Origin of the plateau in the temperature dependence of the normalized magnetization relaxation rate in disordered high-temperature superconductors

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(Received 2 October 2008; revised manuscript received 10 November 2008; published 29 December 2008)

The temperature T dependence of the normalized magnetization relaxation rate S in optimally doped $YBa_2Cu_3O_{7-\delta}$ films with the external dc magnetic field H oriented along the c axis exhibits the well-known plateau in the intermediate T range, associated with the presence of elastic (collective) vortex creep. The disappearance of the $S(T)$ plateau in the high-H domain ($H \geq 20$ kOe) is not completely understood. We show that in the case of high-temperature superconductors with significant quenched disorder the $S(T)$ plateau is directly related to a crossover in the vortex-creep process generated by the macroscopic currents induced in the sample. In dc magnetization measurements the creep-crossover temperature decreases rapidly with increasing H, reaching the low-T region where the magnetization decay is dominated by micro flux jumps. Consequently, at high H no well-defined elastic-creep domain is present and the $S(T)$ plateau disappears.

DOI: 10.1103/PhysRevB.78.212508 PACS number(s): 74.78.Bz, 74.72.Bk, 74.25.Qt

The relaxation of the irreversible magnetization of high-temperature superconductors (HTSs) remains an essential tool for the investigation of vortex dynamics and the vortex phase diagram. In order to reduce the intrinsic ambiguity of flux-creep measurements,[1,2] many magnetization relaxation studies have focused on the analysis of the normalized magnetization relaxation rate. At low T, where the relaxation is weak, one usually determines a normalized magnetization relaxation rate averaged over a fixed relaxation-time window, $S=-\Delta \ln(|M|)/\Delta \ln(t)$, where M is the irreversible magnetization (proportional to the current density J of the macroscopic currents induced in the sample) and t is the relaxation time.

In the case of HTSs with relevant random quenched disorder, the $S(T)$ variation at relatively low H exhibits three distinct regions.[2–9] At high T, $S(T)$ increases with increasing T. At intermediate T values a plateau in $S(T)$ develops, whereas at low T a slight decrease in S with decreasing T appears. The rapid $S(T)$ increase in the high-T region is commonly attributed to thermal fluctuations, whereas the $S(T)$ behavior at low T is associated with quantum vortex creep[10] or with a crossover toward an exponential time dependence of the irreversible magnetization.[11] The apparent universality of S in the plateau region (with values clustered around $S \sim 10^{-2}$) was explained in Ref. 3 by assuming the existence of an elastic vortex glass,[12] with strongly nonlinear current-voltage characteristics. It was shown[3] that the S universality comes from the fact that in the elastic vortex-glass domain S depends on the collective (elastic) creep exponent[13] and only logarithmically on time parameters (which do not change very much with the system and experiment). When the disorder degree is weak, the above interpretation remains essentially the same, owing to the well-established collective pinning behavior in the Bragg glass domain at low H.[14]

For disordered HTSs, when the $S(T)$ plateau is associated with the existence of an elastic vortex glass a question which arises is related to the disappearance of the $S(T)$ plateau at high H. In the case of optimally doped $YBa_2Cu_3O_{7-\delta}$ (Y-123) films, for example, the elastic vortex-glass transition temperature reported from the scaling of current-voltage characteristics remains above 70 K for H as high as 100 kOe,[15] whereas the plateau disappears above a significantly lower H.[7]

In this work, we address the above issue through a detailed analysis of magnetization relaxation in high-quality Y-123 films. It is shown that the appearance of the plateau in $S(T)$ is directly related to a crossover in the vortex-creep process generated by the macroscopic currents induced in the sample during magnetization measurements: plastic creep at high T (low J) and elastic creep at intermediate T (J). The $S(T)$ plateau corresponds to the elastic-creep regime, in agreement with Ref. 3. However, in the case of dc magnetization measurements the creep-crossover temperature decreases strongly with increasing H, reaching the low-T domain where the magnetization decay is highly influenced by the occurrence of micro flux jumps. At high H no well-defined elastic-creep domain is thus present and the $S(T)$ plateau disappears.

