Magnetic relaxation in an isotropic extreme type-II superconductor

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(Received 21 May 1990)

The time dependence of the zero-field-cooled and thermoremanent magnetization of the isotropic extreme type-II superconductor $PbMo_6S_8$ has been measured at temperatures $0.7 \leq T/T_c \leq 0.95$ and for fields $1 \leq H \leq 300$ G. The results show that the time dependence of the decay can be divided into three different field regimes: (1) A low-field regime with a closely logarithmic decay and a well-established critical state. The relaxation rate $[\partial M/\partial \ln(t)]$ increases as H^5 for the zero-field-cooled magnetization and as $H^{2.4}$ for the thermoremanent magnetization. The reduction in magnetization over 5 decades in time ($0.3-3\times10^4$ s) is, at most, a few percent. (2) A crossover-field regime with a change from a logarithmic to a nonlogarithmic decay. (3) A high-field regime where the time dependence of the decay is well described by the expression $\delta M \sim [\ln(t/\tau_0)]^{-\alpha}$. The reduction in magnetization in the crossover- and high-field regimes is as large as 25–30 % in the measured time interval. The results in the high-field regime are interpreted within the theory of collective flux creep. Comparisons are made with the magnetic relaxation found in anisotropic high-T_c superconductors.

INTRODUCTION

Soon after the discovery of high-T_c superconductors,[1] it was observed that sintered, as well as single-crystalline, materials showed large relaxation effects in their magnetization. The observation of an irreversibility line,[2] which showed some resemblance with the Almeida-Thouless line in spin glasses, together with the memory effects observed in measurements of the magnetic relaxation by Rossel *et al.*[3] have led to the suggestion of a glassy behavior in high-T_c superconductors. Hetzel and Morgenstern[4] have, by using a disordered and frustrated XY model to model the behavior of weakly coupled superconducting grains, managed to reproduce several of these experimental findings. Others have argued that the observed time decay of the magnetization, generally referred to as being logarithmic, can be understood in terms of a conventional flux-creep model with thermally activated flux motion.[5] The unusually large magnetic relaxation of high-T_c superconductors is in such a model explained by small values for the activation energies, a direct consequence of the small coherence length ξ and high measurement temperatures. Fisher[6] has argued that, for an extreme type-II superconductor (Ginzburg-Landau parameter $\kappa=\lambda/\xi \gg 1$), pinning of vortices, by e.g., impurities, destroys the long-range order of the flux lattice. At sufficiently low temperatures, a new thermodynamic phase is predicted, a vortex-glass superconductor, while at higher temperatures thermal fluctuations of the vortex-line density will result in an unpinned vortex liquid. Fisher predicts that the magnetization (and the metastable currents) in the vortex-glass state will decay as $\delta M \sim [\ln(t)]^{-\alpha}$, where $\alpha > 1$. A similar time dependence is proposed by Feigel'man *et al.*[7] within their theory of collective flux creep developed in the case where the screening current density (j) is much smaller than the critical current density $j \ll j_c$. Nelson and Seung[8] have also presented theoretical arguments supporting that the vortex lattice is melted over a significant portion of the H-T plane. The large anisotropy of high-T_c superconductors, resulting in a much smaller value for the elastic tilt modulus of the vortices as compared to the case of an isotropic superconductor, will, according to these authors, result in a heavily entangled flux liquid. At present, there is no complete theory predicting what effect an entangled flux liquid will have on the vortex dynamics.

To get a better understanding of the magnetic relaxation effects of extreme type-II superconductors, it is of paramount importance to experimentally examine the relaxation effects in a wide field and temperature regime. Furthermore, to single out the effect of a small coherence length on the time decay of the magnetization in high-T_c superconductors, it is valuable to investigate the magnetic relaxation in an isotropic, extreme type-II superconductor.

In this paper we present results of the zero-field-cooled and thermoremanent magnetic relaxation of the isotropic, extreme type-II, superconductor $PbMo_6S_8$. Our results show that the relaxation behavior in our experimental time window ($0.3-3\times10^4$ s) can be divided into three different field regimes.

(i) A low-field regime, $H < H^*(T)$, with a closely logarithmic magnetic relaxation [$H^*(T)$ is the field at which full flux penetration in the sample is reached]. The observed decay increases rapidly with the applied magnetic field. The reduction of the magnetization in the experimental time window is, at most, a few percent. The decay of the thermoremanent magnetization is always larger than the decay of the zero-field-cooled magnetization.

