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Magnetization relaxation in the flux-creep annealing regime across the second magnetization peak of disordered $YBa_2Cu_3O_{7-x}$ crystals

L. Miu [a,*], C.C. Almasan [b], D. Miu [c], H. Adrian [d]

[a] *National Institute of Materials Physics, P.O. Box MG-7, Bucharest-Magurele, Romania*
[b] *Department of Physics, Kent State University, Kent, OH 44242, USA*
[c] *National Institute of Laser, Plasma, and Radiation Physics, Bucharest MG-6, Romania*
[d] *Department of Physics, University of Mainz, D-55099 Mainz, Germany*

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Abstract

The relaxation of the irreversible magnetization of disordered $YBa_2Cu_3O_{7-x}$ crystals measured in the “flux-creep annealing” regime reveals that across the second magnetization peak (SMP) the barriers against flux motion remain finite at low current densities, which supports the existence of a crossover to a dissipation process involving the plastic deformation of the vortex system. In our experiments, the vortex creep process appears to be exclusively controlled by collective pinning barriers (diverging at low current densities) only below the onset of the SMP, where the vortex system is stable against dislocation formation. The (elastic) collective pinning barriers observed for magnetic field values close to the onset of the SMP (where the plastic barriers are high) could be related to the recently proposed collective pinning of individual dislocations. The proliferation of dislocations across the SMP leads to liquid-like behavior of the disordered vortex phase in the vicinity and above the peak field.

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Keywords: High-temperature superconductors; Magnetization relaxation; Flux-creep annealing; Elastic creep; Plastic creep

1. Introduction

Over the last few years it has become apparent the existence of at least two vortex phases in high-temperature superconductors (HTS): the vortex solid and the vortex liquid [1]. In “clean” HTS single crystals, the vortex lattice at low magnetic fields undergoes a well-documented thermally induced first-order melting transition [2–4]. When the random quenched disorder becomes significant, in the region that is assumed to correspond to the vortex solid there occurs a transition between a low field quasi-ordered vortex phase (the Bragg glass, stable against dislocation formation [5,6]) and a disordered vortex phase at higher fields [7–12]. The disordered vortex phase (in which a better accommodation of vortices to the pinning centers is expected) is characterized by an enhanced apparent critical-current density J_c, leading to the well-known second magnetization peak (SMP) on the magnetization curves of HTS [13–19]. In highly disordered HTS, the SMP is

* Corresponding author. Fax: +40-1-493-0267.
E-mail address: elmiu@alpha1.infim.ro (L. Miu).

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observed up to very close to the critical temperature T_c, and the peak field is shifted to lower values [20].

The dynamic behavior of the disordered vortex phase above the onset of the SMP is still unclear. This could behave as an elastic vortex glass, where the pinning barriers diverge when the current density $J \to 0$, and, consequently, the linear electrical resistivity vanishes. The collective pinning theory [21] gives for $J \ll J_c$ an elastic barrier $\propto (J_c/J)^{\mu}$, with the collective pinning exponent $\mu > 0$. In this case, one should expect a thermally induced elastic vortex glass–vortex liquid transition [22] in the vicinity of the irreversibility line, as often reported for disordered HTS, based on the scaling of the current–voltage characteristics in agreement with the vortex–glass transition theory [23–26]. Alternatively, the dissipation processes involving the plastic deformation of the vortex system can be dominant. This was first reported in Ref. [27], for $YBa_2Cu_3O_{7-x}$ (Y:123) single crystals, and appears to be a general behavior [28–30].

A large amount of plastic vortex creep, characterized by non-diverging barriers, would be in conflict with the existence of a thermally induced elastic vortex glass–vortex liquid transition (predicting zero linear resistivity in the limit $J \to 0$ for any temperature T below the glass temperature) located well above the onset of the SMP [31].

