux flow and the pair-breaking effect, as predicted by traditional theories. This paper concentrates on the dissipation and Hall behavior in the free flux flow state.

In this work, we have employed short-duration ($\sim 1\mu s$) current pulses to explore superconducting transport behavior at current levels above the onset of dissipation J_c and beyond the depairing limit J_d[1]. It was shown[2] that by applying a sufficiently large Lorentz driving force on flux vortices, and driving their motion into the state of free flux flow, simple Bardeen-Stephen[3] behavior could be verified for the first time in any superconductor: the transport characteristic becomes Ohmic (total resistivity $\rho = E/J$ is current independent) and the dissipation obeys:

$$\rho_{FFF}/\rho_n \approx B/H_{c2}, \tag{1}$$

where ρ_n and H_{c2} are the normal-state resistivity and the upper-critical field respectively *for that temperature*. Fig. 1 shows longitudinal resisitivity in free flux flow for two YBCO samples. Recently, similar saturation of $\rho(J)$ leading to an Ohmic plateau was also observed in the low-T_c low-pinning materials 2H-$NbSe_2$ by Higgins et al.[4] and in $Mo_{77}Ge_{23}$ by Hellerqvist et al.[5].

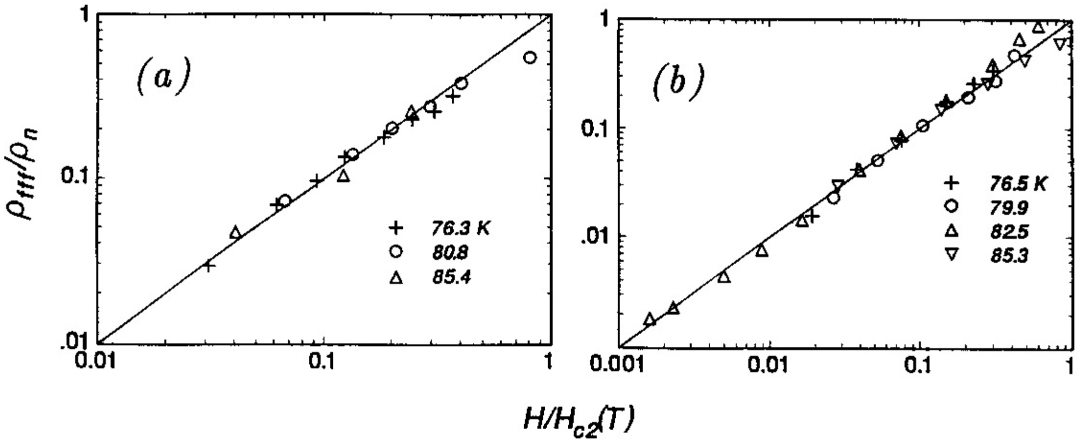

Figure 1: Normalized free-flux-flow resistivity versus H for two YBCO specimens.

At even higher current densities we were able to observe current-induced pair-breaking[6], where T_c was found to shift as:

$$\Delta T_c(H,J)/T_c(H,0) = [J/J_0]^{\frac{2}{3}}. \tag{2}$$

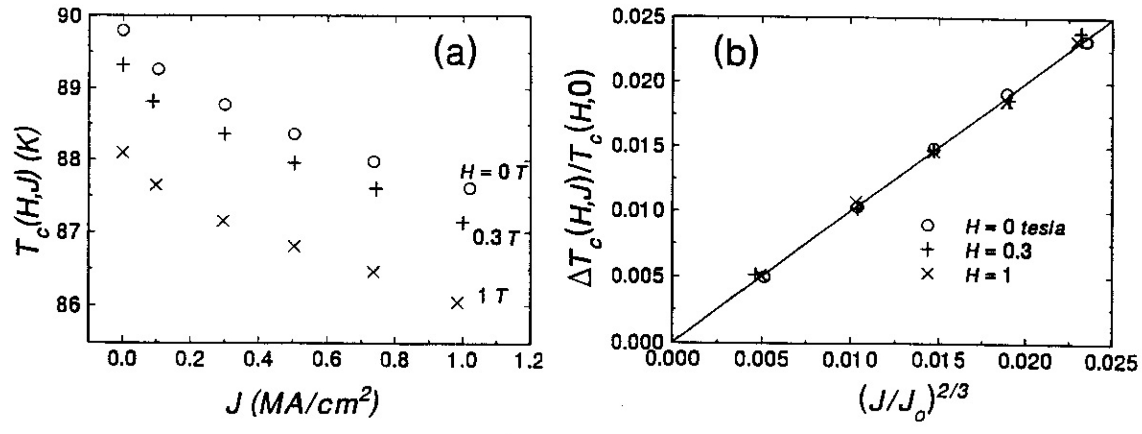

Figure 2: (a) Mid-transition T_c's at different H and J values for a YBCO film. (b) Normalized shifts in T_c as a function of $(J/J_0)^{2/3}$. $J_0 \approx 4 \times 10^8$ A/cm^2 is a field-independent constant.

This is shown in Fig. 2.

The mixed-state Hall effect has been a topic of great current interest and one surrounded by controversy. Central to the controversy is the observed sign reversal of the tangent of the Hall angle $\tan\alpha = \rho_{xy}/\rho_{xx}$, where ρ_{xx} and ρ_{xy} are the longitudinal and transverse (Hall) resistivities respectively[1]. One of the main questions that has plagued the interpretation of the Hall effect is whether the sign reversal is caused by pinning[7]. Our previous work on dissipation and free flux flow had shown the effectiveness of high currents in suppressing the effect of pinning. We were able to apply this same principle to investigate the role of pinning in the Hall effect. Fig. 3 shows the Hall angle versus temperature at different fixed magnetic fields[8]. The upper panel (a) was measured at low continuous currents (5 kA/cm^2), the lower panel at much higher pulsed currents ($\sim$ 1 MA/cm^2). As is evident, there is a dramatic enhancement of the sign reversal showing that pinning hinders the sign reversal rather than causes it. The data in Fig. 3 was one of the earliest and most decisive proofs that the sign reversal is not caused by pinning. Recently, Kang et al.[9] were able to show qualitatively similar behavior by varying the pinning by introducing artificial defects into their samples. An advantage of overcoming pinning using a high current as opposed to varying the defect concentration is that the process is reversible and continuously variable. Also there is little change in other parameters, such as the carrier concentration, which can however be dependent on defect concentration.

In addition to the sign reversal, there is another feature that has been observed by us and other groups[10, 11, 12, 8]. This is the decomposition of the Hall angle, α, into a field-proportional component α_n, and a field-independent component α_M. The field-proportional component arises from the Hall effect of the normal core of vortices, whereas the field-independent component α_M, can arise from the hydrodynamic (Magnus) force on the "body"of the vortex. The latter can be treated through the time-dependent-Ginzburg-Landau (TDGL) theory of the dynamics of the superconducting order parameter (as shown in more detail below). α_M can be estimated by subtracting from the total Hall angle, the normal-core component $\alpha_n(T)$ found by fitting to its $\sim H/T^2$ behavior above T_c and extrapolating this to temperatures below T_c. Fig. 4 shows α_M for our data measured on a YBCO film for 3 different fields[8]. Notice that as the temperature is lowered below T_c, as long as pinning is not influential, α_M seems to saturate to a value that does not depend on the field or current (As long as these are large enough to ensure free flux flow.)

[1]Since the Hall angle is typically small, we will not always distinguish between the angle and its tangent in the remainder of this discussion.