(ii) A crossover-field regime, $H \sim H^*(T)$, with a gradual change from a logarithmic to a nonlogarithmic magnetic decay. The relaxation rate of the zero-field-cooled magnetization ($\partial M/\partial \ln t$) exhibits a maximum at a certain field-dependent observation time. It will be argued below that full flux penetration in the sample is achieved at the observation time where the relaxation rate curve exhibits a maximum. The important implication of this result is that full flux penetration in the sample can be achieved both with increasing magnetic field and with increasing *observation time*. The observed decay can be as large as 25–30 % of the magnetization value detected at the instant at which the recording of the magnetic relaxation was started.

(iii) A high-field regime, $H > H^*(T)$, with a nonlogarithmic magnetic decay. The time dependence of the relaxation is well described by an expression of the type $\delta M \sim [\ln(t/\tau_0)]^{-\alpha}$. The decay of the thermoremanent magnetization is independent of field while the decay of the zero-field-cooled magnetization decreases with increasing magnetic field. The reduction of the magnetization is of the same order as in the crossover-field regime.

The results of the magnetic relaxation measurements on the $PbMo_6S_8$ compound will be compared with results obtained from magnetic relaxation experiments on high-T_c superconductors.

EXPERIMENT

A single crystal of the superconductor $PbMo_6S_8$ with a hexagonal-rhombohedral crystal structure, was used in the present measurements. The single crystal has the size $1\times0.8\times0.3$ mm^3 and the weight $m \approx 0.0015$ g. A coherence length of $\xi \sim 25-30$ Å, giving $\kappa \approx 140-150$, has previously been determined for this compound from measurements of H_{c2} at low temperatures. Furthermore, an anisotropy ratio of only 1.2 has been reported for H_{c2} measured perpendicular and parallel to the ternary axis.[9] This anisotropy is negligible when compared to the corresponding anisotropy ratios determined for the $YBa_2Cu_3O_7$ and $Bi_{2.2}Sr_{1.7}CaCu_2O_8$ high-T_c superconductors.[10] Rossel *et al.*[11] have previously shown, using single crystals of $PbMo_6S_8$ from the same batch as the crystal used in the present study, that the $PbMo_6S_8$ superconductor exhibits similar irreversibility and magnetic relaxation properties as those generally attributed to high-T_c superconductors.

A dc SQUID magnetometer was used to study the time decay of the zero-field-cooled (ZFC) and the thermoremanent (TRM) magnetization in the observation time window $0.3-3\times10^4$ s. The magnetic field, oriented randomly with respect to the lattice, is produced by a superconducting NbTi magnet operating in persistent mode. The switch time for the magnet is shorter than 0.01 s. The ZFC (TRM) decay is measured by cooling the sample in zero field (an applied field) to the measurement temperature. The magnetic field is then applied (removed) and the relaxation of the magnetization is measured as a function of observation time. Having completed a measurement of the time decay of the ZFC (TRM) magnetization, the sample is heated to a temperature above T_c (≈ 14.5 K) and the absolute value of magnetization at the measurement temperature is established. This procedure for measuring the magnetic relaxation was followed for several temperatures in the range 10 ($T/T_c \approx 0.7$)–13.8 K ($T/T_c \approx 0.95$). At each temperature the relaxation was measured at more than 20 different fields in the range 1–300 G.

RESULTS AND DISCUSSIONS

Figure 1 shows the ZFC and TRM magnetization as a function of field at the temperatures 13.0 and 13.5 K. The different curves have been obtained from measurements of the time decay of the ZFC and TRM magnetization and each curve illustrates the field dependence of the magnetization at a specific observation time in the relaxation experiment. The arrows indicate values of H_{c1} obtained from the expression

$$H_{c1}(T) \approx 1.73 H_{c1}(0)(1 - T/T_c)$$

using $H_{c1}(0) = 59$ G.[9,12] Figure 2 shows the relaxation rate $\partial M/\partial \ln(t)$ for the ZFC and TRM magnetization versus field at the same temperatures as in Fig. 1. The arrows indicate the field strengths where a change from logarithmic to nonlogarithmic relaxation of the magnetization occurs. The relaxation rates plotted in the field regime with a nonlogarithmic decay are deduced at an observation time of ≈ 30 s. While the decay of the ZFC magnetization starts approximately at $H_{c1}(T)$ (defining the onset of the decay as the lowest field at which the relaxation is experimentally detectable), a decay of the remanent magnetization is detectable even at fields below $H_{c1}(T)$, since, in this case, flux is trapped in the superconductor. In the field regime with a logarithmic decay of the magnetization, the relaxation rate increases as $\partial M/\partial \ln(t) \sim H^5$ for the ZFC magnetization while it increases as $\partial M/\partial \ln(t) \sim H^{2.4}$ for the TRM magnetization. In the field regime with a nonlogarithmic decay, the relaxation rate of the TRM magnetization becomes field independent, while for the ZFC magnetization it goes through a maximum and then decreases with increasing field. The magnitude of the decay in the field regime with a logarithmic decay of the magnetization is, at most, a few percent of the magnetization value while the magnetization decays as much as 30% in the field regime with a nonlogarithmic decay. From Figs. 1 and 2 it is seen that the time decay of the TRM magnetization is larger than (or equal to) the time decay of the ZFC magnetization for all fields. One possible explanation for a smaller decay of the ZFC magnetization at low fields is that there exists a surface barrier to flux motion into the sample. With an

