Collective Pinning of a Frozen Vortex Liquid in Ultrathin Superconducting YBa$_2$Cu$_3$O$_7$ 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

The linear dynamic response of the two-dimensional (2D) vortex medium in ultrathin YBa$_2$Cu$_3$O$_7$ films was studied by measuring their ac sheet impedance $Z$ over a broad range of frequencies $\omega$. 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.

PACS numbers: 74.60.Ge, 74.76.Bz, 74.25.Nf

It is generally accepted that the effect of weak disorder on flux lattices in superconductors results in some sort of glassy state \cite{14}. In a rigorous sense, however, the existence of a truly superconducting vortex glass phase, such that the linear resistance vanishes, is impossible in two dimensions \cite{11}. 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 point-like defects. To the best of our knowledge, only limited experimental evidence for the absence of the vortex-glass transition in two dimensions has been found so far \cite{12,13,16}.

Relying on sheet impedance measurements, in this Letter we demonstrate that, at typical laboratory time scales, the linear response of the two-dimensional (2D) vortex medium in ultrathin YBa$_2$Cu$_3$O$_7$ (YBCO) films to a small driving ac field drastically changes 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. 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 ($\omega$) and temperature ($T$) dependences are found to be consistent with the idea of a collectively pinned vortex medium, in which large portions of the vortex manifold are fluctuating between pairs of metastable states in the random potential landscape provided by the film microstructure \cite{12,16}.

The crossover to the collective pinning regime has been observed over the whole frequency range explored in our experiments. Quite remarkably, it 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, from a strongly disordered vortex solid). On the other hand, a study of the liquid-like response well above $T^*(\omega)$ as a function of the applied magnetic field $B$ reveals that in this regime the vortex activation process is controlled by a surface barrier mechanism \cite{11}.

It is worth mentioning that the classical theory of flux creep in the collective pinning regime \cite{15,16} and its ac extensions \cite{12,17} ignore the periodic nature of the vortex lattice, the vortex medium being treated as an elastic manifold interacting with a random potential. Therefore, this approach should provide a reliable description of a disordered vortex solid (or frozen vortex liquid) as long as the frequency of observation is high enough for the motion of defects with respect to the driven vortex medium to appear to be frozen.

A possible criterion to distinguish between a genuine continuous phase transition to a vortex glass and a simple dynamic crossover to a collective pinning regime could rely on studies of $T^*(\omega)$ extending at ultralow frequencies. If a true phase transition occurs, $T^*(\omega)$ should saturate with decreasing $\omega$ at the (nonvanishing) vortex-glass transition temperature, whereas in the opposite case the decrease of $T^*(\omega)$ should be unlimited. In our experiments $T^*(\omega)$ decreases by only 5% over three decades in frequency, thereby preventing us from drawing any conclusion with regard to the existence of a genuine phase transition. Qualitatively, in the frequency range explored in our experiments the ac response of ultrathin YBCO layers is identical to that of thick films, for which signatures of a true 3D vortex-glass phase were reported by several authors. Thus, from a practical point of view the glass-like features we observe below $T^*(\omega)$ turn out to be indistinguishable from those of a genuine vortex glass.

The experiments were performed on two films YBCO-2 and YBCO-4, respectively 2 and 4 unit cells thick (i.e., with thicknesses $d = 2.4nm$ and $d = 4.8nm$), grown epitaxially (c-axis oriented) by laser ablation onto (100) SrTiO$_3$ substrates. Both films were sandwiched between

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

nonsuperconducting buffer and cover PrBa$_2$Cu$_3$O$_7$ layers. For comparison, a thick ($d = 110nm$) YBCO film was also studied. Their complex sheet impedance $Z$ was extracted from the mutual-inductance change of a drive-receive coil system [12] operated with a conventional lock-in detector allowing inductances to be measured with a sensitivity of $\sim 10pH$ between 30Hz and 100kHz. The bulk in-plane magnetic penetration depth $\lambda_{ab}(T)$ was inferred from measurements of the inverse sheet kinetic inductance $1/L_0 \equiv \omega \text{Im}[1/Z(B = 0)] = d/\mu_0 \lambda_{ab}^2$ in zero magnetic field and found to fit well, over a wide temperature range, a parabolic dependence [13] $\lambda_{ab}^2(T) = \lambda_{ab}^2(0)[1 - (T/T_c)^2]$ with $\lambda_{ab}(0) = 550nm$ and $T_c = 48.3K$ for YBCO-2 and $\lambda_{ab}(0) = 340nm$ and $T_c = 73.7K$ for YBCO-4.

The inverse sheet inductance $1/L_g = \omega \text{Im}(1/Z)$ (which measures the degree of superconducting phase coherence in the system) and the resistive component of the sheet impedance $R_z = \text{Re}(Z)$ (which measures dissipation) of YBCO-4 in a perpendicular field of 0.01T are plotted logarithmically in Fig. 1 as a function of $1/T$ for a set of 4 representative frequencies. As highlighted by the dotted straight 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 $R_z(T)$ is almost frequency independent, whereas $1/L_g(T)$, as shown in the insert of Fig. 1, is a $B^{-1/2}$ power law. These features are in excellent agreement with what one expects from barrier limited diffusion of non interacting particles.

