Collective Pinning of a Frozen Vortex Liquid in Ultrathin Superconducting $YBa_2Cu_3O_7$ Films

M. Calame,[1] S. E. Korshunov,[2] Ch. Leemann,[1] and P. Martinoli[1]
[1]*Institut de Physique, Université de Neuchâtel, CH-2000 Neuchâtel, Switzerland*
[2]*L. D. Landau Institute for Theoretical Physics, Kosygina 2, 117940 Moscow, Russia*
(Received 25 August 2000)

The linear dynamic response of the two-dimensional (2D) vortex medium in ultrathin $YBa_2Cu_3O_7$ films was studied by measuring their ac sheet impedance Z over a broad range of frequencies ω. With decreasing temperature the dissipative component of Z exhibits, at a temperature $T^*(\omega)$ well above the melting temperature of a 2D vortex crystal, a crossover from a thermally activated regime involving single vortices to a regime where the response has features consistent with a description in terms of a collectively pinned vortex manifold. This suggests the idea of a vortex liquid which, below $T^*(\omega)$, appears to be frozen at the time scales $1/\omega$ of the experiments.

DOI: 10.1103/PhysRevLett.86.3630 PACS numbers: 74.60.Ge, 74.25.Nf, 74.76.Bz

Strictly speaking no crystal is really a solid. At any nonzero temperature the equilibrium concentration of pointlike defects (vacancies and interstitials) as well as their mobility is finite, thereby providing a mechanism for mass flow unrelated to the motion of the crystal lattice. From a practical point of view, however, the overwhelming majority of crystals behaves as a solid at all relevant time scales, the (linear) plastic flow mediated by the motion of pointlike defects being too small to be of any relevance. In this Letter we show that the *two-dimensional* (2D) vortex medium in weakly disordered superconducting films exhibits a somewhat similar behavior.

It is generally accepted that the effect of weak disorder on flux lattices results in some sort of glassy state [1,2]. In a rigorous sense, however, a truly superconducting vortex-glass (VG) phase, such that the linear resistance vanishes, does not exist in two dimensions [3–5]. Although the random potential associated with disorder can quench the motion of the vortex medium as a whole, in two dimensions (in contrast to three) the flow of magnetic flux can be mediated by the motion of thermally created dislocation pairs, which are pointlike defects topologically equivalent to vacancies or interstitials.

Relying on sheet impedance studies, in this Letter we demonstrate that, although both theory [3–5] and experiment [6] point to the absence of a genuine vortex glass in two dimensions, the linear response of the 2D vortex medium in ultrathin $YBa_2Cu_3O_7$ (YBCO) films to a small ac field exhibits glasslike features at temperatures for which the motion of defects with respect to the driven vortex system appears to be frozen at the time scales set by the measurement. This "dynamic freezing" of the 2D vortex medium is reflected in the drastic change of the film response with decreasing temperature. At high temperatures the dissipative component of the impedance is frequency independent and exhibits a strong (exponential) temperature dependence, which can be attributed to the thermally activated motion of single vortices in a liquid subject to pinning. However, at lower temperatures, this contribution becomes too small to be observed. In this regime the dynamic response is governed by a different process, whose frequency and temperature dependences are found to be consistent with the idea of a collectively pinned vortex medium [7–10].

Quite remarkably, the crossover to the collective pinning regime is located at a temperature $T^*(\omega)$ which turns out to be much higher than the estimated melting temperature of a 2D vortex crystal in the absence of disorder. Thus, the response we observe well below $T^*(\omega)$ does not originate from a collectively pinned vortex crystal, but from a collectively pinned (dynamically) frozen vortex liquid (or, in other terms, a strongly disordered vortex solid). In this connection, it is worth mentioning that the classical theory of flux creep in the collective pinning regime [2–5] and its ac extensions [7–10] ignore the periodic nature of the vortex lattice and should therefore provide a reliable description of the dynamic response of a frozen vortex liquid.