We investigated the magnetization relaxation for disk-shaped (2.5 mm in diameter) optimally doped Y-123 films (~250 nm thick) prepared by high-pressure dc sputtering onto (001)-oriented $SrTiO_3$ substrates. The inductively mea-

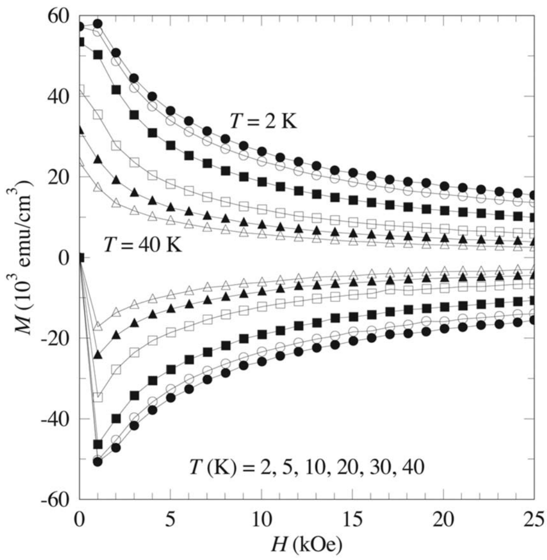

FIG. 1. Characteristic dc magnetization curves $M(H)$ of disk-shaped optimally doped Y-123 films ($T_c \sim 91.5$ K) for T between 2 and 40 K. One can see a limitation of M at low T.

sured critical temperature T_c is around 91.5 K for $H=0$, and the transition width is of the order of 0.3 K. The dc magnetization M was measured using a commercial Quantum Design magnetic property measurement system (MPMS), with H oriented along the c axis and always applied in zero-field-cooling conditions. In the considered (H,T) domain M was identified with the irreversible magnetization. The relaxation time t was taken to be zero when magnet charging was finished, and the first data point was registered at $t_1 \sim 100$ s. The remnant state was created by applying a magnetic field $H=10$ kOe at $T>T_c$, then the desired $T<T_c$ was stabilized, and finally H was decreased to zero. The H sweeping rate in our measurements was of ~ 100 Oe/s.

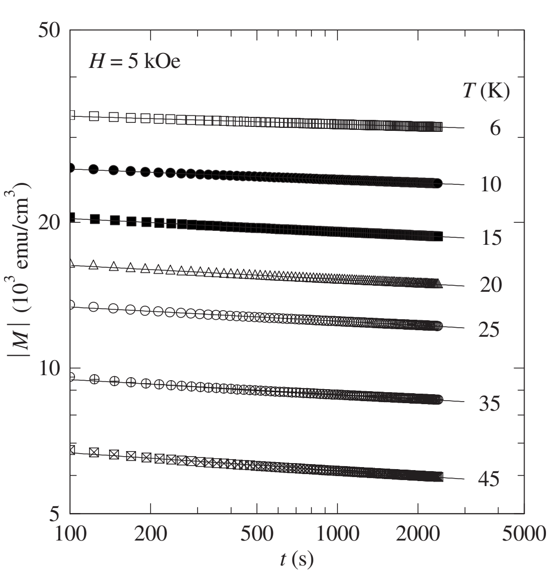

FIG. 2. Characteristic magnetization relaxation curves $M(t)$ in double-logarithmic scales for $H=5$ kOe and several T values well below T_c.

