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Plastic vortex creep above the second magnetization peak in $Bi_2Sr_2CaCu_2O_{8+\delta}$ single crystals

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Abstract

The analysis of the activation energy in the magnetization relaxation of $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) single crystals reveals plastic vortex creep above the second magnetization peak (SMP). Above the peak field, the magnetic field dependence of the activation energy $U(B)$ is very close to the form $U(B) \propto B^{-1/2}$ and the intrinsic variation of U with the current density is weak. The existence of a crossover from elastic to plastic vortex creep across the SMP of high-temperature superconductors (HTSC) seems to be a general behavior, and may have important consequences on the nature of the thermally induced vortex solid–vortex fluid transition at high magnetic fields.

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Vortex matter in high-temperature superconductors (HTSC) is a field of intensive study [1]. Despite remarkable experimental and theoretical efforts, the vortex phase diagram is still far from being completely elucidated. Over the last few years, it has become apparent that at least two distinct vortex phases can exist: the vortex solid and the vortex fluid. In clean HTSC, the vortex lattice at low magnetic fields undergoes a thermally induced first-order melting [2–4] or sublimation [5,6] transition, as evidenced by a discontinuity in the equilibrium magnetization [7,8], a simultaneous sharp hysteretic drop in the electrical resistance [9], and by specific heat measurements [10]. When the quenched disorder is significant, this transition is expected to be replaced by a second-order vortex-glass transition [11]. The nature of the thermally induced vortex solid–vortex fluid transition at high magnetic fields is not yet fully understood.

An additional phase boundary was discovered within the region that is assumed to correspond to the vortex solid [12,13], and manifests itself by a peculiar increase of the absolute value of the irreversible magnetization with increasing magnetic field, known as the second magnetization peak (SMP). The SMP was attributed to surface barriers [12], sample inhomogeneities [13], a crossover from surface barriers to bulk pinning [14], dynamic effects

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[15], a dimensional transition [16], a weak first-order vortex-lattice melting [17], layer decoupling [18], or vortex stacking [19].

By considering the competition between the elastic energy of the vortex system and the pinning energy generated by the quenched disorder, it has been suggested [20–24] that the SMP can result from a disorder-induced transition of the low-field quasi-ordered vortex phase (or Bragg glass [20–23]) to a disordered vortex solid at higher fields [25,26]. This transition is located at the onset of the SMP [27]. Evidence for the existence of two distinct vortex-solid phases was previously obtained in neutron diffraction experiments [28], for example. However, the nature of the disordered vortex phase above the SMP is still unclear. The increase of the effective pinning suggests that this should be a vortex glass [26,29], with the pinning barrier exhibiting the specific increase at low current densities [1]. Alternatively, the proliferation of defects in the vortex system across the SMP can lead to a dissipation process dominated by plastic vortex deformations [30,31].

In this work, we address the above issue through DC magnetization measurements on $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) single crystals. We observed a change in the current density and magnetic field dependence of the activation energy in the magnetic relaxation process, revealing a crossover from elastic to plastic vortex creep at the peak field. The vortex phase above the SMP behaves like a plastic vortex solid, rather than a vortex glass.

The DC magnetization M of several BSCCO single crystals grown by a self-flux technique was measured as a function of the applied magnetic field H and time t. The crystals obtained by this method contain growth defects, which inhibit the fluctuations of vortices, up to a certain extent, and diminish the influence of geometrical and surface barriers. The results presented here are for a $0.56 \times 0.55 \times 0.025$ mm^3 sample. The measurements were performed in a Quantum Design SQUID magnetometer, with the external magnetic field applied along the c-axis, using the persistent mode of the superconducting magnet and the 3 cm scanning length. The onset of diamagnetism is detected below a temperature $T_c \approx 87$ K. The T_c and the peak-field values (see below) show that the measured crystal is slightly overdoped [32].

Fig. 1 shows the $M(H)$ variation in increasing H at several temperatures ($T = 20$, 25, and 30 K) and the location of the onset field H_{on} and the peak field H_p. The SMP is broader in standard global magnetization measurements than in local magnetic field measurements [25]. This is a consequence of sample inhomogeneity and a non-uniform magnetic induction inside the crystal. As can be seen, the amplitude of the SMP decreases with increasing T, and the peak is within our experimental noise for $T \geq 35$ K. It was argued that, at high temperatures, thermal fluctuations become sufficiently strong to smear out the effect of a weak quenched disorder, and the SMP phase boundary in the vortex solid merges with the thermally induced first-order melting of the low-field Bragg glass [25].