The experiments were performed on two films, YBCO-2 and YBCO-4, respectively, two and four unit cells thick (i.e., with thicknesses $d = 2.4$ nm and $d = 4.8$ nm), grown epitaxially (c-axis oriented) by laser ablation onto (100) $SrTiO_3$ substrates. Both films were sandwiched between nonsuperconducting buffer and cover $PrBa_2Cu_3O_7$ layers. For comparison, a 110 nm-thick YBCO film was also studied. Their complex sheet impedance $Z \equiv R_z + i\omega L_z$ was extracted from the mutual-inductance change of a drive-receive coil system [11] operated with a conventional lock-in detector allowing inductances to be measured with a sensitivity of $\sim$10 pH between 30 Hz and 100 kHz. The bulk in-plane magnetic penetration depth $\lambda_{ab}(T)$ was inferred from measurements of the inverse sheet kinetic inductance $1/L_0 \equiv \omega\,\mathrm{Im}[1/Z(B = 0)] = d/\mu_0\lambda_{ab}^2$ in zero magnetic field and found to fit well, over a wide temperature range (excluding the critical region), a parabolic dependence [12] $\lambda_{ab}^{-2}(T) = \lambda_{ab}^{-2}(0)[1 - (T/T_c)^2]$ with $\lambda_{ab}(0) = 550$ nm and $T_c = 48.3$ K for YBCO-2 and $\lambda_{ab}(0) = 340$ nm and $T_c = 73.7$ K for YBCO-4.

The inverse sheet inductance $1/L_g = \omega\,\mathrm{Im}(1/Z)$ (which measures the degree of superconducting phase coherence

 0031-9007/01/86(16)/3630(4)$15.00

in the system) and R_z of YBCO-4 in a perpendicular field of 0.01 T are plotted logarithmically in Fig. 1 as a function of $1/T$ for a set of four representative frequencies. As highlighted by the straight dotted lines, at high temperatures both $1/L_g(T)$ and $R_z(T)$ exhibit an Arrhenius-like behavior, thereby pointing to a thermally activated process involving single vortices. In this regime (where $R_z \gg \omega L_z$) $R_z(T)$ is almost frequency independent, whereas $1/L_g(T)$, as shown in the inset of Fig. 1, is proportional to ω^2. These features are in excellent agreement with what one expects from barrier limited diffusion of noninteracting particles [13].

The linearity of the Arrhenius plots in the thermally activated regime implies that the activation energy $U(T,B)$ is either constant or linearly dependent on temperature. We expect the latter to be the case, since $U(T,B)$ should be proportional to the basic energy scale of vortex matter, $\epsilon(T) = \phi_0^2/4\pi\mu_0\lambda_{ab}^2(T)$ (ϕ_0 is the superconducting flux quantum and μ_0 the induction constant), which varies as $(T_c - T)$ in the temperature range of interest just below T_c. Then, noticing that the temperature dependence of λ_{ab}^{-2} discussed above leads to $U(T,B) \approx 2U(0,B) \times [1 - (T/T_c)]$ near T_c, our data allow one to explore the dimensionality of the vortex medium by studying the field dependence of the zero-temperature activation energy $U(0,B)$.

The results for the three YBCO samples studied in this work are shown in Fig. 2. While the energy barrier for the thickest (reference) film [Fig. 2(a)] obeys a power law $U(0,B) \propto B^{-\alpha}$ with an exponent $\alpha \approx 0.40 \pm 0.05$ fairly close to the prediction $\alpha = 1/2$ for plastic vortex motion in a 3D vortex liquid [5], for the thinnest YBCO-2 layer [Fig. 2(b)] $U(0,B)$ exhibits, over the entire field range covered by our experiments, a logarithmic field dependence signaling 2D behavior. This interpretation is corroborated by a quantitative comparison with the theoretical predictions for flux flow in two dimensions controlled by surface barriers [14]. In this regime the prelogarithmic factor of $U(0,B)$ is of the form $C\epsilon(0)d$, with $C = 1/2$. Inspection of the slope of $U(0,B)$ for YBCO-2 in Fig. 2(b) gives $\partial U(0,B)/\partial \ln(1/B) \approx 71$ K, corresponding to $C \approx 0.45$, a value in excellent agreement with that calculated for the surface barrier mechanism. Note that the response of samples of macroscopic size as those studied in this work would be dominated by edge barriers only at extremely low frequencies, far below those accessible to our experiment. However, as evidenced by atomic force microscope images of the film microstructure, a consequence of the "unit cell by unit cell" growth of the YBCO layers is the formation of linear defects (steps), related to thickness variations in multiples of a unit cell, separating flat islands ~0.2 μm in size.