applied field, there are two separate forces which flux lines feel close to the surface of the sample.[13,14] One force, the image force, is imposed by the boundary condition of zero current normal to the surface. This condition is satisfied by placing an image flux line, of opposite sign, outside the surface and the force will thus be attractive to the surface and the flux line energy due to this force increases with increasing distance from the surface of the sample. The other force is created by the external field which will be parallel to the flux lines in the sample. This force is repulsive from the surface and contributes to the total energy of the flux lines in an opposite sense as compared to the image force. Thus, in the ZFC magnetization, there can exist a surface barrier to flux motion into the sample whereas no surface barrier will exist for flux motion out of the sample for the TRM magnetization since such measurements are performed in zero field. It should be noted, however, that the surface barrier for the ZFC magnetization should only exist up to some maximum field above which the repulsive field force dominates on all length scales.[13]

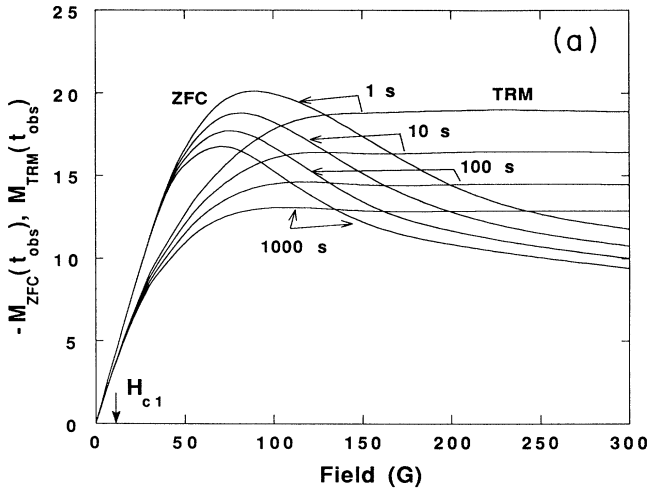

FIG. 1. ZFC ($-M_{ZFC}$) and TRM (M_{TRM}) magnetization vs field. The scale on the y axis corresponds to the measured SQUID voltage. The curves have been obtained from measurements of the time decay of the ZFC and TRM magnetization and each curve illustrates the field dependence of the magnetization at a specific observation time of the relaxation experiment. The arrows indicate calculated values of H_{c1}. $T_c = 14.5$ K. (a) $T = 13.0$ K, $H_{c1} \approx 10$ G. (b) $T = 13.5$ K, $H_{c1} \approx 7$ G.

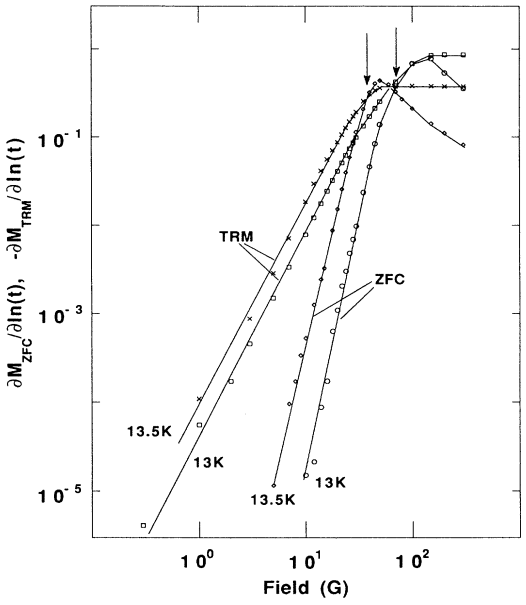

FIG. 2. ZFC and TRM relaxation rate $[\partial M/\partial \ln(t)]$ vs field in a log-log diagram at two different temperatures. The rate is given in the same units as in Fig. 1. The background relaxation rate $\partial M/\partial \ln(t) \approx 1 \times 10^{-5} H$ has been subtracted from the plotted data. The arrows indicate the field strengths where a change from logarithmic to nonlogarithmic relaxation occurs. The relaxation rates plotted in the field regime with a nonlogarithmic decay correspond to $\partial M/\partial \ln(t)$ at an observation time of ≈ 30 s.