Fig. 2 illustrates the relaxation of the irreversible magnetization $M_{irr}(t)$, in a double logarithmic plot, for $T = 25$ K and several H values. The irreversible magnetization was determined from the magnetic hysteresis curves $M(H)$, by the standard procedure, $M_{irr}(H) = [M_+(H) - M_-(H)]/2$, where $M_+(H)$ and $M_-(H)$ represent the magnetization measured in increasing and decreasing H, respectively. As a

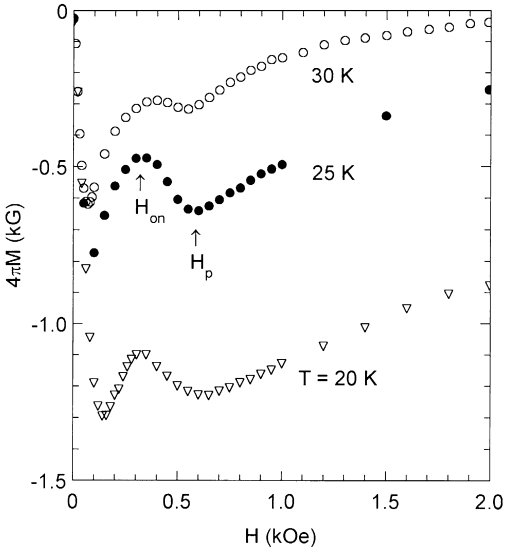

Fig. 1. Magnetic field H dependence of the magnetization M of BSCCO single crystals measured in zero-field-cooling conditions and increasing H at several temperatures T, revealing the occurrence of the second magnetization peak. The onset field H_{on} and the peak field H_p are indicated by arrows.

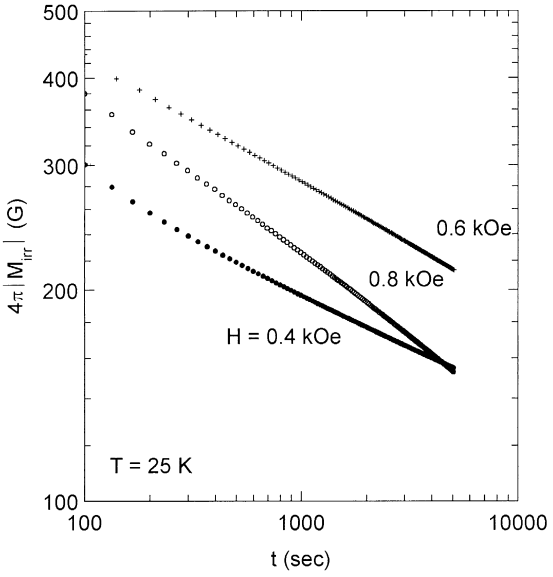

Fig. 2. The absolute value of the irreversible magnetization M_{irr} vs. time t in a double logarithmic plot, for $T = 25$ K and several H values in the second magnetization peak region.

rule, the first data point on the $M(t)$ curve [as well as $M_+(H)$ and $M_-(H)$] was taken $t_1 \approx 100$ s after the field was applied. To obtain $M_{irr}(t)$, the measured $M(t)$ curve was shifted on the y-axis by $M(t_1) - M_{irr}(t_1)$, which contains the (non-relaxing) reversible magnetization of the sample and the magnetization of the sample holder (also reversible). The activation energy U relevant for our study can be determined as [33,34]

$$U = -T\left[\mathrm{d}\ln|M_{irr}|\mathrm{d}\ln(t)\right]^{-1}. \quad (1)$$

In a first approximation, one can assume that $M_{irr}(t)$ is linear in $\ln(t)$. In this case, the variation of the activation energy with H can be rapidly obtained from the $M(H)$ curves, as those illustrated in Fig. 1, measured at $t_1 = 100$ s and $t_2 = 1000$ s (as considered here) after H was applied. These are shown in Fig. 3a, for H in the region of the SMP at $T = 25$ K. Thus, $U(H) = T\ln(t_2/t_1)M_{irr}(H)/\Delta M(H)$, where $\Delta M(H) = M(H,t_1) - M(H,t_2)$ and the irreversible magnetization $M_{irr}(H) = [M_+(H,t_1) - M_-(H,t_1)]/2$. The variation of U with the magnetic induction B is illustrated in Fig. 3b, where $B(H) = H + 4\pi M(H)(1 - D)$. The demagnetization factor $D \approx 0.93$ was extracted from the initial slope of the magnetization curves (Fig. 1). As can be seen, above the peak induction $B_p(25\ \mathrm{K}) = B(H_p) \approx 0.52$ kG,