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, $\epsilon(T) = \phi_0^2/4\pi\mu_0 \lambda_{ab}^2(T) (\phi_0$ is the superconducting flux quantum and $\mu_0$ 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 $\lambda_{ab}^2$, discussed above leads to $U(T,B) \approx 2U(0,B)[1 - (T/T_c)]$ near $T_c$, our data allow 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 [14], 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 $C = 1/2$. Figure 2. (a) Log-log and (b) lin-log plots of the zero-temperature activation energy as a function of the magnetic field. The dashed lines are fits through the data. The dotted line in (a) is a $B^{-1/2}$ power law. $B^*$ indicates the 2D-3D crossover of the vortex medium in YBCO-4.
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 AFM 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 $\sim 0.2 \mu m$ in size. It can be then 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 $B^* \approx 1T$, 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 linear-log plot of Fig. 2(b). By comparing the activation energies for 3D and 2D vortex liquids, it is possible to estimate the crossover field as $B^* \approx k\phi_0 (\gamma/d)^2$, where $\gamma$ 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 $R_c(T)$ 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 $R_c(T)$ would continue to decrease exponentially with $1/T$, the crossover in $R_c(T)$ 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 $R_c(T)$-curves in Fig. 1, the particular choice of the criterion being irrelevant for our conclusions. At $B = 0.01T$, $T^*(\omega)$ 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, $T_m = \phi_0 d/32\pi^2 \sqrt{3k_B \mu_0} \lambda_2^2(T_m) \approx 15 K$ ($k_B$ is the Boltzmann constant), at which the 2D vortex crystal in YBCO-4 would melt due to dislocation unbinding 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. [8,9] 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 $l_{z\nu}(T,\omega)$ to the specific inductance of the system can be shown to be, quite generally, of the form:

$$l_{z\nu}(T,\omega) \sim B^2 \int_{r_c}^{r_\nu} d\nu(r) \frac{V^2 u^2}{E}$$  \hspace{1cm} (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 and $r_\nu \sim r_c[(k_B T/\mu_0)\ln(1/\omega \tau)]^{1/\chi}$ a frequency-dependent length scale setting the maximum size of the vortex bundles which appreciably contribute to the response ($U_0$ and $\tau$ are, respectively, a characteristic energy and a relaxation time related to disorder, $\chi$ 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 $c_{11}(q) \approx (B^2 d/\mu_0 \lambda_{ab}^2) q^{-2}$ in the regime of interest, one obtains $E_c \sim (B^2 d/\mu_0 \lambda_{ab}^2) S^2(u/r)^2$, where $S = V/d \propto L_{r\perp}$. 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 from Eq. (1) we find:

$$L_z(T,\omega) \approx L_0(T) \{1 + C' \ln[(k_B T/\mu_0)\ln(1/\omega \tau)]\},$$  \hspace{1cm} (2)$$

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

$$R_z(T,\omega) \approx C' (\pi/2) \omega L_0(T)/\ln(1/\omega \tau).$$  \hspace{1cm} (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 line 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) $1/L_0(T)$-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 $\tau$ ($10^{-12} - 10^{-7}s$). 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 $3T$), 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 from the analysis of the frequency dependence of our data. As shown in Fig. 1(a), at low temperatures, well below $T^*(\omega)$, $1/L_0$ 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 $\omega$ 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 $\sim 0.7$ at the lowest temperature ($65K$) 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 non-frozen pinned vortex liquid for which one would expect a crossover to an $\omega^2$-dependence at sufficiently low temperatures.

We would like to thank L. Baselgia Stahel, S. Blaser and B. Schnied for their contribution during the early stages of this work, A. Daridon for the AFM images and A. Riufenacht for assistance in the analysis of the data. 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, V.B. Geshkenbein, A.I. Larkin and V.M. Vinokur, Phys. Rev. Lett. 63, 2303 (1989).
[3] M.V. Feigel’man, V.B. Geshkenbein and A.I. Larkin, Physica C167, 177 (1990).
[4] V.M. Vinokur, P.H. Kes and A.E. Koshelev, Physica C168, 29 (1990).
[5] G. Blatter, M.V. Feigel’man, V.B. Geshkenbein, A.I. Larkin and V.M. Vinokur, Rev. Mod. Phys. 66, 1125 (1994).
[6] C. Dekker, P.J.M. Wöltgens, R.H. Koch, B.W. Hussey, and A. Gupta, Phys. Rev. Lett. 69, 2717 (1992).
[7] Z. Sefrioui, D. Arias, M. Varela, J.E. Villegas, M.A. López de la Torre, C. León, G.D. Loos, and J. Santamaria, Phys. Rev. B 60, 15423 (1999).
[8] A.E. Koshelev and V.M. Vinokur, Physica C 173, 465 (1991).
[9] D.S. Fisher, M.P.A. Fisher and D.A. Huse, Phys. Rev. B 43, 130 (1991).
[10] S.E. Korshunov, cond-mat/0007385 (2000).
[11] I. Burlakov, V.B. Geshkenbein, A.E. Koshelev, A.I. Larkin, and V.M. Vinokur, Phys. Rev. B 50, 16770 (1994).
[12] B. Jeanneret, J.L. Gavilano, G.-A. Racine, Ch. Leemann, and P. Martinoli, Appl. Phys. Lett. 55, 2336 (1989).
[13] M. Tinkham, Introduction to Superconductivity, McGraw Hill, New York, 1996, 2nd edition, p. 381.
[14] A.I. Larkin and Yu.N. Ovchinnikov, J. Low Temp. Phys. 34, 409 (1979).
[15] L. Lundgren, P. Svedlindh and O. Beckman, J. Magn. Magn. Mat. 25, 33 (1981).