To study the effect of temperature on the measured relaxation, the ZFC magnetization was measured in a wider temperature interval $0.7 \leq T/T_c \leq 0.93$. The results are presented in Fig. 3 where $\partial M/\partial \ln(t)$ is plotted versus H in a log-log diagram. Only the low-field regime with a logarithmic decay of the ZFC magnetization is shown. The most significant temperature-dependent features of the measured relaxation are that the onset of relaxation is shifted to higher fields and that the maximum field where a logarithmic decay is found increases when the temperature decreases. Since the decay is logarithmic and since the magnitude of the decay is small in comparison to the magnetization value (the reduction in magnetization is, at most, a few percent measured over 4 decades in time), it is reasonable to assume that a critical state as defined by Anderson[15] is established in the sample. In the critical state model, the flux lines creep along the flux density gradient in the sample by thermally activated hopping between pinning centers. Measurements of the decay of

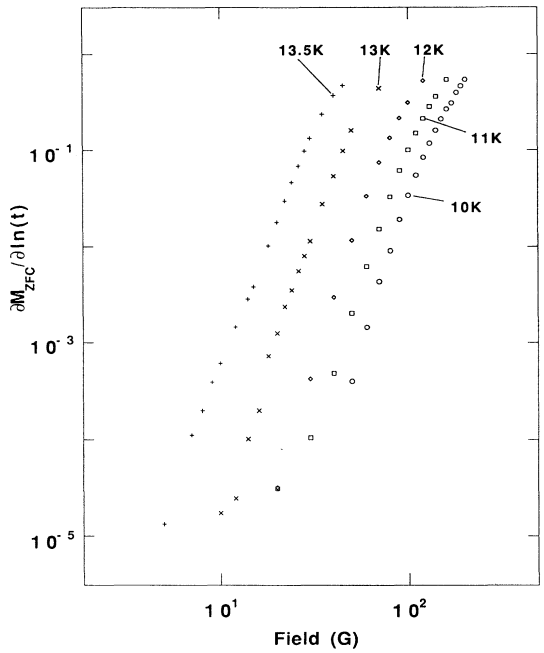

FIG. 3. ZFC relaxation rate $[\partial M_{\mathrm{ZFC}}/\partial \ln(t)]$ vs field in a log-log diagram at five different temperatures in the interval $0.7 < T/T_c < 0.93$. The rate is given in the same units as in Fig. 1. Only the low-field regime with a logarithmic decay is shown. The background relaxation rate $\partial M/\partial \ln(t) \approx 1 \times 10^{-5} H$ has been subtracted from the plotted data.

the magnetization are often used to obtain information on the apparent pinning energies.[5,11] Generally such measurements are performed in the field regime $H \gg H^*$ (where H^* is the field at which the sample is fully penetrated by the field). The results presented in Fig. 3 are, however, from measurements in the field regime $H < H^*$. Yeshurun *et al.*[5] have extended the Bean critical-state model by introducing a field-dependent critical current density $j_c = j_{c1}(H_{c1}/h)^n$, where j_{c1} is the maximum critical current density at a given temperature, h is the local field, and n is associated with the pinning strength. Using this model, the following expression can be derived for the apparent pinning potential:

$$\frac{U_0}{kT} = \ln\left[\frac{t}{\tau_0}\right] + \left[\left[\frac{\partial M}{\partial \ln(t)}\right] \Big/ (M+H)\right]^{-1}, \quad H_{c1} \le H \le H^*, \qquad (1)$$

where τ_0 is a microscopic time usually assumed to be in the range $10^{-6} < \tau_0 < 10^{-10}$ s. This expression differs from the expression valid for $H \gg H^*$ in that $(M+H)$ is used in the denominator instead of just M when calculating the normalized relaxation rate. We have used Eq. (1) to calculate the apparent pinning potential from our data at different temperatures and different fields. At all temperatures, the apparent pinning potential first decreases with increasing field, up to a field ~ 4 times larger than $H_{c1}(T)$, after which it levels off to a constant value. Using $\tau_0 = 10^{-10}$ s, this constant level is ≈ 40 meV for all temperatures shown in Fig. 3. For fields closer to H_{c1}, the calculated values of U_0 are in the range of 100–200 meV (the lower limit corresponding to the higher temperatures shown in Fig. 3). A field dependence of U_0 is not foreseen in the model, the decrease of U_0 with increasing field is therefore rather surprising. It has been suggested[16] that, at sufficiently high fields when the vortex lattice spacing is significantly smaller than the superconducting penetration depth, there will be a crossover to a new kind of collective pinning behavior where the apparent pinning potential depends on field as $U_0 \propto B^{-1}$. The present results have, however, been obtained using fields which do not fulfill this requirement for the vortex lattice spacing and hence such an explanation is not valid in our case. Xu *et al.*[17] have argued that the apparent pinning potential determined from magnetic relaxation experiments will not yield the true depth of pinning potential U_p. In the flux-creep model one customarily assumes that the effective pinning potential is a linear function of the flux density gradient. If this approximation fails, the measured apparent pinning potential will not yield the true value of U_p and results on U_0 will, in this case, not be representative of the true pinning potential and how it depends on temperature and field. The existence of a surface barrier to flux motion into the sample will also have an influence on the extracted values of the apparent pinning potential. It can be shown, using the simple model for the surface barrier discussed above,[13] that the surface barrier decreases with increasing field. With such an explanation, the results presented above imply that flux motion in the low-field limit is partly controlled by the surface barrier while at higher fields, where the extracted values of U_0 levels off, the flux motion is controlled by the intrinsic pinning barriers. The validity of this explanation can be tested by relaxation measurements on a powdered sample of $PbMo_6S_8$, where any influence of surface barriers should be negligible.

