SQUID Milliattovoltometry of $YBa_2Cu_3O_{7-x}$ Thin Films: Dissipation in Low Magnetic Fields

F. C. Wellstood
Center for Superconductivity Research, Department of Physics, University of Maryland, College Park, Maryland 20742-4111

M. J. Ferrari
AT&T Bell Laboratories, Murray Hill, New Jersey 07974

J. J. Kingston, T. J. Shaw, and John Clarke
Department of Physics, University of California, and Materials Sciences Division, Lawrence Berkeley Laboratory, University of California, Berkeley, California 94720
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The observation of magnetic flux noise in superconducting thin-film flux transformers of $YBa_2Cu_3O_{7-x}$ enables us to measure a very small, frequency-dependent resistance $R_v(f)$ in the films, caused by vortex motion. In static magnetic fields below 10^{-6} T and at voltages down to 10^{-20} V Hz$^{-1/2}$, R_v is proportional to the frequency f at which the measurement is made, is approximately independent of temperature, and decreases with increasing static current in the transformer. This behavior cannot be explained by certain currently accepted models of dissipation arising from vortex motion, but is consistent with a model in which independent vortices hop between two potential wells in a confined region.

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There has been continuing debate about the nature of the high-field vortex state in high-transition-temperature (T_c) superconductors such as $YBa_2Cu_3O_{7-x}$ (YBCO) and in particular over the existence of a novel vortex-glass state [1]. The properties of and phase transitions into the vortex state have been probed with transport [2–6], magnetization [7–11], mechanical oscillator [12], and magnetic decoration techniques [13–16]. Of these, the transport measurements have shown the clearest evidence for a vortex-glass state. Recently, Gammel, Schneemeyer, and Bishop [4] extended transport measurements down to the nV/m electric field range and Sandvold and Rossel [11] used the magnetization decay of thin-film rings to infer characteristics at the fV/m level. In this Letter, we report an analysis of magnetic flux noise data in YBCO thin films at low magnetic fields that yields the resistance at voltages in the milliattovolt regime, corresponding to electric fields below 10^{-18} V Hz$^{-1/2}$m^{-1}. The measured resistance cannot be explained by the dissipation mechanisms associated with the vortex-glass model [1], thermally activated flux creep [17], or collective creep [18], implying that our experiments are in a hitherto inaccessible regime. We explain our data in terms of a vortex-hopping model in which each vortex hops independently between two pinning sites in a spatially confined region.

We have previously [19–21] used a dc superconducting quantum interference device (SQUID) to measure flux motion in YBCO thin-film flux transformers. Our flux transformer is a closed superconducting loop consisting of a relatively large area, single-turn pickup loop of inductance L_p and a smaller area, multiturn input coil of inductance L_i coupled to the SQUID via a mutual inductance M_i (Fig. 1). When a magnetic flux is applied to the pickup loop, a persistent current is induced in the transformer and a fraction of the flux is coupled to the SQUID. The predominant noise source is the hopping of vortices between pinning sites in the input coil [19]. Because of flux quantization, this vortex motion generates a fluctuating current $I(t)$ in the transformer which, in turn, produces a flux noise $\Phi(t) = M_i I(t)$ in the SQUID. The SQUID is operated in a flux locked loop that produces an output proportional to $\Phi(t)$. The transformer is cooled in an ambient magnetic field of less than 10^{-6} T. Typically [19–21], the measured flux noise has a spectral density $S_\Phi(f,T)$ that scales as T/f, where f is the frequency. When we apply a static magnetic field $|B| \lesssim 5\ \mu$T at tem-

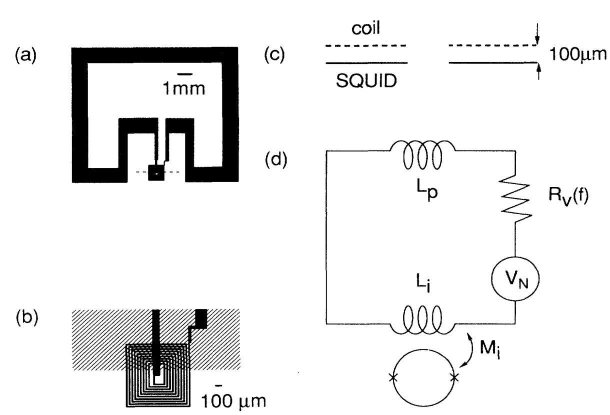

FIG. 1. YBCO flux transformer inductively coupled to low-T_c planar dc SQUID. (a) Plan view; (b) expanded view of 10-turn input coil; (c) section through dashed line in (a); (d) lumped circuit model including resistance R_v, due to vortex motion, and its associated Nyquist noise voltage $V_N(t)$.

peratures $T < T_c$, a persistent current is induced in the transformer that completely screens the applied flux. In this way, we can measure the current dependence of the noise. In Fig. 2 we show the current dependence of S_Φ(1 Hz) at two temperatures; the noise is suppressed symmetrically as the static current is increased from zero.