In this work, we address the above issue through magnetization relaxation experiments at low J levels. The J dependence of the pinning barriers involved in the dissipation process is given by the activation energy $U(J)$ in the relaxation of the irreversible magnetization M_{irr}, if the extrinsic contribution to $U(J)$ is small. We analyzed the magnetization relaxation in the flux-creep annealing regime [32] across the SMP of highly disordered Y:123 crystals. Our results suggest that the creep process can be exclusively controlled by collective pinning barriers (diverging for $J \to 0$) only below the onset of the SMP, where the vortex system is stable against dislocation formation.

2. Experimental results and discussion

Single-grain fully oxygenated Y:123 specimens were cut from a bulk material obtained by melt texturing, containing $\approx$15% Y_2BaCuO_5 precipitates. The samples, similar to those investigated in Ref. [33], [1] are suitable for flux-creep annealing experiments, since the effect of the surface barriers is small, not very close to the irreversibility line. The presented results are for a $3 \times 3 \times 1.5$ mm^3 crystal (with the smaller dimension along the crystallographic c axis), having $T_c \approx 91$ K and a magnetically determined transition width of $\approx$0.5 K. The in-plane electrical resistivity at $T = 95$ K is relatively high ($\approx$0.8 mΩ cm), due to the presence of quenched point disorder, growth defects, Y_2BaCuO_5 precipitates, and randomly oriented twin boundary structures.

The magnetization M was measured using a Quantum Design SQUID magnetometer, with the external magnetic field H oriented along the c axis. The sample was first cooled in $H = 0$ from above T_c up to a certain temperature T_A, and then H was applied. In our flux-creep annealing experiments, the sample was warmed up to a temperature T_B, and then cooled back to T_A, where $M(t)$ was measured. The thermal cycle took 15–25 min, and the relaxation time t was considered to be zero at the moment when the magnet charging was finished. The procedure was repeated for a higher T_B (which was increased with a step of 0.3 K) after the sample was heated up to $T > T_c$. T_A was stabilized with an accuracy of $\approx$5 mK.

$M_{irr}(H, T_A)$ was extracted from the magnetic hysteresis curves $M(H, T_A)$ by the standard procedure [34]. As a rule, the first data point on the $M(t)$ and $M(H)$ curves was taken $t_1 = 100$ s after H was applied, to avoid the influence of flux redistribution in the initial stage of the relaxation process [35]. $M_{irr}(t)$ is the measured $M(t)$ curve shifted by $M_{irr}(t_1) - M(t_1)$.

$M_{irr}(H)$ at $T = 80$, 82.5, and 85 K, revealing the increase of $|M_{irr}|$ between the onset field H_{on} and the peak field H_p, is shown in Fig. 1. The flux-creep annealing procedure was applied for $H > H_{on}$, at $T_A = 75$ and 80 K. The results are very similar, and below are presented those obtained across the

[1] The Y-123 crystals investigated by us were grown by C. Beduz at the Institute of Cryogenics, Southampton.

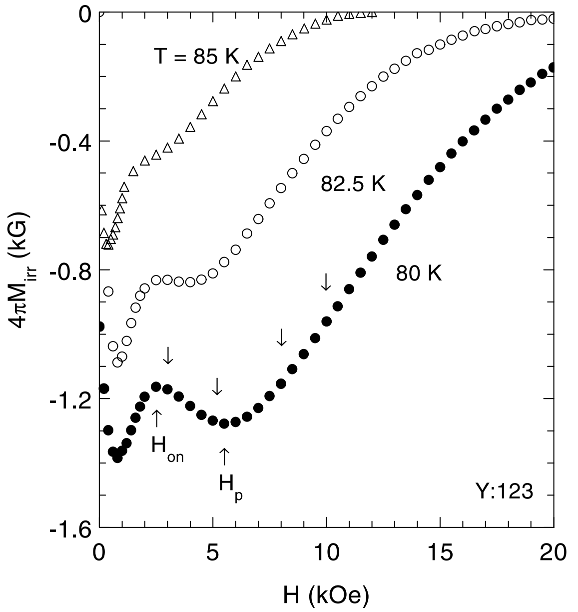

Fig. 1. Magnetic field H dependence of the irreversible magnetization M_{irr} of highly disordered $YBa_2Cu_3O_{7-x}$ (Y:123) crystals, revealing the appearance of the SMP at high temperatures T. The onset field H_{on} and the peak filed H_p are indicated by arrows, as well as the H values for which the flux-creep annealing experiments were performed.