Figure 4 shows the time decay of the ZFC magnetization [Fig. 4(a)] and the relaxation rate [Fig. 4(b)] at $T = 13.5$ K. The different curves correspond to different fields in the field regime where there is a gradual crossover with time from a logarithmic to a nonlogarithmic decay. Consulting Fig. 1(b), we see that this crossover behavior occurs at fields slightly larger than the field at which the magnetization obtains its maximum value. The nonlogarithmic character of the decay is particularly evident from Fig. 4(b) where the relaxation rate at constant field exhibits a maximum at a certain field-dependent observation time ($t_{\max}$). With increasing field this maximum is pushed towards shorter observation times in such a way that there is a linear relationship between the applied field and $\ln(t_{\max})$. Since the reduction in magnetization is as large as 25–30 % over 4.5 decades in time and since the relaxation is nonlogarithmic, it would be improper to conclude that the material is in a critical state. Nevertheless, we will use arguments from the critical-state model to interpret the results shown in Fig. 4. Using the extended Bean model[5] (or some other flux-creep model with a local field-dependent critical current[18,19]), it is possible to show that the relaxation rate

of the ZFC magnetization should have a maximum at H^*. It can also be shown, using the same model,[5] that full flux penetration will occur at a field slightly larger than the field (H_m) at which the magnetization obtains its maximum value

$$H^* = [(n+1)H_m^{n+1} + H_{c1}^{n+1}]^{1/(n+1)} .$$

Although not conclusive, remembering the weakness of using the critical-state model to explain the observed behavior, we therefore propose that the observation time where the relaxation rate curve at a given field exhibits a maximum, as shown in Fig. 4(b), corresponds to the time when full flux penetration in the sample is achieved. The important implication of the results shown in Fig. 4(b) is then that full flux penetration can be achieved both with increasing *observation time* (on experimental time scales) and with increasing field. The observation of a linear relationship between the applied field H^* and $\ln(t_{\max})$ lends further support to such an interpretation. It can be shown using the Bean model that H^* will depend on $t_{\max}$ like

$$H^* \propto 1 - (k_B T/U_0)\ln(t_{\max}/\tau_0) ,$$

in accordance with our experimental results.

Although the decay of the TRM magnetization shows a gradual change from a logarithmic to a nonlogarithmic time dependence in the same field regime as the ZFC decay, the relaxation rate curves do not show any maxima as a function of observation time in this case. Instead, with increasing observation time in the crossover-field regime, the TRM relaxation asymptotically approaches the decay of the high-field, field-independent TRM magnetization [see Fig. 5(b)]. This is due to the fact that a partially penetrated flux state is never attained in TRM measurements since, prior to the removal of the applied field, the sample is in a fully penetrated flux state (the Meissner signal is only $\approx 30\%$ of its ideal value at $H = 10$ G).

Figure 5 shows the decay of the ZFC [Fig. 5(a)] and the TRM [Fig. 5(b)] magnetization at 13.5 K in the high-field