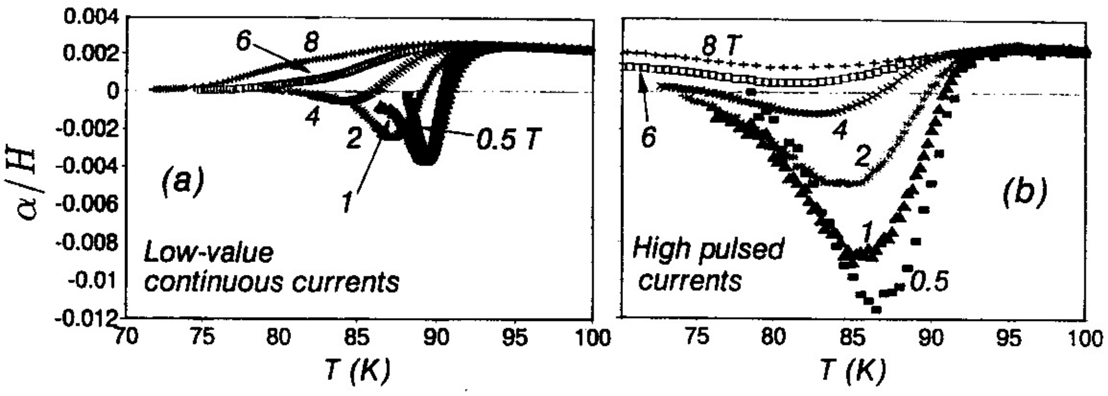

Figure 3: (a) Temperature dependence of α/H for indicated fixed fields measured at a continuous current of J=5 kA/cm^2 in a YBCO film. (b) Similar data measured with pulsed current at $J = 0.7 MA/cm^2$ (H = 0.5, 1, 2, and 4 T), $J = 1.1 MA/cm^2$ ($H = 6T$) and $J = 1.5 MA/cm^2$ ($H = 8T$).

Moreover the saturation or maximum negative value (about 0.017 in our Fig. 4) is about the same order as that reported by other groups in various materials: 0.005–0.007 for $Tl_2Ba_2CaCu_2O_8$ (references [10] and [11]) and 0.009–0.017 for YBCO (references [12] and [8]).

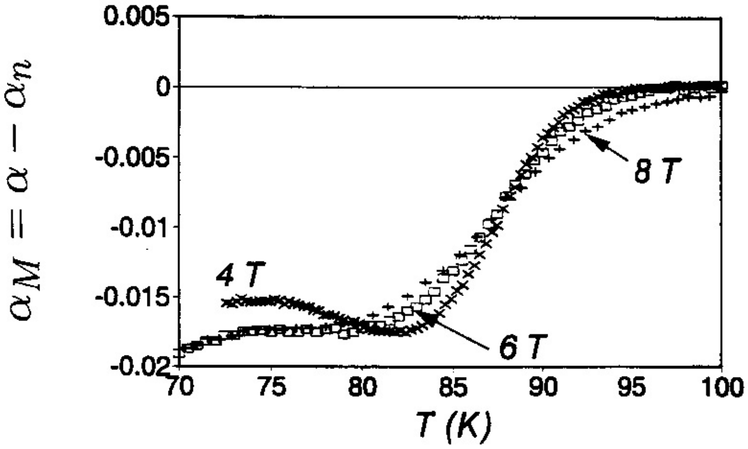

Figure 4: Magnus-force component of the high-field Hall angle $\alpha_M = \alpha - \alpha_n$, for the data of Fig. 3. Here, $\alpha_n = 21.9 \times H/T^2$ is the observed T- and H-dependent normal-state contribution.

We now show that the decomposition $\tan\alpha \approx \tan\alpha_n + \tan\alpha_M$ follows from the dynamics of the superconducting order parameter as calculated in the framework of time-dependent-Ginzburg-Landau (TDGL) theory[13]. Moreover the magnitude of $\tan\alpha_M$ obtained in this way (0.014) is in rough accord with the experimental observations ($\sim$ 0.005–0.02).

The vortex viscosity η, that governs the longitudinal conductivity consists of two parts: the Bardeen-Stephen component, which arises from dissipation of normal carriers inside the vortex core, and the Tinkham component, which arises from the dynamics of the order parameter. The net longitudinal conductivity in free flux flow is [13]

$$\sigma \approx \sigma_n \left(\frac{\alpha u}{2}\right)\left(\frac{H_{c2}}{B}\right), \tag{3}$$

where σ_n is the normal-state conductivity. Similarly, the transverse conductivity can be written as

$$\sigma^H \approx C\sigma_n^H \left(\frac{\alpha u}{2}\right)\left(\frac{H_{c2}}{B}\right) + sign(e)\sigma_n \left(\frac{\zeta\beta u}{2}\right)\left(\frac{H_{c2}}{B}\right), \tag{4}$$

where C is a constant of order unity. The first term arises from the Hall effect of normal carriers in the core, and the second term is the contribution from order-parameter dynamics. There is no contribution from quasiparticles embedded in the superfluid outside the core at zero frequency. $\alpha \sim 0.0.5$, $\beta \sim 0.27$, and $\zeta \approx T_c/E_F \sim 90/4000$ are constants of the TDGL theory[13] (E_F was taken from Wolf and Kresin [14].) The constant α should be distinguished from the Hall angle α.

