Nature of magnetic relaxation in a superconducting $Pr_{1.85}Ce_{0.15}CuO_{4-y}$ single crystal

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Magnetic relaxation in a $Pr_{1.85}Ce_{0.15}CuO_{4-y}$ single crystal has been studied as a function of field and temperature. Two different regimes are evident in the relaxation process. The usual logarithmic decay of the magnetization is observed at low fields. A crossover from logarithmic to exponential relaxation takes place at higher fields. The magnetization decay becomes exponential as H approaches $H_{irr}(T)$, and eventually any trace of relaxation disappears. The logarithmic decays have been used to estimate the field dependences of the critical currents and the relevant activation energies at high driving currents. The functional dependence $U(J)$ has been evaluated by using Maley's method. The extracted values for $U(J,T,B)$ are discussed in terms of the available models for flux motion; we conclude that the observed behavior might be a signature of a crossover in the nature of the thermally activated process, from jumps of elastically correlated flux volumes to plastic flux motion.

INTRODUCTION

Magnetic relaxation in type-II superconductors is related to the decay of the persistent currents; it is thought to be originated by the thermally activated jumps of (bundles of) vortices over their pinning barriers, the so-called flux-creep process. Anderson and Kim[1] showed that, when the effective potential barrier U depends linearly on the driving current J and $U/kT \gg 1$, the magnetization should decay logarithmically with time. Beasley, Labusch, and Webb[2] observed this logarithmic decay and studied the relation between relaxation rate and effective pinning energy. They showed that a linear dependence of U on J is a good approximation when J is of the order of the critical current $J_c(T,B)$. This limit is usually achieved in experiments with classical superconductors. However, nonlogarithmic magnetic relaxations have been observed in high-temperature superconductors (HTSC), and even in some type-II alloys. This behavior cannot be easily described in the framework of the Anderson-Kim (AK) regime and several models have been proposed to account for it. All of these models are based on the observation that HTSC exhibit unusually large magnetic-relaxation rates, which are caused by higher measurement temperatures and lower effective potential barriers.[3] These high-relaxation rates imply that induced currents can decay to rather small fractions of the critical current $J_c(T,B)$ within the experimental time window, and thus new phenomena associated with this effect may become observable.

In the framework of vortex-glass or collective-creep theory[4,5] it is argued that the potential barriers grow in a highly nonlinear way when decreasing J, thus leading to a decay law $m(t)\alpha[\ln(t)]^{-1/\mu}$ for $t\rightarrow\infty$, where μ depends on the dimensionality of the flux line (bundle) as well as on temperatures, applied magnetic fields, and driving currents. Such nonlinear J dependences of U have already been observed in I-V curves[6] and magnetic-relaxation measurements,[7] where a $U\alpha\ln(J)$ relationship was found. Vinokur, Feigel'man and Geshkenbein[8] have shown that this logarithmic dependence is a good approximation to the collective-creep dependence $U\alpha(J/J_c)^{-\mu}$, for small values of μ. On the other hand, the TAFF[9] model predicts a crossover from the Anderson-Kim logarithmic decay to an exponential one at small driving currents J, associated with diffusive flux motion, provided U does not diverge when $J\rightarrow 0$.[10] It has been shown[11] that TAFF motion of defects in the flux-line lattice is likely to occur at low enough J, where the energy barriers associated with plastic motion (which do not depend on J) become smaller than the ones of the collective-creep or vortex-glass model.

The electron-doped superconductors, $Ln_{2-x}Ce_xCuO_{4-y}$, are significantly different from the other high-temperature superconducting cuprates.[12] Their critical temperatures and upper critical fields are much lower. Consequently, the in-plane coherence lengths are larger and the pinning energies are expected to be higher than in other HTSC. The activation barriers measured from ac susceptibility[13,14] and resistivity[15] in single crystals of this family are close to the values reported for $YBa_2Cu_3O_7$ (YBCO) single crystals, and somewhat higher than those for $Bi_2Sr_2CaCu_2O_8$ (BSCCO). Hence, as the measuring temperatures are much lower, the relaxation rates should be closer to those in classical type-II superconductors. This fact seems to be confirmed by the narrow reversible region in the H-T plane, which has been observed from both inductive and transport measurements.[13–16]

The study of magnetic relaxation in the electron-doped superconductors is interesting from two points of view. First, it provides another determination of the activation energies relevant in thermally activated flux motion, and their dependence on the driving current J. Consequently, information on the nature of the moving flux volumes, and thus about the characteristics of the mixed state, can be gained. Second, since the values of the upper critical field along the c-axis direction are $\mu_0 H_{c2}(0)\approx 6$ T (Refs.

15, 17, 18), all the points in the (T,H) diagram of these materials are experimentally accessible. This fact is in strong contrast with the case of the p-type HTSC cuprates, in which the very high values of $H_{c2}(0)$ make impossible the study of flux-motion processes near $H_{c2}(T)$ at low temperatures.