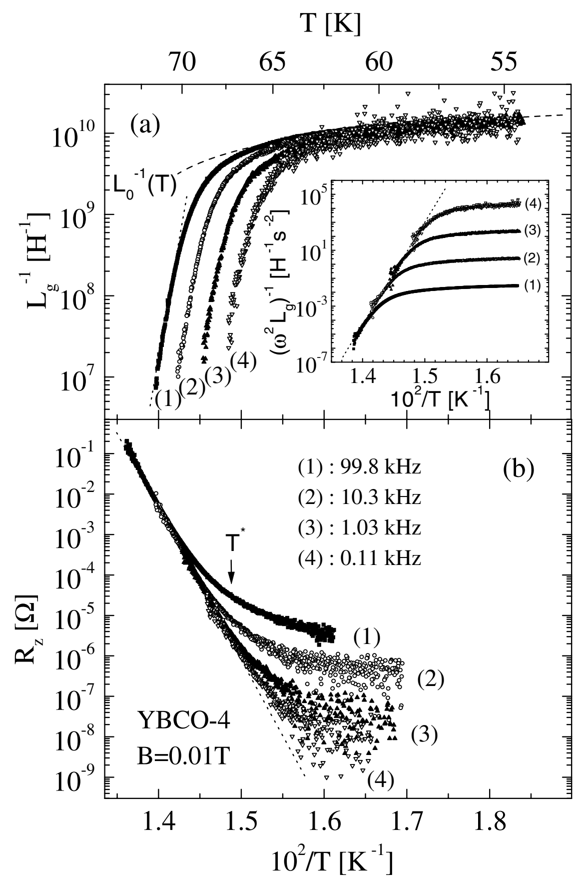

FIG. 1. Arrhenius plots of (a) $1/L_g$ and (b) R_z of YBCO-4 at $B = 0.01$ T for four representative frequencies. The dashed curve in (a) is the inverse kinetic inductance in zero field. T^* indicates the crossover temperature of R_z at 99.8 kHz.

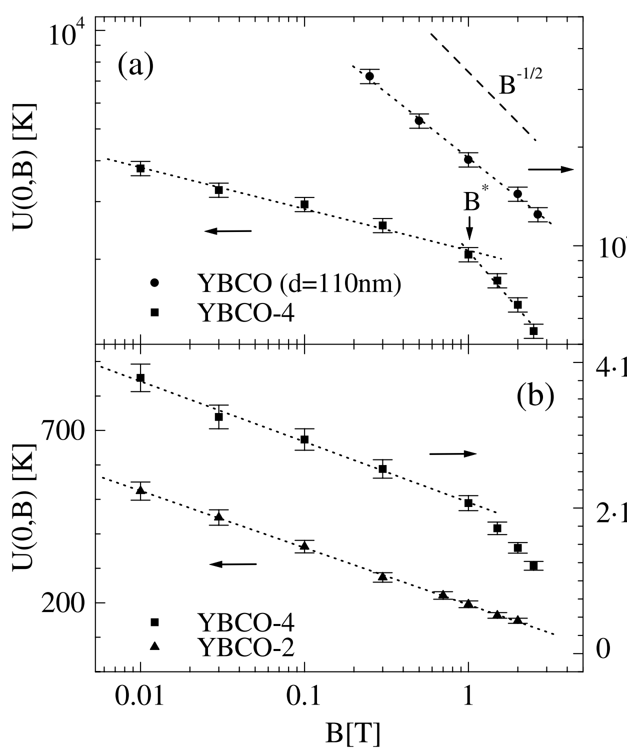

FIG. 2. (a) Log-log and (b) lin-log plots of the zero-temperature activation energy as a function of the magnetic field. The dotted lines are fits through the data. The dashed line in (a) is a $B^{-1/2}$ power law. B^* indicates the 2D-3D crossover of the vortex medium in YBCO-4.

It can then be expected that such steps play a role similar to that of sample edges, as they would also provide barriers against vortex motion exhibiting, in two dimensions, a logarithmic field dependence.