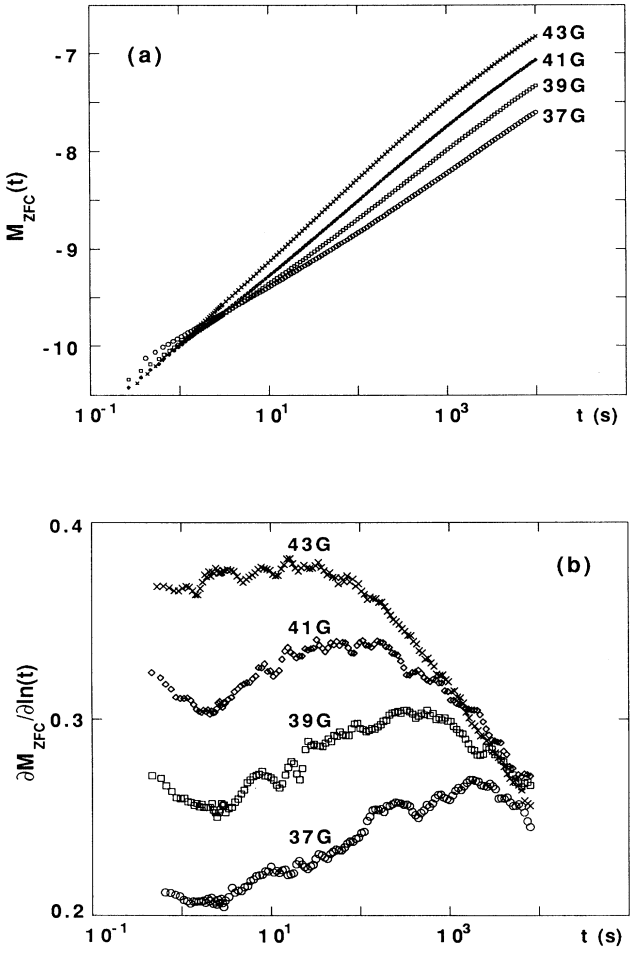

FIG. 4. Time decay of the ZFC magnetization at different fields and the corresponding relaxation rates vs $\log(t)$ at $T = 13.5$ K. The time decay is given in the same units as in Fig. 1. (a) $M_{ZFC}(t)$. (b) $\partial M_{ZFC}/\partial \ln(t)$.

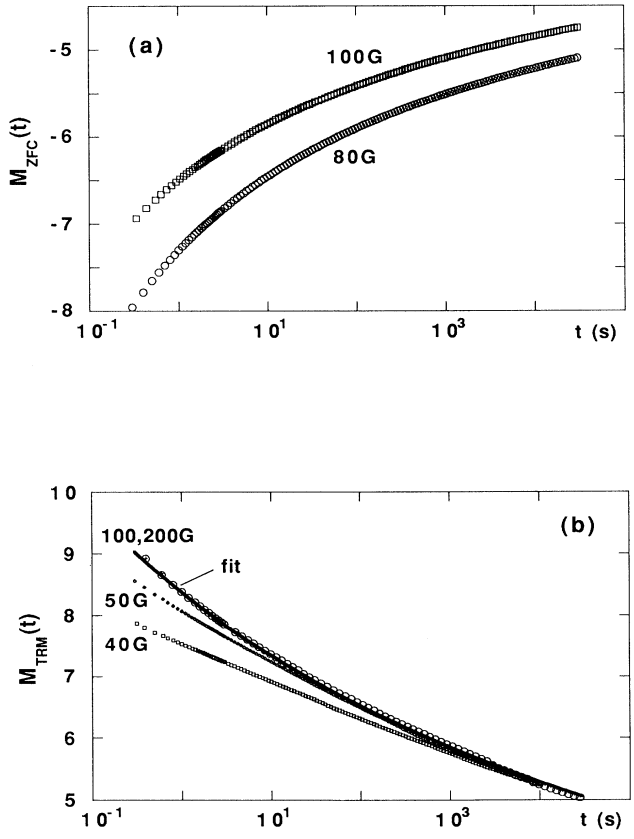

FIG. 5. Time decay of the ZFC (M_{ZFC}) and the TRM (M_{TRM}) magnetization vs $\log(t)$. The time decay is given in the same units as in Fig. 1. The different curves correspond to the decay for different fields in the field regime with a nonlogarithmic decay. (a) $M_{ZFC}(t)$. $T = 13.5$ K. (b) $M_{TRM}(t)$. $T = 13.5$ K. The time decay of $M_{TRM}(t)$ is field independent for $H > 60$ G. The solid line represents the best fit of the experimental data for $H > 60$ G to the expression $M_{TRM} = A[\ln(t/\tau_0)]^{-\alpha}$.

regime. The decay curves of the TRM magnetization for two fields in the intermediate field regime have also been included in Fig. 5(b). The high-field regime is defined from the character of the decay of the magnetization.

(i) The magnetization has a *nonlogarithmic* time dependence.

(ii) The reduction of the magnetization at a given temperature and field is 25–30 %.

(iii) The magnitude of the decay and the absolute value of the TRM magnetization are independent of field while the same quantities for the ZFC magnetization smoothly decrease with increasing field.