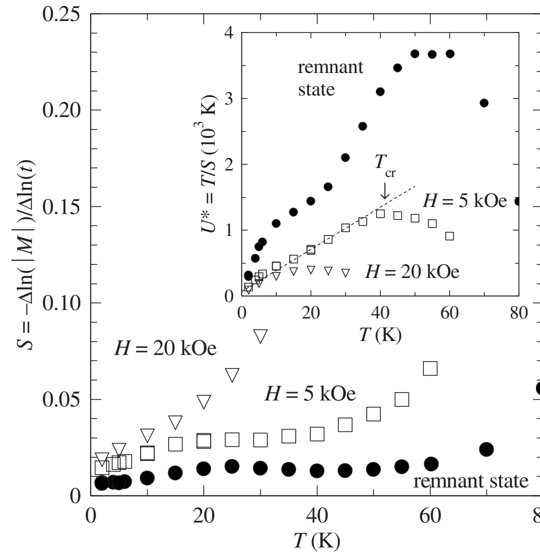

FIG. 3. Temperature variation in the normalized magnetization relaxation rate $S=-\Delta \ln(|M|)/\Delta \ln(t)$ averaged over the relaxation-time window from Fig. 2. At low H (5 kOe and remnant state) the $S(T)$ plateau develops above ~ 20 K, but it disappears for $H=20$ kOe. The inset illustrates the corresponding normalized vortex-creep activation energy $U^*(T)=T/S$, showing a maximum at the crossover temperature T_{cr} (indicated by an arrow). The T interval for the $S(T)$ plateau is located just below the crossover temperature T_{cr}. The linear $U^*(T)$ variation in the plateau region has a small intercept $U^*(0)$. The linear fit for $H=5$ kOe (dashed line) leads to $U^*(0)=70\pm 30$ K.

Figure 1 illustrates the dc magnetization curves $M(H)$ for T between 2 and 40 K. One can see the "limitation" of M at low T, which is attributed to the presence of micro flux jumps, as will be discussed below. Characteristic magnetization relaxation curves $M(t)$ in double-logarithmic scales are presented in Fig. 2, and we determined the normalized magnetization relaxation rate S averaged over the (fixed) relaxation-time window t_w (100–2500 s) and the related normalized vortex-creep activation energy $U^*=-T\Delta \ln(t)/\Delta \ln(|M|)=T/S$.

The $S(T)$ variation for $H=5$ and 20 kOe and for the remnant state is plotted in the main panel of Fig. 3, whereas the corresponding $U^*(T)$ is shown in the inset. For the remnant state and $H=5$ kOe the $S(T)$ plateau develops above ~ 20 K. The T domain for the $S(T)$ plateau shrinks with increasing H and the plateau disappears for $H=20$ kOe. The T interval for the $S(T)$ plateau is located just below the crossover temperature T_{cr} where $U^*(T)$ exhibits a maximum. The linear $U^*(T)$ in the plateau region has a small intercept $U^*(0)$.

The discussion of the results from Fig. 3 first requires the precise meaning of U^*, which is sometimes called "the effective pinning energy." While U^* cannot actually be identified with the effective pinning barrier, this is very useful for detecting changes in the vortex-creep process. Following Ref. 16, the actual vortex-creep activation energy U can be written as

$$U(T,H,J) = (U_c/p)[(J_c/J)^p - 1], \tag{1}$$

where U_c is a characteristic pinning energy, whereas exponent p is identified with the (positive) collective pinning exponent μ in the case of elastic (collective) vortex creep,[13] and $p<0$ for plastic creep. With Eq. (1) and $J \propto |M|$ one can derive $U^*(J)$ using the general creep relation[17] $U = T\ln(t/t_0)$, where t_0 is a macroscopic time scale for creep[1] or the microscopic vortex hopping "attempt" time.[3] For the elastic-creep domain one obtains

$$U^*(J) = U_{ce}(J_c/J)^{\mu}, \tag{2}$$

where U_{ce} is the characteristic pinning energy for elastic creep. In the plastic creep regime U^* has an opposite variation with J,

$$U^*(J) = U_{cp}(J_c/J)^{p}, \tag{3}$$

where U_{cp} is the characteristic pinning energy for plastic creep. With a fixed t_w and T well below T_c, the above equations lead to