Unlike prior static measurements of resistivity or flux creep, which are sensitive only to net transport of vortices, our flux noise measurements detect vortex motion in thermal equilibrium (this equilibrium may be metastable because the system is in a state of quenched disorder) [1]. Thus, according to the fluctuation-dissipation theorem, [22] *the flux noise must be associated with dissipation which couples to the measured variable.* In our experiment this variable is the current in the transformer, so that the dissipation is an electrical resistance R_v produced by the vortex motion. Since we have found that cutting the flux transformer causes a large reduction in the noise [19], the resistance is in series with the two inductances [Fig. 1(c)]. According to the Nyquist theorem [23], the spectral density of the voltage noise in the flux transformer is $S_v(f) = 4k_BTR_v$. This voltage noise produces a flux noise in the SQUID with spectral density

$$S_\Phi(f) = 4k_BTR_vM_i^2/[R_v^2 + \omega^2(L_i+L_p)^2] , \qquad (1)$$

where $\omega \equiv 2\pi f$. Experimentally, we find $S_\Phi(f,T) = C_0T/f$, where C_0 is only weakly dependent on temperature and frequency. Our previous work [20,24] analyzed these weak dependences, but we may safely ignore them here. Solving Eq. (1) for R_v in the physically relevant limit $R_v \ll \omega(L_i+L_p)$ yields a frequency-dependent resistance

$$R_v(f) = \pi^2 fC_0(L_i+L_p)^2/k_BM_i^2 . \qquad (2)$$

This equation shows (i) $R_v(f) \propto f$, (ii) $R_v(f) \propto C_0$, and (iii) $R_v(f)$ is approximately independent of T [25]. We note that $R_v(0) = 0$, so that the transformer can sustain a static supercurrent; from the observed upper limit on the decay of this current during the time scale of our experiments ($> 10^3$ s), we find $R_v(0) \leq 2\times10^{-16}$ Ω. Using the experimental values $L_i \approx 75$ nH, $L_p \approx 20$ nH, $M_i \approx 3$ nH, and $C_0 \approx 10^{-8}\,\Phi_0^2/\mathrm{K}$ in units of the flux quantum $\Phi_0 = h/2e$, we find $R_v(1\ \mathrm{Hz}) \approx 30$ pΩ; for the geometry of our input coil, we infer a resistivity ρ_v(1 Hz) of less than 10^{-20} Ω m. The frequency dependence of $R_v(f)$ and $\rho_v(f)$ is plotted in Fig. 3(a), and R_v(1 Hz) and ρ_v(1 Hz) are plotted versus temperature in Fig. 3(b). Although there is considerable scatter in the data, we see that R_v(1 Hz) does not increase markedly with temperature. Since the magnitude of the noise, and hence C_0, can be strongly sample dependent [24], we expect $R_v(f)$ to depend on sample microstructure. In addition, the noise measurements in Fig. 2 imply that $\rho_v(f)$ and $R_v(f)$ decrease with increasing current.

Our measurement of $R_v(f)$ involves electric fields which are generated thermally and are orders of magnitude smaller than those employed by previous workers. At 50 K and 1 Hz the rms noise current in the flux transformer is approximately $(4k_BTR_v)^{1/2}/\omega(L_i+L_p) \approx 5\times10^{-10}\ \mathrm{A\,Hz^{-1/2}}$, corresponding to a current density of $80\ \mathrm{A\,Hz^{-1/2}\,m^{-2}}$ in the $20\times0.3\ \mu\mathrm{m}^2$ films of the input coil. The fluctuating voltage induced across R_v [Fig. 1(c)] by this current is thus $1.5\times10^{-20}\ \mathrm{V\,Hz^{-1/2}}$, corresponding to an electric field of about $6\times10^{-19}\ \mathrm{V\,Hz^{-1/2}\,m^{-1}}$, where we have used the total length of the input coil, 24 mm.