As mentioned above, the concept of a critical state is only well defined if the time decay is logarithmic and if the reduction in magnetization is small, i.e., if the screening current density (j) associated with the flux gradient in the sample is of the same order as the critical current density ($j_c - j \ll j_c$. The theory of collective flux creep[7] also includes the case when $j \ll j_c$ and can therefore be employed in a wider field regime. The main result of this theory is that, for $j \ll j_c$, the activation barrier $[U(j)]$ grows with decreasing screening current density as $U(j) \propto j^{-1/\alpha}$. Since $U[j(t)]$ depends on observation time as[20]

$$U[j(t)] \propto T \ln(t/\tau_0) ,$$

it follows that the screening currents will depend on observation time as

$$j(t) \propto [T \ln(t/\tau_0)]^{-\alpha} .$$

In the following it will be assumed that the decay of the screening currents in the limit $j \ll j_c$ is reflected in the time dependence of the magnetization so that

$$\delta M \propto [T \ln(t/\tau_0)]^{-\alpha}$$

[as is the case in the critical-state model in the limit $(j_c - j) \ll j_c$]. The time dependence of both the ZFC and TRM magnetization has been fitted to this expression. In the case of the ZFC decay, one must include the equilibrium magnetization value M_{eq} as a fitting parameter. Although it is possible to obtain good fits, different combinations of fitting parameters yield equally good fits. In the case of the TRM magnetization, the equilibrium magnetization value is zero. It is then possible to construct the following derivative:

$$\{\partial \ln[M_{TRM}(t)]/\partial \ln(t)\}^{-1} ,$$

which should be a linear function of $\ln(t)$, with a slope of $-(1/\alpha)$ and a $\ln(t)=0$ intercept of $(1/\alpha)\log_{10}(\tau_0)$, if the above expression for time dependence of the magnetization is correct. Plotting this derivative, at 13 K for $H > 150$ G and at 13.5 K for $H > 60$ G, as a function of $\ln(t)$ gives straight lines and the following values of the parameters: $\alpha = 1.1 \pm 0.1$ and $\log_{10}(\tau_0) = -8 \pm 1$ at 13 K, and $\alpha = 1.25 \pm 0.1$ and $\log_{10}(\tau_0) = -9 \pm 1$ at 13.5 K. The fit for the decay of the TRM magnetization has been included as a solid line in Fig. 5(b). We have also tried to fit a power law and a stretched exponential to the experimental decay of the magnetization. Neither of these functional forms can give a good description of the observed decay. We believe that collective flux creep with $j \ll j_c$ is progressively reached when full flux penetration of the sample is achieved (at H^*) since, at that point, no large screening currents ($j \sim j_c$) can flow; the critical state breaks down. According to the proposed theory,[7] the exponent α changes for different regimes of collective flux creep. As j decreases below j_c, a first region with $\alpha = 7$ is entered,

$$[j_c > j > j_1 \sim j_c (L_c/a_0)^{7/5}] ,$$

where collective pinning of single vortices occurs. The condition $\xi < L_c < a_0$ is satisfied, L_c being the longitudinal size of the short-range correlation volume V_c of the flux line lattice (flux line lattice with weak disorder) and a_0 the spacing between flux lines. For screening currents

$$j_1 > j > j_2 = j_1 (a_0/\lambda)^2 ,$$

the collective creep is dominated by hopping of flux bundles of size $R_\perp$, $a_0 < R_\perp < \lambda$, where λ is the London penetration depth. In that case, the exponent α is $\frac{2}{3}$. At even lower screening currents $j < j_2$, the condition $R_\perp > \lambda$ is achieved and $\alpha = \frac{9}{7}$. The last exponent is close to the exponent extracted from the experiments $\alpha \sim 1.1$–1.2. This result might suggest that the flux dynamics in $PbMo_6S_8$ at fields higher than H^* is dominated by hopping of large bundles of flux lines ($R_\perp > \lambda$).

A qualitatively similar behavior, with three distinctly different field regimes, has also been observed in relaxation measurements on single crystals of $Bi_{2.2}Sr_{1.7}CaCu_2O_8$.[21] In the low-field regime, the decay is logarithmic and both the ZRC and TRM relaxation rates increase as $\sim H^2$. The crossover-field regime is even more explicit in the $Bi_{2.2}Sr_{1.7}CaCu_2O_8$ superconductor; with increasing observation time the relaxation rate increases by as much as a factor of 5 and obtains pronounced maxima when full flux penetration in the crystal is reached. In the high-field regime and in the same reduced temperature range as in the present measurements ($T/T_c > 0.7$), the time decay of the magnetization for the high-T_c superconductor is well described by the same expression as for the $PbMo_6S_8$ crystal, albeit with an exponent α less than 1 (α is of the order 0.3–0.4 at the highest temperatures measured but increases to 0.6–0.7 at $T/T_c \approx 0.6$). The magnitude of the decay for the $Bi_{2.2}Sr_{1.7}CaCu_2O_8$ crystals is, in this case, of the same size as for the $PbMo_6S_8$ superconductor. At low temperatures, $T/T_c < 0.5$, and in the high-field regime, it is not possible to fit an expression of the type

$$\delta M \propto [T \ln(t/\tau_0)]^{-\alpha}$$

to the observed decay of the ZFC and TRM magnetization. Furthermore, the magnitude of the decay can be as large as 70–80 % of the absolute value of the magnetization at the beginning of the measurement. These results obtained for the anisotropic, extreme type-II supercon-

ductor $Bi_{2.2}Sr_{1.7}CaCu_2O_8$ might indicate a more involved H-T diagram for this superconductor with the possibility of distinctly different thermodynamic vortex phases.[6,8]