In this paper we report a study of magnetic relaxation, as a function of magnetic field and temperature, in a highly irreversible $Pr_{2-x}Ce_xCuO_{4-y}$ single crystal. The hysteresis loops do not show any reversible regime within our experimental resolution, and thus they indicate that the irreversibility line $H_{irr}(T)$ should lie very close to the upper critical field.[19] We have observed logarithmic relaxations at the lowest fields and temperatures, crossing to an exponential decay at higher fields and longer times. At fields close enough to $H_{irr}(T)$ no relaxation is observed. The implications of these facts, together with the height of the activation barriers and their dependence upon J, are discussed.

EXPERIMENT

The single crystal used in the present study was grown by the self-flux method[20] and has an estimated composition $Pr_{1.85}Ce_{0.15}CuO_{4-y}$. It is a platelet of dimensions $1.80\times0.95\times0.03$ mm^3, the shortest length being the one parallel to the c axis. The critical temperature was extracted from the onset of diamagnetism in the dc magnetization, measured after a zero-field cooling process (ZFC) at $H=1.6$ G. It is $T_c=19.2$ K, with a transition width of about 1.5 K.

The measurements were performed by means of a Quantum Design superconducting quantum interference device (SQUID) magnetometer with a scan length of 3 cm. The relaxation data were taken by cooling the crystal down to the desired temperature in zero field and then the field parallel to the c axis was applied. The first data point is taken typically 50 s after the application of the field. Fields ranging from $\mu_0H=0.1$ T up to $\mu_0H=2$ T were used and the magnetization was recorded for 3 h. The data analyzed here correspond to the irreversible regime, above the full penetration field $H_p(T)$, for temperatures $T=4.5$, 6, and 9 K. The boundaries of this zone were determined from magnetic hysteresis measurements.[19]

RESULTS

Figures 1(a) and 1(b) show some typical curves of magnetization decay at $T=4.5$ K, for different fields. Clearly, the data corresponding to $\mu_0H\leq1.2$ T are well described by a linear expression $m(t)\alpha\ln(t)$. At higher fields, even though an initial logarithmic decay is observed, a crossover to a nonlogarithmic decay becomes apparent. This second issue will be addressed later.

Let us concentrate now in the low-field regime. The deviations observed at $t\leq300$ s can be attributed to transitory effects or to the difficulty in defining properly the origin of time; in any case, these effects may be taken into account by adding a constant time t_i. This time determines the transient period after which the logarithmic relaxation is observed.[21] It is a macroscopic, nonintrinsic constant, related to the nonlinear flux diffusion through the sample, which essentially depends on initial conditions, particularly the field sweep rate $\dot{H}$ used to establish the measuring field. In Fig. 2 we have depicted some fits to the experimental data by using the equation,

$$m(t)=p\ln(t+t_i)+q\ , \tag{1}$$

where p and q stand, respectively, for the slope and a constant term in the logarithmic law describing the data. It can be appreciated that the addition of t_i provides a reasonable description of the first measured points. It turns out that in all the cases a value of about $t_i\approx30\pm10$ s is obtained. An approximate expression for this time constant is[21] $t_i\approx kT/[U_c(T,B)]\{[m_c(T,B)]/\dot{H}\}$, where $m_c(T,B)\equiv m(t=0,T,B)$ is the nonrelaxed irreversible magnetization, which is proportional to the critical current $J_c(T,B)$ in a critical-state model. $U_c(T,B)$ is the linear extrapolation to $J=0$ of the effective energy barrier at $J\approx J_c(T,B)$. Using $m_c(T,B)\approx2$ T and $U_c/kT\approx30$ (see below), a typical value $\mu_0\dot{H}=10^{-3}$ T/s provides an estimate of $t_i\approx70$ s, which is in reasonable accordance with the extracted values from the fits.

Assuming a flux-creep relaxation process in the critical state, and taking a linear dependence on J for the energy barrier, $U(T,B,J)=U_c(T,B)(1-J/J_c)$ (which is a