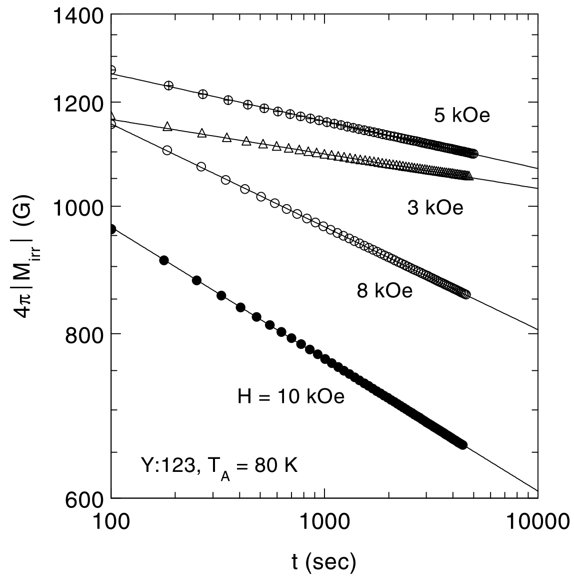

Fig. 2. Absolute value of the irreversible magnetization $|M_{irr}|$ at moderate relaxation levels vs. time t in a log–log plot at $T_A = 80$ K, for the H values across the SMP indicated in Fig. 1. In double logarithmic scales, a straight line can fit the data.

SMP at $T_A = 80$ K ($H = 3$, 5, 8, and 10 kOe), with $T_B \leqslant 82.1$ K.

The annealing temperature T_B is limited by several factors, such as the T stability at T_A (in the case of very low relaxation rates), some possible history effects (due to the $H_p(T)$ variation, for example), and by the influence of the surface barriers close to the irreversibility line. We first investigated the variation of an activation energy U_0 determined at moderate relaxation levels as $-T\,\mathrm{d}\ln(|M_{irr}|)/\mathrm{d}\ln(t)$ with the magnetic induction $B \approx H + 4\pi M(1-D)$ (where $D \approx 0.64$ is the demagnetization factor), for $T = 77.5$, 80, and 82.5 K, and $B \leqslant 1.5$ T. No qualitative change in the $U_0(B)$ dependence was observed, which rules out a significant contribution of the surface barriers for $T_B \leqslant 82.5$ K.

The relaxation of M_{irr} at T_A, without flux-creep annealing, is illustrated in Fig. 2. In logarithmic scales, a straight line can fit the data. With the general equation $U(J) = T\ln(t/\tau)$ [34] (where τ is the "effective" hopping attempt time) and $J \propto |M_{irr}|$, this shows that $U(J)$ obeys the logarithmic model, $U(J) = U_0 \ln(J_c/J)$ [36]. The linear fit gives U_0 (8 kOe, 80 K) =1040 K, for example.

In the case of a relaxation curve obtained after flux-creep annealing at T_B, the "equivalent relaxation time at T_A" t_{eq} can be correctly determined if $U(J)$ remains essentially the same at lower J, by assuming the continuity of the magnetization data [34]. However, as can be seen in Fig. 3(a), when J is lowered, the relaxation of M_{irr} becomes faster than that predicted by the linear fit of the $M_{irr}(t)$ data at moderate relaxation levels in the representation from Fig. 2, if H is higher than H_p. For H close to H_{on} (Fig. 3(b)), by decreasing J, the relaxation of M_{irr} is first slower than that predicted by the linear fit from Fig. 2, and then becomes faster than it. Similar results are obtained when the fit of $M_{irr}(t)$ with the interpolation formula from the collective pinning theory [34] is used. This supports the occurrence of a crossover from elastic creep to a creep process involving the plastic vortex deformation [27–30].