We now briefly compare our measured dissipation with the predictions of thermally activated flux creep [17] and of the collective flux creep [18] and vortex glass models, in which dissipation arises from the nucleation and growth of vortex rings [1,26]. None of these mechanisms

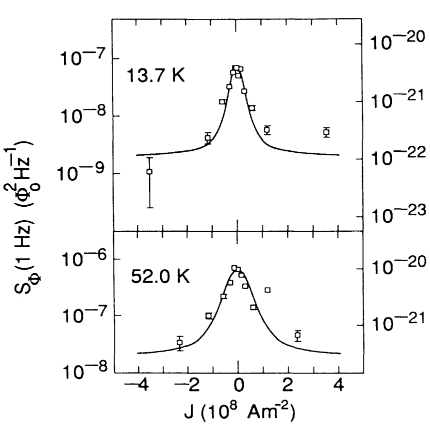

FIG. 2. Flux noise S_Φ(1 Hz) and corresponding resistivity ρ_v(1 Hz) vs static current density in input coil of transformer 1 at two temperatures. Points are experimental data; solid lines, fit to theory (Ref. [19]).

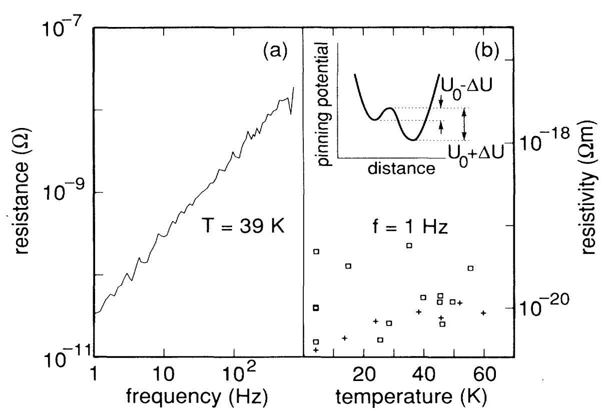

FIG. 3. (a) Vortex resistance $R_v(f)$ and resistivity $\rho_v(f)$ vs frequency at 39 K, computed from Eq. (2) and noise data in Ref. [19]. (b) R_v(1 Hz) and ρ_v(1 Hz) vs temperature for flux transformers 1 (crosses) and 2 (open squares) (Ref. [18]). Inset: Two-state vortex-hopping model in the presence of a static current with $\Delta U = I\Phi_0\bar{l}/2w$.

produces dissipation that is consistent with our data, which show that $R_v(f)$ scales with frequency, is approximately independent of temperature, and decreases with increasing current. For example, in flux creep, the resistance is frequency independent, falls exponentially as the temperature is lowered, and increases with increasing current. In the case of the nucleation and growth of vortex rings, the resistance is predicted to be proportional to [1,18] $\exp[-(J_T/J)^\mu]$, where J is the current density, $J_T \propto 1/k_B T$ is a characteristic current density, and μ is a model-dependent exponent. For $T \neq 0$, the resistance increases as the current is increased, whereas our data show that $R_v(f)$ *decreases* with increasing static current. Furthermore, the resistance is predicted to vanish only in the limit of zero temperature or zero current, whereas Eq. (2) implies that $R_v(0)$ is strictly zero even for nonzero temperature and current. We note that in the vortex ring picture, the characteristic size of the ring and hence the energy barrier determining the dissipation scales inversely as a (model dependent) power of J. In the low-J limit, this length is cut off by the film thickness, which only serves to increase the predicted resistivity over that for an infinite sample. We believe that these qualitative disagreements with our data arise because neither the flux creep nor the vortex ring model treats the dominant source of dissipation *in the regime of small measuring current.* One common characteristic of these models is a single energy barrier [27] for a given current density, leading to thermally activated resistance. By contrast, a broad distribution of activation energies leads naturally to a temperature-independent dissipation that depends on frequency.

Guided by these considerations, we have constructed a model in which each vortex in the flux transformer is confined to two symmetric potential wells and hops between them via thermal activation (inset to Fig. 3). Each such process changes the flux applied to the transformer by approximately $(\bar{l}/w)\Phi_0$, where $\bar{l}$ is the average vortex hopping distance and w is the linewidth of the coil. This change in flux induces a supercurrent in the transformer and hence a flux $\Delta\Phi = M_i \bar{l}\Phi_0/w(L_i+L_p)$ in the SQUID. Note that we model each vortex as a rigid rod, a justifiable approximation given that low-current-density measurements probe vortex correlation lengths of the same order as the sample thickness [18], and that our SQUID is sensitive only to the positions at which the ends of the vortex emerge from the sample. To make our calculation more tractable, we make the following simplifying assumptions: The activation barrier U_0 separating each pair of wells is temperature independent, the ensemble of such wells has a uniform number $D(U_0)$ of barriers per energy interval, $\bar{l}$ is independent of U_0, and the vortex displacements have no component parallel to the transport current. We emphasize that none of these approximations is essential to our argument, and, indeed, they were not made in our earlier work [20,24]. Under the influence of a static current I_0 and a current I_ω oscillating at frequency ω, generated, for example, by an applied magnetic field with static and oscillating components, we can describe the thermally activated motion of each vortex by a first-order differential equation for the probability $P(t)$ of finding the vortex in a given well at time t:

$$\frac{dP}{dt} = -\gamma_0 P(t)\exp\left(-\frac{U_0}{k_B T}\right)\exp[\alpha(I_0+I_\omega)] + [1-P(t)]\gamma_0\exp\left(-\frac{U_0}{k_B T}\right)\exp[-\alpha(I_0+I_\omega)]\,, \quad (3)$$

where $\alpha \equiv \Phi_0\bar{l}/2wk_BT$ and γ_0 is the attempt frequency [20] for the vortex in each well. The total voltage induced in the transformer by the motion of a single vortex is then

$$V = (L_i+L_p)dI/dt - (\Phi_0\bar{l}/w)dP/dt\,. \quad (4)$$

For arbitrary I_ω, Eqs. (3) and (4) can be solved numerically and yield a nonlinear frequency-dependent I-V characteristic. For small I_ω, however, these equations can be solved analytically, and we find the remarkable result that each vortex simply responds linearly at frequency ω as a resistance R_v^0 in series with an inductance L_v^0:

$$R_v^0 = \frac{1}{2k_BT\cosh\alpha I_0}\left(\frac{\Phi_0\bar{l}}{w}\right)^2 \frac{\gamma_0\exp(-U_0/k_BT)\omega^2}{\omega^2+[2\gamma_0\exp(-U_0/k_BT)\cosh\alpha I_0]^2}\,, \quad (5)$$

$$L_v^0 = \frac{1}{k_BT}\left(\frac{\Phi_0\bar{l}}{w}\right)^2 \frac{\gamma_0^2\exp(-2U_0/k_BT)}{\omega^2+[2\gamma_0\exp(-U_0/k_BT)\cosh\alpha I_0]^2}\,. \quad (6)$$

Integrating over an ensemble of uncorrelated (or randomly correlated) vortices, we find, for $\omega \ll \gamma_0$,

$$R_v = [\pi D(U_0)(\Phi_0\bar{l}/w)^2/8\cosh^2\alpha I_0]\omega\,, \quad (7)$$

$$L_v = \frac{D(U_0)}{8\cosh^2\alpha I_0}\left(\frac{\Phi_0\bar{l}}{w}\right)^2 \ln\{1+[2(\gamma_0/\omega)\cosh\alpha I_0]^2\}\,. \quad (8)$$

Equation (7) shows that R_v scales as ω, is independent of T for $I_0=0$, and decreases with increasing I_0 according to $R_v(I_0,f) = R_v(0,f)/\cosh^2\alpha I_0$, in good agreement with our experimental results. Thus, this model predicts all of the essential features of R_v measured experimentally. We note that Eqs. (7) and (8) would be difficult to test in a conventional transport measurement. Since $R_v \ll \omega(L_i$

$+L_p$) is itself proportional to ω, precise measurements of the phase of the complex impedance would be necessary to distinguish it from a small parasitic inductance. The value of L_v(1 Hz) predicted with R_v(1 Hz) = 30 pΩ is of the order of 0.1 nH, much smaller than the geometrical inductance of the flux transformer with which it is in series; it is virtually frequency independent at our experimental frequencies $\omega \ll \gamma_0$.

In conclusion, we have shown that the resistance of YBCO measured in the regime of low magnetic fields with very low or even zero current densities can be explained by a model in which each vortex is independently thermally activated between two pinning sites. The application of a static current does not induce a flow of these vortices, which would increase the dissipation, but rather causes each vortex to dwell for a longer time in one of the potential wells, thereby decreasing its hopping rate and its contribution to the dissipation. The introduction of a distribution in the barrier heights between the wells leads to a resistance that scales with frequency. The resistance is also approximately independent of temperature because, as the temperature is varied, a different set of barrier heights contributes to the active processes. This mechanism is in contrast to those processes used to account for dissipation in higher magnetic fields and current densities, and it would be of particular interest to investigate the behavior in the intermediate regime between the two limiting cases. Finally, we note that our measurements are made in a regime where the sample supports a persistent current and that this current can be as large as the critical current of the sample. This regime, and not one with a net flow of vortices, is of technological interest because most applications of superconductors require them to sustain a persistent current.

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