CONCLUSIONS

In summary, we have shown that the time decay of the ZFC and TRM magnetization of an isotropic extreme type-II superconductor can be divided into three different field regimes: a low-field regime with a closely logarithmic time decay of the magnetization, a crossover-field regime with a change from logarithmic to nonlogarithmic relaxation, and a high-field regime where the time dependence of the decay is well described by

$$\delta M \propto [T \ln(t/\tau_0)]^{-\alpha} .$$

Our results indicate that conventional flux-creep models are not sufficient to describe the observed behavior of the decay. Instead, the theory of collective flux creep, which also includes the case when the screening current density is much smaller than the critical current density, seems to be applicable, as shown here. The consequences of large thermal fluctuations and quenched disorder on the metastability of the vortex lattice have also been studied by Fisher.[6] Their theory proposes the same nonlogarithmic decay law as observed in these experiments, and suggests the existence of a vortex glass phase in the H-T phase diagram.

ACKNOWLEDGMENTS

Financial support from the Swedish Natural Science Research Council (NFR) is gratefully acknowledged. We thank V. M. Vinokur and V. B. Geshkenbein for useful discussions. Chimie Minérale B, Université de Rennes is Unité associé au CNRS No. 254.

[1]J. G. Bendorz and K. A. Müller, Z. Phys. B **64**, 189 (1986).
[2]K. A. Müller, M. Takashige, and J. G. Bednorz, Phys. Rev. Lett. **58**, 1143 (1987).
[3]C. Rossel, Y. Maeno, and I. Morgenstern, Phys. Rev. Lett. **62**, 681 (1989).
[4]R. Hetzel and I. Morgenstern, *Proceedings of the NATO Advanced Workshop on Relaxation and Related Topics in Complex Systems, Italy, 1989* (Plenum, New York, 1989).
[5]Y. Yeshurun, A. P. Malozemoff, F. Holtzberg, and T. R. Dinger, Phys. Rev. B **38**, 11 828 (1988).
[6]Matthew P. A. Fisher, Phys. Rev. Lett. **62**, 1415 (1989); Daniel S. Fisher, Matthew P. A. Fisher, and David A. Huse (unpublished).
[7]M. V. Feigel'man, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Phys. Rev. Lett. **63**, 2303 (1989); V. B. Geshkenbein, A. I. Larkin, M. V. Feigel'man, and V. M. Vinokur, Physica C **162-164**, 239 (1989).
[8]D. R. Nelson and H. S. Seung, Phys. Rev. B **39**, 9153 (1989).
[9]See, for example, M. Decroux and Ø. Fischer, in *Superconductivity in Ternary Compounds,* Vol. II of *Topics in Current Physics 34,* edited by Ø. Fischer and M. B. Maple (Springer, Berlin, 1982), p. 57.
[10]M. Naughton, R. C. Yu, P. K. Davies, J. E. Fischer, R. V. Chamberlin, Z. Z. Wang, T. W. Jing, N. P. Ong, and P. M. Chaikin, Phys. Rev. B **38**, 9280 (1988).
[11]C. Rossel, E. Sandvold, M. Sergent, R. Chevrel, and M. Potel, Physica C **165**, 233 (1990).
[12]Ø. Fischer, Appl. Phys. **16**, 1 (1978).
[13]C. P. Bean and J. D. Livingston, Phys. Rev. Lett. **12**, 14 (1964).
[14]M. Foldeaki, M. E. McHenry, and R. C. O'Handley, Phys. Rev. B **39**, 2883 (1989).
[15]P. W. Anderson, Phys. Rev. Lett. **9**, 309 (1962).
[16]Y. Yeshurun and A. P. Malozemoff, Phys. Rev. Lett. **60**, 2202 (1988).
[17]Youwen Xu, M. Suenaga, A. R. Moodenbaugh, and D. O. Welch, Phys. Rev. B **40**, 10 882 (1989).
[18]Y. B. Kim, C. F. Hempstead, and A. R. Strnad, Phys. Rev. Lett. **9**, 306 (1962); Phys. Rev. **129**, 528 (1963).
[19]P. Chaddah and G. Ravikumar, Physica C **162-164**, 347 (1989).
[20]V. B. Geshkenbein and A. I. Larkin, Zh. Eksp. Teor. Fiz. **95**, 1108 (1989).
[21]K. Niskanen, P. Norling, P. Svedlindh, P. Nordblad, L. Lundgren, C. Rossel, and G. V. Chandrashekhar (unpublished).