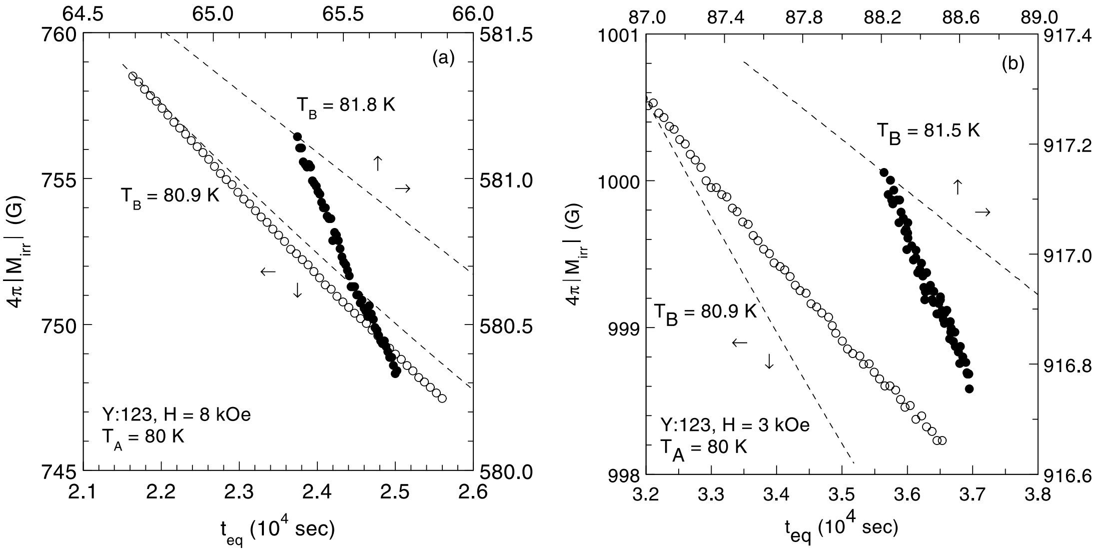

Fig. 3. $|M_{irr}|$ measured at $T_A = 80$ K after flux-creep annealing at a temperature T_B vs. the "equivalent relaxation time at T_A", $t_{eq} = t + t_0$. The offset time t_0 was determined by assuming the continuity of the magnetization data. The relaxation curve measured after flux-creep annealing was translated along the time axis so that the first data point to lie on the dashed line representing the fit of the data obtained without flux-creep annealing (Fig. 2) extrapolated to longer time values. (a) $H = 8$ kOe $> H_p$, $T_B = 80.9$ and 81.8 K; (b) $H = 3$ kOe $< H_p$, $T_B = 80.9$ and 81.5 K.

The above behavior implies the deviation of the $U(J)$ dependence at low J from the logarithmic variation usually observed at moderate relaxation levels, and reinforces the idea that extrinsic factors (such as the barrier distribution and/or the spatial distribution of J_c) contribute to the upward curvature of $U(J)$ for J close to J_c [37–40]. For this reason, we considered the flux-creep annealing procedure more appropriate than the method proposed in Ref. [41], for the investigation of $U(J)$ at low J. This is because the latter involves relaxation data at different T for J not sufficiently far from $J_c(T)$, and the extrinsic contribution to $U(J)$ will be seen in the whole J range. In the flux-creep annealing experiments, the intrinsic $U(J)$ dependence should appear at low J, since the above distribution is expected to be cut off at a certain value.

From the magnetic relaxation data, the electric field–current density (E–J) characteristics at low E levels were obtained. The electric field at the sample surface $E \approx \pi r(1 - D)\,\mathrm{d}|M_{irr}|/\mathrm{d}t$, where r is the characteristic sample size perpendicular to H, and J is extracted from $|M_{irr}|$ with the Bean model [34].

The resulting E–J curves across the SMP at $T_A = 80$ K are plotted in Fig. 4, in log–log scales. The logarithmic $U(J)$ variation at high J leads to $E \propto J^n$, where $n = U_0/T \gg 1$. The intersection of the $E(J)$ curves for $H = 3$ and 5 kOe reflects the interesting time evolution of the SMP, due to different relaxation rates below and above H_p.