Additional evidence for the 2D-3D dimensional crossover of the vortex medium is provided by a study of the activation energy for YBCO-4. As shown in Fig. 2(a), for this film $U(0,B)$ exhibits algebraic behavior with $\alpha \approx 0.58 \pm 0.06$ at high fields, but crosses over, at $B^* \approx 1$ T, to a low-field regime characterized by a much smaller exponent ($\alpha \approx 0.1$) pointing to a logarithmic field dependence, which is indeed demonstrated by the lin-log plot of Fig. 2(b). By comparing the activation energies for 3D [5,14] and 2D [14] vortex liquids, it is possible to estimate the crossover field as $B^* \approx k\phi_0(\gamma/d)^2$, where γ is the anisotropy ratio and k a numerical constant of order unity. Using $\gamma \approx 1/7$ for YBCO, one obtains $B^* \approx 1$ T for YBCO-4 by choosing $k \approx 1/2$.

As shown by Fig. 1, with decreasing temperature both $1/L_g(T)$ and $R_z(T)$ cross over to a regime where their temperature dependence becomes much weaker than in the activated regime. While the change in behavior in $1/L_g(T)$ would be present even if $R_z(T)$ would continue to decrease exponentially with $1/T$, the crossover in $R_z(T)$ points to the onset of a different regime. In order to elucidate its physical nature, it is convenient to define a crossover temperature $T^*(\omega)$ and to compare it with the melting temperature of the vortex medium. We *ad hoc* identify $T^*(\omega)$ as the temperature corresponding to the maximum curvature of the $R_z(T)$ curves in Fig. 1, the particular choice of the criterion being irrelevant for our conclusions. At $B = 0.01$ T, $T^*(\omega)$ varies approximately from 67 K at the upper limit to 64 K at the lower end of the frequency spectrum explored in our measurements [the maximum curvature in the $1/L_g(T)$ curves occurs at about the same temperatures]. Quite remarkably, these crossover temperatures are much higher than the temperature, $T_M = \phi_0^2 d/32\pi^2\sqrt{3}\,k_B\mu_0\lambda_{ab}^2(T_M) \approx 18$ K (k_B is the Boltzmann constant), at which the 2D vortex crystal in YBCO-4 would melt due to dislocation unbinding [15] in the absence of disorder (notice that the data of Fig. 1 were taken in the 2D regime well below B^*). Thus, the picture emerging from the impedance measurements is that of a vortex liquid which, well below $T^*(\omega)$, appears to be frozen at the time scales, $1/\omega$, of our measurements and whose response, as is shown below, can be described by ac extensions of the theory of collective pinning.

The low-frequency linear dynamic response of a collectively pinned elastic vortex manifold has been investigated in Refs. [7–9] and recently, in a more systematic way, in Ref. [10]. In this approach large (in general anisotropic) portions of the vortex medium (vortex bundles) are assumed to be fluctuating between pairs of metastable states in the uncorrelated random potential. Treating these pairs of states as current-driven two-level systems with a size distribution $\nu(r)$, the vortex contribution $l_{zv}(T,\omega)$ to the specific inductance of the system can be shown [7,10] to be, quite generally, of the form:

$$l_{zv}(T,\omega) \sim B^2 \int_{r_c}^{r_\omega} dr\, \nu(r)\, \frac{V^2u^2}{E}\,, \tag{1}$$

where u is the average vortex displacement inside a bundle in the direction of the current-induced force, r the size of a bundle in the same direction, V its volume, E the energy scale which can be associated with it, r_c the collective pinning length [5], and $r_\omega \sim r_c[(k_BT/U_0)\ln(1/\omega\tau)]^{1/\chi}$ a length scale setting the maximum size of the vortex bundles which appreciably contribute to the response at times of the order of $\sim 1/\omega$ (U_0 and τ are, respectively, a characteristic energy and a relaxation time related to disorder, χ the energy barrier exponent). It turns out to be convenient to estimate E by considering the compressive contribution E_c to the energy. In doing this, one has to take into account the dispersive nature of the compression modulus $c_{11}(q)$, which in thin films ($d \ll \lambda_{ab}$) is always nonlocal. Using $c_{11}(q) \approx (B^2d/\mu_0\lambda_{ab}^2)q^{-2}$ in the regime of interest, one obtains [5] $E_c \sim (B^2d/\mu_0\lambda_{ab}^2)S^2(u/r)^2$, where $S = V/d \propto rr_\perp$. Then, setting $\nu(r) \propto 1/r^3$ as imposed by the presence of a hierarchical distribution of quasi-isotropic ($r_\perp \propto r$) two-level systems [10], from Eq. (1) we find