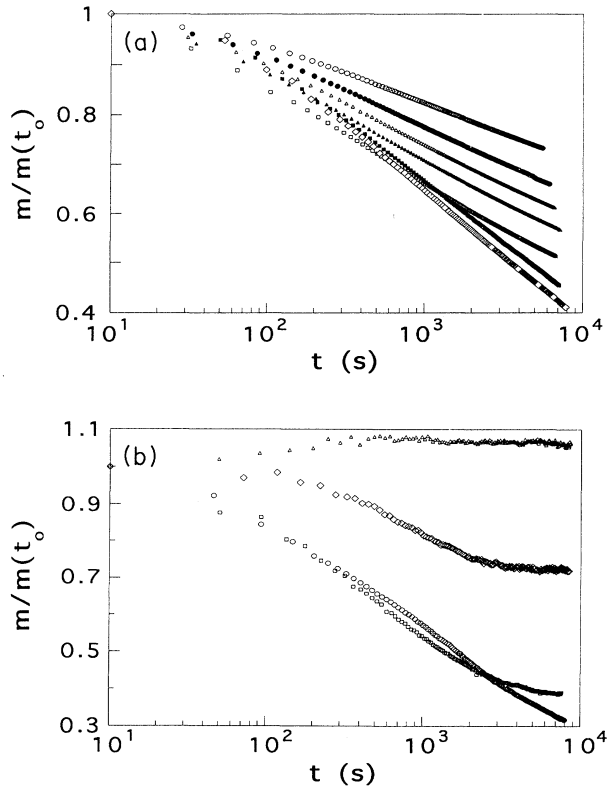

FIG. 1. Magnetization decays at $T=4.5$ K for different fields: (a) $\mu_0H=0.2$, (○), 0.3 (●), 0.4 (△), 0.5 (▲), 0.7 (□), 1 (■), and 1.2 (◇) T; (b) $\mu_0H=1.4$ (○), 1.5 (□), 1.7 (◇), and 1.9 (△) T. $m(t_0)$ is the first measured point at each field.

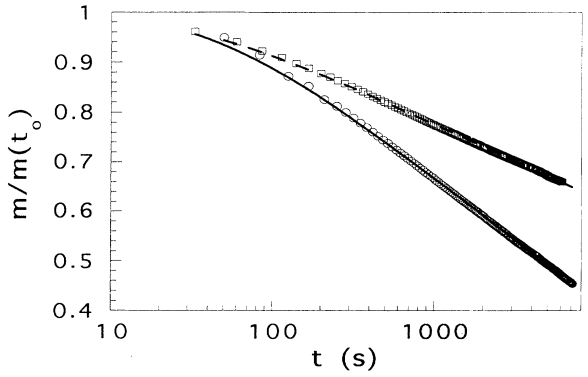

FIG. 2. Relaxation curves at $T=4.5$ K, for $\mu_0 H=0.3$ T (□) and 1 T (○); the lines are fits to Eq. (1) in the text.

reasonable approach when $U \gg kT$), Anderson and Kim[1] showed that magnetization should decay in time according to a law,

$$m(t,T,B)-m_{eq}(T,B) = m_c(T,B)\left[1-\frac{kT}{U_c(T,B)}\ln\left[\frac{t+t_i}{\tau}\right]\right], \quad (2)$$

where τ is the characteristic relaxation time and $m_{eq}(T,B)$ is the equilibrium magnetization. In a type-II superconductor $m_{eq}(T,B)$ is usually much smaller than $m_c(T,B)$, provided we are far enough from the reversible region; we will thus neglect this contribution[17,19] in the analysis of the low-field relaxations.

Comparing Eqs. (1) and (2), it may be appreciated that the fits shown in Fig. 2 provide two relationships between the three parameters involved in the relaxation process, namely, $m_c(T,B)$, $U_c(T,B)/kT$, and τ. Consequently, we have a set of two equations and three unknown factors.[22] In order to get a unique determination of these three factors, we need either a new relation between them or the value of one of these unknown factors evaluated from some other analysis. Since $\ln(\tau)$ is not expected to change appreciably with field and temperature, the evaluation of the normalization time τ seems to be a good starting point in order to achieve this goal. According to Feigel'man, Geshkenbein, and Vinokur,[23] a lower limit for this characteristic time can be obtained from

$$\tau^{-1}=\frac{2\rho_f}{\mu_0 d^2}\frac{U_c(T,B)}{kT}, \quad (3)$$

where $\rho_f \approx \rho_n B/B_{c2}$ is the flux-flow resistivity and d is a sample dimension. The method of Maley, to be described below, shows that there is a constant value $U_c(T,B)/kT \approx 30$ for all tested temperatures and fields. Estimates of τ using Eq. (3) with the above values for the activation energy, the normal-state resistivity in these materials $\rho_n=3\times10^{-5}$ Ω m (Ref. 15), an effective sample dimension $d=0.74$ mm, and the upper critical field $B_{c2}=6$ T (Refs. 17, 18), gives a value for the characteristic relaxation time $\tau \sim 10^{-9}-10^{-10}$ s. The nonlogarithmic relaxation data and the analysis of Maley also point to a value of τ of this order of magnitude, as will be shown later. We will thus take the value $\tau \sim 10^{-9}$ s as a fixed parameter in the study of the logarithmic fits, observed for $\mu_0 H \leq 1.2$ T.

Once the value of τ has been fixed, the fits of the experimental $m(t,T,B)$ data to Eq. (1) can provide a straightforward estimate of $U_c(B)$ and $m_c(B)$ at 4.5 K; results are shown in Fig. 3. The potential barrier shows a very smooth field dependence which can be roughly described by a power law $U_c(T,B) \approx U_0(T) B^{-0.15}$, with $U_0(T=4.5$ K, $B=1$ T$) \approx 13$ meV. The values of these energies are almost one order of magnitude smaller than those reported from ac-susceptibility and resistivity measurements on other electron-doped superconducting single crystals.[12–14] The field dependences of the activation barriers determined from these different methods are also in contrast with the one determined here. This discrepancy, common to other superconducting oxides, may be related to the fact that different driving currents are involved in the relaxation processes and ac-susceptibility or resistivity measurements. In the framework of a model with a highly nonlinear dependence of the effective energy barrier on J, such as those of vortex-glass or collective-creep models, the use of these different experimental techniques would imply that very different regions of the $U(J)$ curve are sampled in each experiment, and therefore, different results should be expected. We will deal with this issue later in the discussion.