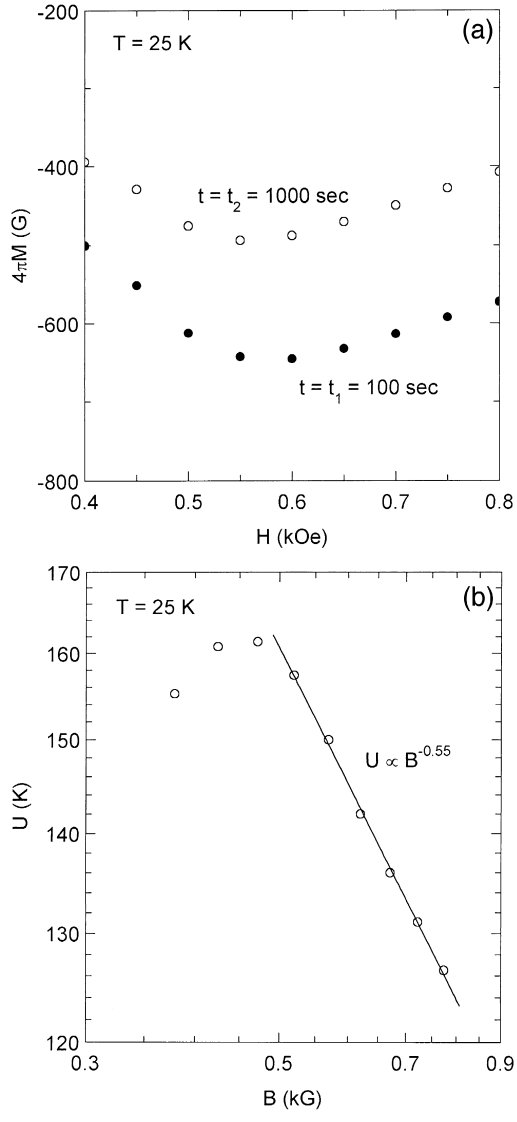

Fig. 3. (a) Magnetic field dependence of the magnetization M across the second magnetization peak at $T = 25$ K measured 100 s (●) and 1000 s (○) after the field was applied. (b) The mean activation energy U vs. the magnetic induction B in a double logarithmic plot at $T = 25$ K. Above the peak induction $B(H_p) \approx 0.52$ kG, $U(B) \propto B^{-0.55}$, as indicated by the linear fit represented by the continuous line.

the $U(B)$ dependence is very close to $U(B) \propto B^{-1/2}$, suggesting the plastic deformation of the vortex system [30,31]. Plastic deformations, including vortex cutting and reconnection, are on the scale of the mean intervortex spacing $a_0 \approx (\Phi_0/B)^{1/2}$, leading to a plastic barrier $U_{pl} \propto B^{-1/2}$ [1].

However, the slope of the relaxation curves from Fig. 2 exhibits some variations, especially for H below H_p. A new aspect is revealed when U determined with Eq. (1) is plotted as a function of $|M_{irr}|$, and the latter is substituted by the current density J, using the Bean model. Fig. 4a shows the resulting $U(J)$ dependence for several H values between H_{on} and H_p, at $T = 25$ K. For this field range, there is a first rapid increase of U with decreasing J, resembling elastic vortex creep. In the case of elastic vortex creep, the elastic barrier at low J is $U_{el}(J) = U_c(J_c/J)^{\mu}$ [1], where U_c is the collective pinning barrier, J_c is the critical-current density, and $\mu \approx 1$. The evolution of the shape of the $U(J)$ curves in the elastic creep region shows that U_c increases between H_{on} and H_p. By further decreasing J, the relatively strong $U(J)$ dependence is replaced by a slower $U(J)$ variation. This is characteristic for plastic vortex creep, since U_{pl} has a weak dependence on J at low J. The crossover takes place when $U_{el} = U_{pl}$.

For $H > H_p$, only the weak $U(J)$ variation was observed (Fig. 4b). It is worth noting that instead of a slow increase of U with decreasing J, as expected for plastic vortex creep, a weak decrease of U with decreasing J appears. This is because the relaxation of M_{irr} leads to an increase of B inside the sample, and the plastic barrier is a decreasing function of B.

As shown above, recent models for the destruction of the low-field Bragg glass induced by the quenched disorder [20–24] are based on the competition between the elastic energy of the vortex system and the pinning energy, in the temperature interval where the thermal energy is negligible. By equating the elastic energy and the pinning energy, it was possible to describe the $H_{on}(T)$ variation [27,31]. In a similar way, the equality between the collective pinning barrier and the plastic barrier at H_p leads, in the case of less anisotropic HTSC, to a $B_p(T)$ dependence in good agreement with that experimentally determined [34]. The observed crossover from elastic to plastic vortex creep across the SMP of HTSC does not contradict these models, but it emphasizes the significant role of the proliferation of defects in the vortex system above H_{on}.