$$L_z(T,\omega) \approx L_0(T)\{1 + C'\ln[(k_BT/U_0)\ln(1/\omega\tau)]\}, \tag{2}$$

where C' is a numerical constant. Since $L_z(\omega)$ depends only logarithmically on ω, a simplified Kramers-Kronig relation [16] can be used, which leads to

$$R_z(T,\omega) \approx C'(\pi/2)\omega L_0(T)/\ln(1/\omega\tau)\,. \tag{3}$$

Notice that in a film with a high density of linear defects (steps) acting as a network of strong pinning lines the regime of collective pinning can be realized by vortices moving along these linear defects.

The main features of the response of our very thin YBCO layers in the frozen vortex-liquid regime below $T^*(\omega)$ are well described by these expressions [note that, below $T^*(\omega)$, $R_z \ll \omega L_z$ and, therefore, $L_g \approx L_z$]. Focusing first on the temperature dependence, one sees that it should be dominated by $L_0(T)$, the logarithmic corrections entering Eqs. (2) and (3) varying too slowly to be of any relevance in the limited temperature interval of our experiments. The dashed curve in Fig. 1(a), which is simply $1/L_0(T)$ as inferred from zero-field measurements, fits nicely the low-temperature (almost frequency-independent) $1/L_g(T)$ curves. On the other hand, the temperature dependence and the order of magnitude of R_z below $T^*(\omega)$ are compatible with Eq. (3) for any reasonable estimate of τ (10^{-12}–10^{-7} s). We have also found that the response in the frozen vortex-liquid regime is only weakly dependent on B in the field range covered by our experiments (up to 3 T), in agreement with the theoretical prediction (derived for $B \ll B_{c2}$).

Further evidence for the interpretation of the response in terms of a collectively pinned vortex manifold emerges

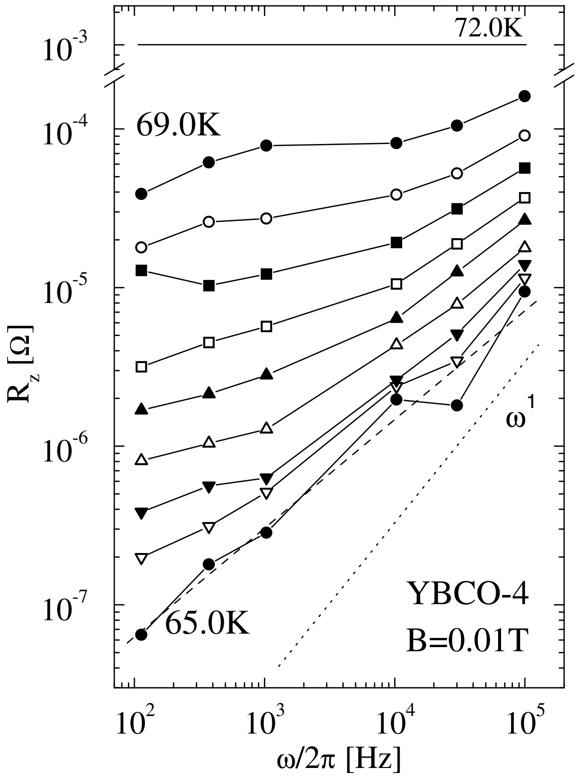

FIG. 3. Sheet resistance vs frequency isotherms taken at 0.5 K intervals between 65 and 69 K for YBCO-4 at $B = 0.01$ T on a log-log plot. The dashed line is a power-law fit of the 65 K data with an exponent ≈0.7. The dotted line is the (almost) linear ω dependence predicted by Eq. (3).