An estimate for the irreversible magnetization at $t=0$, $m_c(T,B)$, can be obtained by taking $\tau \sim 10^{-9}$ s and comparing Eq. (2) with the fits to Eq. (1), as described above. Under these considerations, it may be seen in Fig. 3 that $m_c(T,B)$ exhibits an exponential decrease with magnetic field as follows:

$$m_c(T,B)=m_c(T,B=0)\exp(-B/B_0), \quad (4)$$

with $B_0 \approx 0.47$ T and $\mu_0 m_c(4.5$ K, $B=0) \approx 2.3$ T. Assuming a critical-state model, these $m_c(T,B)$ values can

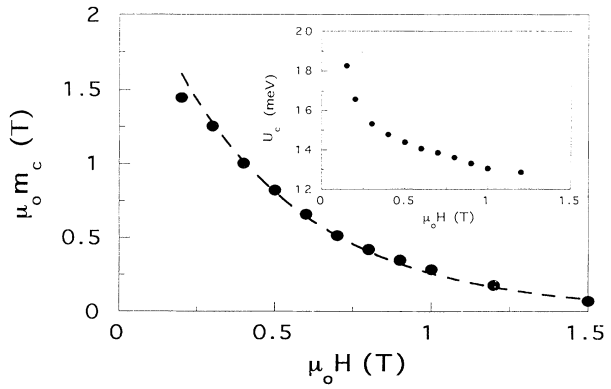

FIG. 3. Field dependence of the nonrelaxed irreversible magnetization m_c at 4.5 K, extracted from the fits to Eq. (1) with $\tau=10^{-9}$ s. The point corresponding to $\mu_0 H=1.5$ T, obtained from the logarithmic decay at short times, is also displayed. The solid line is a fit to $m_c \sim \exp(-B/B_0)$, with $B_0 \approx 0.47$ T. Inset: Field dependence of the activation barrier U_c at 4.5 K, extracted from the same fits.

be used to estimate the critical current of the sample. It turns out that $J_c(4.5\ \text{K}, 1\ \text{T}) \approx 10^8\ \text{A/m}^2$. The observed field dependence of $m_c(T,B)$ [Eq. (4)] is in good agreement with the field dependence of J_c extracted from hysteresis cycles[16] and with the exponential decay of the relaxation rate $S(B) \equiv dm/d\ln(t)$ observed in the same crystal data.[19]

We would like to emphasize the fact that the estimates of $m_c(T,B)$ are nearly twice as high as the first measured point in any run. This means that, at the first measurement, the magnetization has already decayed to an important fraction of its original value. Relaxation effects are thus rather important. It is well known that these effects mask the evaluation of the critical current (its value and temperature and field dependences) from the hysteresis curves, $J_c \alpha \Delta m$, by means of a critical state model.[24] Equivalently, the field dependence of the relaxation rate $S(B)$ should reflect the ratio of the field dependences of $m_c(B)$ and $U_c(B)$.[25] However, the exponential dependence on field in our estimates of $m_c(B)$ [Eq. (4)] is much stronger than the field dependence of the relevant activation energies (Fig. 3). Thus, it dominates the global field dependence of $\Delta m(B)$ and $S(B)$, which may be considered as a good approximation of the actual field dependence of the critical current in these electron-doped superconducting oxides.

We turn now to the high-field relaxation data, shown in Fig. 1(b). As mentioned above, at higher fields ($\mu_0 B \geq 1.4$ T for $T=4.5$ K) the magnetization does not show a logarithmic decay in the whole time window, but crosses to another regime of slower decay at a time t_{cr}, and eventually becomes constant at the highest fields and longest times. There are several theoretical frameworks in which a nonlogarithmic decay of magnetization may be expected. In the simplest case of an effective energy barrier displaying a linear dependence on J (Anderson-Kim model), the logarithmic relaxation would be left at the time when $m(t)$ reached a value much smaller than $m_c(T,B)$. In this situation, $U(J,T,B)$ could be considered as independent of J and the thermally activated flux motion would take a diffusive nature (TAFF regime). Magnetization would thus decay exponentially to its equilibrium value. The crossover time from logarithmic to exponential relaxation can be estimated by[10]

$$\tau_2 \approx \tau \exp[U_c(T,B)/kT] \ . \tag{5}$$

The assumption of a linear $U(J)$ implies that the elastic interactions between segments of vortices are not taken into account. However, the collective-creep and vortex-glass models[4,5] have shown that the elastic nature of the vortex system cannot be ignored. Incorporation of its effect into a thermally activated flux motion provides activation barriers with a power dependence on J as follows:

$$U(J,T,B) = U_c(T,B) \left[\frac{J}{J_c(T,B)} \right]^{-\mu} . \tag{6}$$

These enhanced effective barriers are the result of the fact that the jumping flux-bundle size and jumping distance u increase as $J \to 0$. Evidently, in the presence of potential wells of the type (6), the logarithmic decay predicted by the AK model is not valid any more; instead, magnetization should decay according to a law,