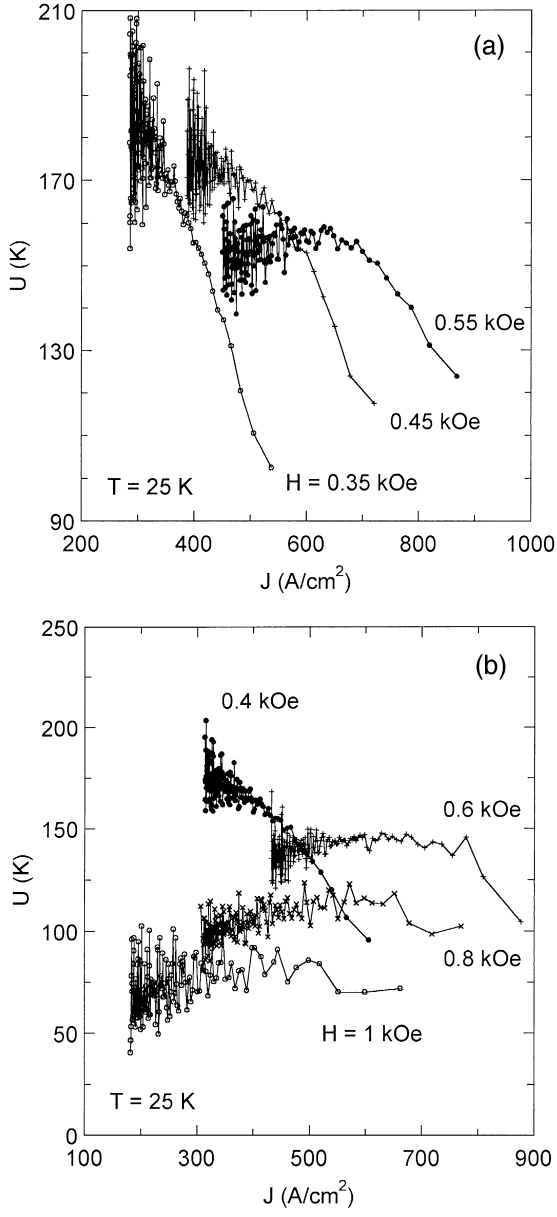

Fig. 4. (a) The activation energy U determined with Eq. (1) vs. the current density J for several H values between H_{on} and H_p at $T = 25$ K. In this H domain, there is a first rapid increase of U with decreasing J, followed by a weaker $U(J)$ variation as J is further decreased. (b) Current density dependence of the activation energy across the second magnetization peak at $T = 25$ K. Above H_p, which is of the order of 0.6 kOe (Fig. 1), only a weak $U(J)$ variation can be observed.

The crossover from elastic to plastic vortex creep across the SMP appears to be a general behavior [30,31,34]. This implies that the mechanism for the occurrence of the SMP in highly anisotropic HTSC, such as BSCCO, and in less anisotropic HTSC, like $Y_1Ba_2Cu_3O_{7-\delta}$, should be the same, related to the increase of the effective pinning for a softer vortex system. The existing differences are generated by the fact that in the case of highly anisotropic HTSC, B_p is limited by the crossover field $B_{cr} \approx \Phi_0/\gamma^2 s^2$, where γ is the anisotropy parameter and s is the distance between the superconducting layers. Around this field value, the elastic moduli of the vortex system suddenly decrease [1]. This could explain the very sharp onset of the SMP observed in local magnetic field measurements performed on relatively clean BSCCO single crystals, as well as the weak $B_p(T)$ dependence [25].

In conclusion, the current density and magnetic field dependence of the activation energy in the magnetization relaxation of slightly overdoped BSCCO single crystals grown by the self-flux method reveals a crossover from elastic to plastic vortex creep at the SMP. Above the SMP, the vortex system behaves like a plastic vortex solid, with $U(B) \propto B^{-1/2}$ and a weak intrinsic $U(J)$ variation. This could have important consequences on the nature of the thermally induced vortex solid-vortex fluid transition above H_p. It was found that the activation energy in the vortex fluid domain of slightly overdoped BSCCO films obeys the same $U(B)$ dependence [35]. These results suggest that above H_p, the thermally induced vortex phase transition should be consistent with the freezing of a viscous vortex fluid into a plastic vortex solid.

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