$$U^* \approx U_{ce} + \mu T\ln(t_w/t_0) \tag{4}$$

for elastic creep and

$$U^* \approx U_{cp} - |p| T\ln(t_w/t_0) \tag{5}$$

in the case of plastic creep. (For simplicity, we have considered the same t_0 for elastic creep and plastic creep.) The maximum in $U^*(T)$ appearing in the inset of Fig. 3 indicates a crossover elastic creep at low T-plastic creep at high T, and, in the approximation t_0=const, Eqs. (4) and (5) lead to a creep-crossover temperature $T_{cr} \approx (U_{cp}-U_{ce})/[(\mu + |p|)\ln(t_w/t_0)]$.

Since for $H \geq 5$ kOe the observed T_{cr} is much lower than T_c, the $U^*(T)$ decrease above T_{cr} cannot be primarily attributed to thermal fluctuations. It was shown in Ref. 18 that the creep crossover in dc magnetization measurements [leading to the nonmonotonic $U^*(T)$ from the inset of Fig. 3] is in fact caused by the macroscopic currents induced in the sample $J(T,t)$. At least for T well below T_c, the main role of the thermal energy is to change the probed J domain. This is due to a different overall relaxation in the interval between the moment when magnet charging was finished ($t \approx 0$) and the time t_1 at which the first data point is taken. When T decreases J increases toward J_c, and below a certain T_{cr} the effective pinning energy becomes lower than the elastic energy E_{el} in the vortex system. In this situation the creep is entirely elastic since some dislocations may heal[19] and the others will be collectively pinned[20] at such high drives. Simple arguments[21] (based on the dynamic energy balance relation[18] $U \propto E_{el} \propto H^{-1/2}$) lead to $T_{cr}(H) \propto H^{-1/2}$. With increasing T above T_{cr} the probed J interval is shifted down relative to J_c, and Eq. (3) indicates a decrease (increase) in $U^*(S)$. When this combines with the effect of thermal fluctuations at high T, a rapid $S(T)$ increase above the $S(T)$ plateau appears.[7]

Figure 4 illustrates U^* vs $1/J$, where J was extracted with the Bean model[22] from $M(t)$ averaged over t_w. For the remnant state and $H=5$ kOe a well-defined elastic-creep domain is present, where $U^*(J)$ is given by Eq. (2), with widely

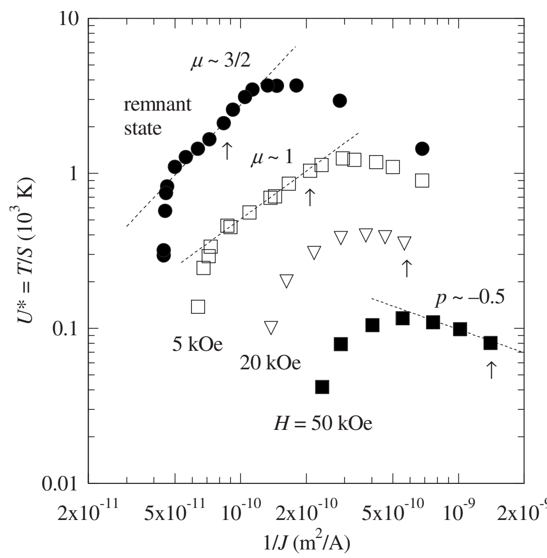

FIG. 4. The normalized vortex-creep activation energy U^* vs $1/J$, where J was extracted with the Bean model from $M(t)$ averaged over the relaxation-time window from Fig. 2. For the remnant state and $H=5$ kOe [for which the plateau in $S(T)$ develops (Fig. 3)] a well-defined elastic-creep domain is present, where $U^*(J)$ is given by Eq. (2), with widely accepted elastic-creep exponent μ values for small (remnant state) and intermediate ($H=5$ kOe) vortex bundle regimes. At $H=20$ and 50 kOe no elastic-creep domain was observed and the $S(T)$ plateau disappears. At high J the plastic-creep regime sets in and the plastic-creep exponent is close to the value proposed in Ref. 23 ($p=-0.5$). The dashed lines represent the linear fit in the log-log plot, and the arrows indicate $T=30$ K.