$$m(t,T,B) - m_{\text{eq}}(T,B) \approx \left[\frac{kT}{U_c(T,B)} \ln \left[\frac{t+t_i}{\tau} \right] \right]^{-1/\mu} , \tag{7}$$

for $t \to \infty$. The value of the exponent μ is predicted to increase from nearly zero at $J \approx J_c$ to $\frac{9}{8}$ or $\frac{3}{2}$ (depending on whether the flux-line lattice is 2D or 3D). Hence, the Anderson-Kim logarithmic decay, which assumes a linear dependence of $U(J)$, is a good approximation at high-current densities J, i.e., in the first stages of the relaxation. But as J/J_c decreases, the linear approximation becomes worse and the description of the data by means of Eq. (7) becomes more suitable. Experimentally, a dependence of the type,

$$U(J,T,B) = U_c(T,B) \ln \left[\frac{J}{J_c(T,B)} \right] , \tag{8}$$

has also been observed in HTSC,[6,7] which would imply that the relaxation process is governed by a law,

$$m(t,T,B) - m_{\text{eq}}(T,B) \approx \left[\frac{t+t_i}{\tau} \right]^{-kT/U_c(T,B)} . \tag{9}$$

Vinokur, Feigel'man, and Geshkenbein[8] have shown that the dependence (8) is a good approximation of (6) for small values of μ, i.e., in the stage that appears immediately after the Anderson-Kim regime. On the other hand, Eq. (6) predicts a divergence of $U(J)$ as $J \to 0$. This implies that no relaxation should be expected at low driving currents, and thus a nondissipative vortex-glass state would appear. Nevertheless, as has been pointed out in Ref. 10, this divergence can be cut off in a 2D flux-line lattice (FLL) or in a 3D vortex-fluid phase. In these cases, plastic excitations and defects, such as dislocations or bending, cutting, and reconnection of flux lines, are likely to appear in the vortex system. The energies $U_{\text{pl}}(T,B)$ associated with creation and motion of these defects are independent of the driving current J. Consequently, a crossover from collective creep to plastic motion (TAFF) is expected to take place as $J \to 0$, at the point where $U_{\text{pl}}(T,B)$ becomes lower than $U(J,T,B)$ and comparable to kT. The barriers $U_{\text{pl}}(T,B)$ are expected not to be excessively high in the layered superconducting oxides, as a consequence of the high anisotropy and large κ values of these materials, which soften the FLL.[11] The plastic TAFF process would, hence, be a real possibility in HTSC. In this case an exponential decay of the magnetization, typical of diffusive phenomena, would be observed. According to this picture, in a magnetic-relaxation experiment which was carried out at long enough times, one should expect three different successive regimes as J/J_c decreases: first, when $J \approx J_c$, the single vortex (bundle) creep dominates the relaxation process, leading to a $m(t) \alpha \ln(t)$ decay (Anderson-Kim regime); second, collective creep would take place when decreasing J, giving a decay governed by Eq. (7); and finally,

when J/J_c was small enough, TAFF of defects would play a major role in the relaxation process, leading to an exponential decay of the magnetization before its equilibrium value was reached. The nonlogarithmic relaxations observed in our data would, hence, indicate that the Anderson-Kim regime has already been surpassed and that expressions (6) or (8), which take into account the nonlinearity of $U(J)$, might be more suitable to describe the relaxation process.

An accurate analysis of the nonlogarithmic relaxation data has been performed in view of the above two interpretations (AK or collective flux-creep and vortex-glass models). First, we have checked to what extent experimental data may be described by Eqs. (7) and (9). What we found is that these expressions are not able to fit the experimental data at long times. Nevertheless, a good description of the data for long-time relaxation, $t > t_{cr}$, and high fields $\mu_0 H \geq 1.4$ T can be obtained by using an exponential decay. Figure 4 shows the data taken at $\mu_0 H = 1.5$ T. The solid line is a fit of the points $t < 1200$ s to the logarithmic law (1) and the dashed line corresponds to a description of the data at $t > 1600$ s by using

$$m_c(t,T,B) - m_{eq}(T,B) \alpha \exp(-t/\tau_2) , \qquad (10)$$

with $\tau_2 \approx 1500$ s; $m_{eq}(T,B) \approx 43$ G is the value of the equilibrium magnetization determined from the hysteresis loops. A clear crossover from the logarithmic to the exponential regime is noted at $t_{cr} \approx \tau_2$. This crossover shifts to shorter times as the field is increased; at $\mu_0 H = 1.7$ T (also shown in Fig. 4) a value $t_{cr} \approx 800$ s is found. An exponential decay of magnetization is observed in the whole time window as H is further increased, until any trace of relaxation disappears at fields close enough to $\mu_0 H_{irr}(T=4.5\ \text{K}) \approx 2$ T, at which the equilibrium magnetization has been obtained [see, for instance, the curve corresponding to $\mu_0 H = 1.9$ T, in Fig. 1(b)]. The same general trend with an even sharper transition between both regimes has been observed for higher temperatures, at lower fields and time ranges.