The important feature is the change of the shape of the E–J curves at low E levels across the SMP. The obvious upward curvature of the E–J curves around and above H_p indicates the trend to thermally assisted flux flow at low J [42]. The downward curvature appearing close to H_{on} (for $E > 10^{-12}$ V/m, Fig. 4) signals the presence of (elastic) collective pinning barriers. Fig. 4 illustrates the crossover from diverging barriers at low J for $H < H_{on}$ to non-diverging barriers for $H > H_{on}$. (It should be noted that H_{on} is well below the elastic vortex glass transition line usually reported from the scaling of the E–J curves obtained in transport measurements, with a limited voltage sensitivity.)

The fit of the $E(J)$ data for $H = 3$ kOe in the range $10^{-12} \leqslant E \leqslant 10^{-10}$ V/m gives a collective

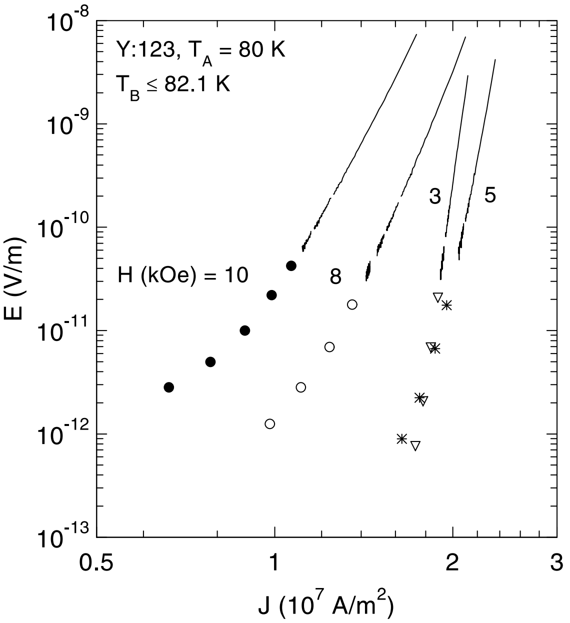

Fig. 4. E–J curves (double logarithmic scales) resulting from the relaxation of the irreversible magnetization across the second magnetization peak at $T_A = 80$ K with the annealing temperature $T_B \leqslant 82.1$ K. For a large $T_B - T_A$ difference, the plotted $E(J)$ (symbols) are averaged values, whereas for a small $T_B - T_A$ difference the whole data set is shown by a continuous line.

pinning exponent $\mu \approx 1$, suggesting that the collective pinning barriers just above H_{on} are related to the collective pinning of individual dislocations [43,44]. (The effect of the correlated disorder in our samples seems to be diminished due to the presence of a large amount of random quenched disorder, in agreement with Ref. [45].) The liquid-like behavior of the vortex assembly in the vicinity and above H_p would then appear as the result of the proliferation of dislocations across the SMP. As pointed out in Refs. [43,44], when the dislocation density becomes of the order of a^{-2}, where a is the mean inter-vortex spacing, the disordered vortex solid and the vortex liquid phase are thermodynamically indistinguishable.

3. Conclusions

The magnetization relaxation experiments in the flux-creep annealing regime performed across the SMP of disordered Y:123 crystals at relatively high T indicate that the barriers against flux motion remain finite at low J. The (elastic) collective pinning barriers observed in our experiments close to the onset field could be related to the recently proposed collective pinning of individual dislocations. However, even in this field range, below a certain J value, the creep process is also controlled by non-diverging barriers (involving the plastic vortex deformation), due to the co-existence of the Bragg glass and the disordered vortex phase. The creep process appears to be exclusively controlled by elastic barriers (diverging in the limit $J \to 0$) only below the onset of the SMP, where the vortex system is stable against dislocation formation. The proliferation of dislocations across the SMP leads to liquid-like behavior of the disordered vortex phase in the vicinity and above the peak field. Due to a large amount of plastic vortex creep, an elastic vortex glass–vortex liquid transition close to the irreversibility line becomes inappropriate. At least in the high-field region, as recently proposed in Ref. [46], a scenario more like a conventional glass transition merits further consideration [47].

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