from the analysis of the frequency dependence of our data. As shown in Fig. 1(a), at low temperatures, well below $T^*(\omega)$, $1/L_g$ is almost frequency independent, a behavior consistent with Eq. (2), where the extremely slow-varying (double-logarithmic) vortex contribution can be hardly expected to be noticeable against the superfluid background $1/L_0$. To discuss the dissipative component, in Fig. 3 we show a family of R_z vs ω isotherms in a log-log plot. With decreasing temperature the isotherms progressively evolve from the frequency-independent regime characteristic of the vortex liquid at high temperatures to an almost algebraic behavior with an exponent ~0.7 at the lowest temperature (65 K) at which dissipation could be studied with sufficient accuracy. Considering the fact that this isotherm reflects, at a time scale $\sim 1/\omega$, the response of a 2D vortex medium "on the verge of freezing" rather than that of a "deep-frozen" liquid, we interpret the general behavior emerging from Fig. 3 and, in particular, the value of the exponent extracted from the "coldest" isotherm, as an indication that $R_z(\omega)$ will likely tend to the almost linear frequency dependence predicted by Eq. (3) well below $T^*(\omega)$. This should be compared with the response of a nonfrozen pinned vortex liquid for which one would expect a crossover to an ω^2 dependence at sufficiently low temperatures [13].

To conclude, for the first time we have shown that, although a vortex glass, strictly speaking, cannot exist in two dimensions, the associated collective pinning behavior is accessible to experimental observation. Earlier investigations of ultrathin YBCO films [6] were devoted to the formal proof of the absence of a VG phase transition in two dimensions inferred from a critical scaling analysis of their current-voltage curves. Unlike these studies, our work focused on a detailed comparison of different dynamic regimes of 2D vortex matter in weakly disordered superconductors. In our ac measurements the difference between a genuine VG phase transition and a simple dynamic freezing could become manifest in the behavior of $T^*(\omega)$ at ultralow frequencies. If a true phase transition occurs, $T^*(\omega)$ should saturate, with decreasing ω, at the (nonzero) VG transition temperature, whereas in the opposite case the decrease of $T^*(\omega)$ should be unlimited. Within the spectral range (covering three decades in frequency) accessible to our method, $T^*(\omega)$ was observed to decrease by only 5%, thereby preventing us from drawing any conclusion with regard to the existence of a genuine phase transition. Qualitatively, the ac response of ultrathin YBCO layers was found to be identical to that of thick films, for which the existence of a true 3D VG phase relies on solid experimental evidence [17]. Thus, from a practical point of view the glasslike features we observe below $T^*(\omega)$, although arising from a liquid, turn out to be undistinguishable from those of a genuine vortex glass.

We thank L. Baselgia Stahel, S. Blaser, A. Daridon, A. Rüfenacht, and B. Schmied for their valuable contributions to the experimental aspects of this work. We are also grateful to G. Blatter and V. B. Geshkenbein for interesting discussions and useful comments. This work was supported by the Swiss National Science Foundation.

[1] M. P. A. Fisher, Phys. Rev. Lett. **62**, 1415 (1989).
[2] M. V. Feigel'man *et al.,* Phys. Rev. Lett. **63**, 2303 (1989).
[3] M. V. Feigel'man *et al.,* Physica (Amsterdam) **167C**, 177 (1990).
[4] V. M. Vinokur *et al.,* Physica (Amsterdam) **168C**, 29 (1990).
[5] G. Blatter *et al.,* Rev. Mod. Phys. **66**, 1125 (1994).
[6] C. Dekker *et al.,* Phys. Rev. Lett. **69**, 2717 (1992).
[7] A. E. Koshelev and V. M. Vinokur, Physica (Amsterdam) **173C**, 465 (1991).
[8] D. S. Fisher *et al.,* Phys. Rev. B **43**, 130 (1991).
[9] K. H. Fischer and T. Nattermann, Phys. Rev. B **43**, 10 372 (1991).
[10] S. E. Korshunov, cond-mat/0007387.
[11] B. Jeanneret *et al.,* Appl. Phys. Lett. **55**, 2336 (1989).
[12] M. Tinkham, *Introduction to Superconductivity* (McGraw Hill, New York, 1996), 2nd ed., p. 381.
[13] C. J. van der Beek *et al.,* Phys. Rev. B **48**, 3393 (1993).
[14] L. Burlachkov *et al.,* Phys. Rev. B **50**, 16 770 (1994).
[15] B. A. Hubermann and S. Doniach, Phys. Rev. Lett. **43**, 950 (1979); D. S. Fisher, Phys. Rev. B **22**, 1190 (1980).
[16] L. Lundgren *et al.,* J. Magn. Magn. Mater. **25**, 33 (1981).
[17] P. J. M. Wöltgens *et al.,* Phys. Rev. B **52**, 4536 (1995), and references therein.