Consequently, the present relaxation data can be described by logarithmic or exponential laws, but no traces of any other regimes (as those expected in collective flux-creep or vortex-glass models) have been observed. Furthermore, the evaluation of the crossover time τ_2 predicted in the AK model [Eq. (4)] by using the value of $U_c(T,B)$ at $B = 1.5$ T, obtained from the fit of the experimental data to Eq. (1), with $\tau \approx 10^{-9}$ s, gives $\tau_2 \approx 10^3 - 10^4$ s, in good accordance with the crossover time observed in the relaxation data. This good agreement gives support to the interpretation of our data as a crossover from single flux creep to diffusive flux motion, within the framework of the Anderson-Kim assumption of a linear $U(J)$. Besides, this estimate gives further support to our earlier assumption of a characteristic time $\tau \approx 10^{-9}$ s.

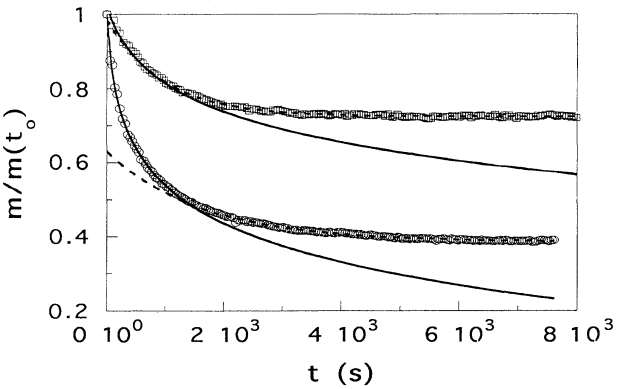

FIG. 4. Magnetization decay at $T=4.5$ K, for $\mu_0 H = 1.5$ T (○) and 1.7 T (□). The solid (dashed) lines are fits to a logarithmic (exponential) law at short (long) times. A change of regime is clearly seen at $t_{cr} \approx 1500$ s for $\mu_0 H = 1.5$ T and $t_{cr} \approx 800$s for $\mu_0 H = 1.7$ T.

The above interpretation, if true, would imply that interactions between vortices could be ignored and thus, collective creep would not take place. Nevertheless, the elastic interactions in the FLL cannot be usually ignored at small driving currents. They should lead to a nonlinear $U(J)$, not compatible with the above interpretation. It would thus be interesting and desirable to get an evaluation of the actual $U(J)$ dependence in our crystal. An estimate can be obtained by means of the analysis of Maley *et al.*[7] This method consists in plotting $U(J,T,B) \equiv kT[C(T,B) - \ln|\partial m/\partial t|]$ vs $m - m_{eq} \alpha J$, at different temperatures, for a given field. The constant C is chosen to get a unique $U(J)$ curve. The actual $U(J)$ function can only be obtained when the temperature dependence of U_c and m_c are accounted for. This can be achieved by choosing C in order to scale $U(J,T,B)/(1-T/T_c)^n$, instead of $U(J,T,B)$.[26] The factor $(1-T/T_c)^n$ here stands for the global dependence on temperature of the activation barrier, coming from both the temperature dependences of the critical current $J_c(T,B) \alpha m_c(T,B)$ and the effective barrier energy $U_c(T,B)$ at $J \approx J_c(T,B)$. Introducing $U_c(T) \sim (1-T/T_c)^\nu$ and $J_c(T) \sim m_c(T) \sim (1-T/T_c)^\eta$ into the effective energy $U(J)$ given by Eq. (6), one has $n = \nu + \mu\eta$. C is assumed to have no dependence on temperature, since this would be very smooth, of the type $C(T) = C_0 + n' \ln(1-T/T_c) \approx C_0$, with n' close to unity.

In the case of electron-doped superconductors, as a consequence of the low values of H_{c2} and T_c,[17] the analysis of Maley is restricted to a very narrow window of fields and temperatures. In the present study, using temperatures between 4.5 and 9 K, the range of accessible fields is $0.2\ \text{T} \leq \mu_0 H < 0.5$ T. The lower limit is the full penetration field at $T=4.5$ K, $H_p(4.5\ \text{K})$, and the upper limit corresponds to the highest field at which exponential relaxations are not yet observed at $T=9$ K.[27] In Fig. 5 we show the results we have obtained by using the analysis of Maley described above. The parameters n and C have been chosen in order to get the best collapse of the experimental data recorded at different temperatures. We have obtained $n = \frac{1}{2}$ and $C = 17 \pm 2$ for all the fields.