From Eqs. 3 and 4 one then gets

$$\tan\alpha = \frac{\sigma^H}{\sigma} = C\frac{\sigma_n^H}{\sigma_n} + sign(e)\frac{\zeta\beta}{\alpha} \sim \alpha_n + 0.014. \tag{5}$$

This expression shows a Hall angle that decomposes into a field-proportional "normal-core" component α_n and a field-independent term $\alpha_M \sim 0.014$, which is comparable to the experimental observations ($\sim$ 0.005–0.02).

We'd like to point out that in a recent Letter[15] Khomskii and Freimuth present an interesting mechanism based on charging of the vortex core that explains the sign-reversing component of the mixed-state Hall effect and also predicts the qualitative decomposition $\tan\alpha \sim \tan\alpha_n + \tan\alpha_M$. Their predicted estimate of α_M (which they call α_q), about 10^{-4}, is however lower than the experimental results by about a factor of 100.

To summarize, we have reviewed our past work that explored the transport behavior of superconductors at current densities well above the critical value, where flux motion converts to free flow and eventually leads into the current-induced pair-breaking regime. In addition we present results on the mixed-state Hall effect, where high J was once again used to decisively rule out pinning as the cause of the sign reversal. In the free-flux-flow limit, we find that the Hall angle shows a decomposition $\tan\alpha \sim \tan\alpha_n + \tan\alpha_M$, also seen in other superconductors. An explanation, based on TDGL theory, describes the qualitative features of the data and also shows rough quantitative agreement.

Some of this research was sponsored by the Division of Materials Sciences, U. S. Department of Energy under Contract No. DE-AC05-84OR21400, with Martin-Marietta Energy Systems, Inc. Partial support for the program was administered by the Oak Ridge Institute for Science and Education.

References

[1] M. N. Kunchur, Mod. Phys. Lett. B **9**, 399 (1995).

[2] M. N. Kunchur, D. K. Christen, and J. M. Phillips, Phys. Rev. Lett. **70**, 998 (1993).

[3] J. Bardeen and M. J. Stephen, Phys. Rev. **140**, A1197 (1965).

[4] M.J Higgins, S. Bhattacharya, and T. V. Ramakrishnan, *Bull. Am. Phys. Soc.* **39**, 730 (1994).

[5] Hellerqvist et al., Phys. Rev. Lett. **76**, 4022 (1996).

[6] M. N. Kunchur *et al.*, Phys. Rev. Lett. **72**, 752 (1994).

[7] Z.D. Wang and T.S. Ting, Phys. Rev. B **46**, 284 (1992); Phys. Rev. Lett. **67**, 3618 (1991).

[8] M. N. Kunchur, D. K. Christen, C. E. Klabunde, and J. M. Phillips, Phys. Rev. Lett. **72**, 2259 (1994).

[9] Kang et al., Phys. Rev. Lett. **76**, 2993 (1995).

[10] S. J. Hagen, C. J. Lobb, R. L. Greene, and M. Eddy Phys. Rev. B **43**, 6246 (1991).

[11] A. V. Samoilov, Z. G. Ivanov, and L.-G. Johansson, Phys. Rev. B bf 49, 3667 (1994).

[12] A. V. Samoilov, Phys. Rev. B **49**, 1246 (1994).

[13] N. B. Kopnin, B. I. Ivlev, and V. A. Kalatsky, J. Low Temp. Phys. **90**, 1 (1993).

[14] S. A. Wolf and V. Z. Kresin, IEEE Trans. on Mag. **27**, 852 (1991).

[15] D. I. Khomskii and A. Freimuth, Phys. Rev. Lett. **75**, 1384 (1995).