accepted μ values.[13] This is in agreement with the linear $U^*(T)$ variation from Eq. (4), where $U_{ce}=U^*(0)$. The intercept determined for $H=5$ kOe in the inset of Fig. 3 indicates that U_{ce} is small, as expected for elastic creep (where the pinning centers do not accommodate vortices). Neglecting U_{ce}, Eq. (4) gives $S(T) \approx T/U^* \approx [\mu \ln(t_w/t_0)]^{-1}$, which is the result for the universal S in the plateau region from Ref. 3. By considering the S value in the plateau region for $H=5$ kOe, $\mu \sim 1$, and t_w of the order of 10^3 s, one obtains t_0 of the order of 10^{-11} s, an accepted value for the microscopic vortex hopping attempt time.[1,3] For the plastic-creep domain above T_{cr}, Eq. (3) suggests a creep exponent p not far from -0.5 (the value proposed in Ref. 23). Equation (5) shows that it is impossible for an $S(T)$ plateau to appear in the plastic-creep regime, where the characteristic pinning energy $U_{cp} > |p| T\ln(t_w/t_0)$ (for vortices better accommodated to the pinning centers at $J \ll J_c$).

In the low-T domain a rapid decrease in U^* with increasing J occurs (Fig. 4). $M(t)$ cannot be related to elastic vortex creep since the resulting μ overcomes any plausible values[13] and/or is strongly J dependent. Together with the M limitation at low T (Fig. 1) this suggests the occurrence of micro flux jumps, preceding the avalanches expected at even lower T and/or a higher field-sweeping rate.[24] For $H=20$ kOe the creep-crossover temperature T_{cr} already reaches 20 K [the end T value for the $S(T)$ plateau], which means that at high H the micro flux jumps set in as soon as the creep becomes elastic. In such a situation, no well-defined elastic-creep domain is present and the $S(T)$ plateau disappears.

Another indication of the presence of micro flux jumps is the upturn in $U^*(T)$ with decreasing T just below the $S(T)$ plateau, contributing to the decrease in S (see Fig. 3 and Ref. 7). In our opinion, this upturn is caused by a (continuous) crossover between elastic-creep-dominated magnetization relaxation and micro-flux-jump-dominated magnetization decay. In the upturn region the micro flux jumps mainly appear in the t interval between $\sim t_0$ and $\sim t_1$, where J is closer to J_c. The measured magnetization relaxation can still be dominated by elastic creep at $t \gtrsim t_1$, but the probed J interval was already shifted to lower values, leading to a higher U^* [see Eq. (2)]. This is somehow similar to what happens in the "flux-creep annealing" process, reducing the magnetization relaxation.[25] In these conditions, the use of pure thermally activated creep and the same t_w become inappropriate for T (H) below (above) the end point of the $S(T)$ plateau in the (H,T) diagram ($T \sim 20$ K and $H \sim 20$ kOe in the case of our Y-123 films). The micro flux jumps seem to be responsible for the existence of a finite S in the low-T limit, as an alternative to quantum vortex creep.

In summary, the existence of the plateau in $S(T)$ for disordered HTSs appears to be directly related to a crossover in the vortex-creep process generated by the macroscopic currents induced in the sample during magnetization measurements. The $S(T)$ plateau corresponds to the elastic-creep regime, in agreement with Ref. 3. In dc magnetization measurements the creep-crossover temperature decreases strongly with increasing H, reaching the low-T domain where the magnetization decay is influenced by the occurrence of micro flux jumps. At high H no well-defined elastic-creep domain is thus present and the $S(T)$ plateau disappears.

This work was supported by the MER at NIMP Bucharest and TU Cluj-Napoca and by the Alexander von Humboldt Foundation at the University of Mainz.

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