The constant $C(T,H)$ can be written in terms of macroscopic quantities as follows:

$$C(T,B) = \ln\left[\frac{m_c(T,B)}{\tau}\frac{kT}{U_c(T,B)}\right] \qquad (11)$$

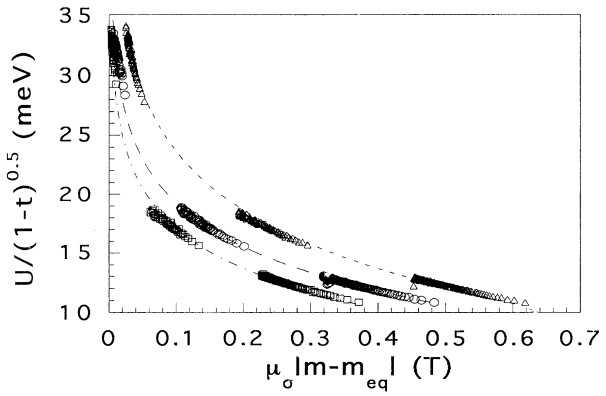

FIG. 5. Plot of Maley for $\mu_0 H=0.2$ (△), 0.3 (○), and 0.4 (□) T. The activation energy is defined by $U(J,T,B) \equiv kT(C-\ln|\partial m/\partial t|)$; a value of $C=17$ has been used for all fields. The lines correspond to logarithmic fits.

for an infinite slab, when $U_c \gg kT$.[28] $U_c(T,B)$ is the linear extrapolation to $J=0$ of the barrier height at $J \approx J_c$ and τ is the macroscopic characteristic relaxation time introduced earlier. The experimental value is in excellent agreement with that obtained from Eq. (11) by using our estimate $\tau=10^{-9}$ s, that is, $C=18$. The incorporation of the field dependences of m_c and U_c (see above) into C gives a field correction of $\Delta C(B)=\ln[\exp(-B/B_0)B^{0.15}]$, which implies $\Delta C(\mu_0 B=0.2\ \mathrm{T})=-0.06$ and $\Delta C(\mu_0 B=0.5\ \mathrm{T})=-2$. These values fall within the error made when determining C, and thus justify the use of a field-independent C.

Regarding the dependence of the activation energy on the driving current J, it can be described by a logarithmic law [Eq. (8)], with $U(B) \approx 11-16$ meV, in agreement with the $U_c(T,B)$ values determined from the logarithmic relaxations. Logarithmic dependences have been reported from relaxation measurements[6] and I-V data[7] in hole-doped HTSC. A potential $U(J)$ dependence has also been observed in $YBa_2Cu_3O_{7-\delta}$, $\delta \leq 0.2$.[29] Recent I-V and magnetoresistance measurements on $Sm_{1.85}Ce_{0.15}CuO_{4-y}$ single crystal have also provided evidence of a logarithmic $U(J)$ dependence.[30] As mentioned above, some theoretical studies[8] have shown that in the case of small exponents μ of the power-law dependence expected in the framework of the collective-creep theory [Eq. (7)], logarithmic approximations may account in a suitable way for the experimental data; actually, our data can also be reasonably described by $U(J,B)=U(B)[(J/J_c)^{-\mu}-1]$ and provide small values of μ, ranging from 0.14 to 0.25, as expected.

DISCUSSION

Our relaxation data show two main features that cannot be easily reconciled: (i) the relevant activation energies for thermally activated flux motion display a nonlinear dependence on the driving current J; (ii) at each temperature, a crossover from logarithmic to exponential decay of the magnetization is observed as the magnetic field is increased. No signatures of other creep decays, associated to the nonlinearity of $U(J)$, have been observed.

The observed nonlinearity of $U(J)$ in our crystal rules out the simplest explanation of the experimental crossover from logarithmic ($J \approx J_c$) to exponential relaxation ($J \to 0$) in terms of the AK model. However, the explanation in the framework of the collective-creep–vortex-glass model has the difficulty of the absence of the collective-creep decay at long times [Eq. (7)] which should become observable between the logarithmic and the exponential decays. In fact, in $YBa_2Cu_3O_{7-\delta}$, $\delta \leq 0.2$, the interpolation formula (7) has provided an adequate description of the relaxation data. An explanation for our apparently contradictory experimental observations might be found by taking into consideration the role played by the characteristic barriers for plastic motion, $U_{\mathrm{pl}}(T,B)$. Let us assume that they are low, of the same order of magnitude as that of $U_c(T,B)$. In this case, the plastic TAFF regime [i.e., the exponential decay of $m(t)$] might take place immediately after the AK logarithmic decay, without the observance of any collective-creep regime. The barrier for plastic motion has been estimated to be $U_{\mathrm{pl}} \approx \Phi_0^2 a_0/(\Gamma^{1/2}\mu_0 8\pi^2\lambda^2)$ (Ref. 11), with $\Gamma=m_c/m_{ab}$ the anisotropy ratio, and λ the in-plane penetration depth. Using $\lambda_{ab}(0)=800$ Å,[17] a value of $U_{\mathrm{pl}}(B) \approx 1.4(\Gamma B)^{-0.5}$ eV is found. Even though many different values have been reported for Γ in these materials, it seems likely that $\Gamma > 1000$;[31] hence, an energy $U_{\mathrm{pl}}(T,B) < 45(1-T/T_c)/B^{-0.5}$ meV is obtained. This is only slightly higher than the experimentally determined values, $U_c(T,B)$. Therefore, not only plastic motion can be possible, but it can also be responsible for the observed crossover. At this point, it is worth mentioning that activation energies of $U(T,B) \approx 40(1-T/T_c)B^{-0.6}$ meV have been reported[13,14] from ac-susceptibility measurements with very low ac-excitation fields; these values are very close to those estimated for $U_{\mathrm{pl}}(T,B)$. Actually, the possibility of plastic motion in the flux-line lattice of these materials has been already considered in Ref. 14.

If a change from the creep of an elastically correlated flux volume to plastic flux motion is indeed responsible for the experimental crossover in magnetization decay, the observed logarithmic relaxations would be a consequence of the fact that the time interval in which $U(J,T,B) < U_{\mathrm{pl}}(T,B)$ is too short for the nonlinearity of $U(J)$ to be observed. In this case, the activation energies and critical currents extracted from the logarithmic fits may not reflect the actual $U_c(T,B)$ and $J_c(T,B)$ coming from collective pinning in an elastically correlated volume V_c. This is clearly appreciated in Fig. 6, where a nonlinear $U(J)$ is depicted. Though in both A and B regions the J dependence of U can be well approximated by a linear function, it can be seen that the determinations of U_c and m_c using region B represent, respectively, an overestimate and an underestimate of the actual values. On the other hand, region A, which lies closer to $J/J_c \approx 1$, would provide much better estimates of the actual U_c and m_c. As a consequence, the experimentally determined energies and critical currents coming from the logarithmic fits and their field dependences have to be

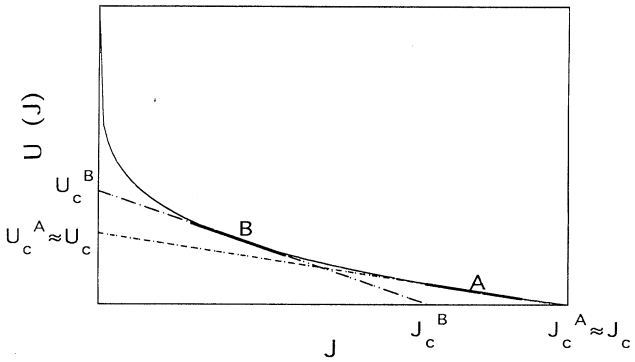

FIG. 6. Nonlinear effective energy $U(J)$. It can be shown that the dependence on J in regions A and B can be well approximated linearly, so that relaxation would be, in both cases, logarithmic; however, the extrapolated values of U_c and J_c are, respectively, larger and smaller than the actual ones; this effect becomes more important as the considered J/J_c values decrease.

considered as approximations to the actual ones; since the relaxation process is expected to become more important for increasing fields and temperatures, the field dependence of U_c (J_c) determined from the logarithmic fits would be smoother (stronger) than the actual one. Nevertheless, as Figs. 3 and 4 are obtained for a low temperature, $T=4.5$ K, we expect this effect not to disturb appreciably our evaluations of $U_c(B)$ and $m_c(B)$ at this temperature.

The above considerations allow one to get a deeper insight into the earlier mentioned discrepancy between the field dependence of the potential barrier obtained from ac susceptibility and the one coming from relaxation measurements. The activation energy extracted from logarithmic relaxation data is related to the creep of single bundles of elastically correlated flux lines. The energies extracted from ac-susceptibility and resistivity measurements, which are usually carried out at low driving currents, may be either related to the same kind of creep, though at different J/J_c values, or to plastic motion, depending on the exact value of $J/J_c(T,B)$ at which they are determined. Both energies are thus expected to display different values and field dependences, as observed. Nevertheless, it comes out from the previous discussion that their orders of magnitude should lie close, the energies coming from ac susceptibility and resistivity being slightly larger; this fact has also been reported.[14,16]

CONCLUSIONS

The magnetic relaxation in a single crystal of the electron-doped superconductor $Pr_{1.85}Ce_{0.15}CuO_{4-y}$ has been studied as a function of field and temperature. Two clearly different regimes have been found in the relaxation process, with a sharp crossover between them. The usual logarithmic behavior is observed at low fields but changes to exponential at high fields, and eventually any trace of relaxation disappears at even higher fields, at which the equilibrium magnetization is reached. The change of regime becomes sharper with increasing temperature, and even at the lowest temperature we have not found any trace of the intermediate behavior in which the interpolation formula for the relaxation process could apply. Such a sharp transition from the Anderson-Kim regime to a diffusive one may be due to the conjunction of the rather large values of ξ and Γ in these materials that may enhance, respectively, pinning and bending of flux lines, leading to rather low values of the ratio $U_{pl}(T,B)/U_c(T,B)$; a crossover from creep of individual (bundles of) correlated vortices to plastic flux motion would thus become more likely. In this sense, it would be interesting to perform relaxation measurements in wider time intervals and at lower temperatures ($T<4.5$ K), where the crossover should be smoother, in order to check whether the nonlinear $U(J)$ dependence in the creep regime becomes evident. This would be manifested in the appearance, between the logarithmic and exponential decays, of the relaxation law predicted in the collective-creep–vortex-glass regime.

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