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Nonlinear $U(j)$ dependence determined directly from low-electric-field E-j_s curves in $YBa_2Cu_3O_{7-\delta}$ thin films

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Low-electric-field E-j_s curves, normalized dynamical relaxation rate $Q(T)=d\ln(j_s)/d\ln(E)$, and experimental critical current density $j_s(E=E_c)$ were measured from 10 to 85 K with the magnetic sweeping method. By taking a general assumption for the activation energy $U(j_s,T)=U_c(T)f[j_s(T)/j_c(T)]$, with $t=T/T_c$, we obtained the $U(j,T=0\text{ K})$ relation by scaling all the E-j_s curves of different temperatures to one smooth curve with $U_c(T)$ and $j_c(T)$ determined by the so-called generalized inversion scheme. It is shown that the important parameter $c=\ln(v_0B)$, which was actually taken as a fitting parameter in the scaling process in previous work, can be predetermined. It is also found that the resulting $U(j,T=0\text{ K})$ relation based on the E-j_s curves for $T\leq 79$ K can be well interpreted with the collective-pinning model. However, the data for $T>79$ K start to deviate from this interpretation, which may be attributed to the reverse hopping of vortices.

I. INTRODUCTION

The dissipation in the mixed state of high-T_c superconductors is enhanced strongly by giant flux creep,[1] which has been of interest to many experimentalists and theorists recently in this field. According to the thermally activated flux motion (TAFM) model, the essential problem is to know the peculiar nonlinear $U(j)$ dependence at fixed temperatures and fields, where U is the activation energy for the hopping of flux bundles or flux lines. The linear model proposed early by Kim and Anderson and co-workers[2] was challenged by some authors when they found the nonlogarithmic relaxation of remanent magnetization with time. Therefore, other models, such as the logarithmic model,[3] collective creep,[4] and vortex glass,[5] which predict different $U(j)$ relations, have been proposed to account for this effect. In order to check the validity of these models, it is necessary to obtain the $U(j_s,T)$ relation directly from the experiment, which principally can be achieved by measuring the $E(j_s)$ curves at a fixed temperature, since the TAFM model shows that $U(j_s,T)=k_BT\ln[v_0B/E(j_s,T)]$, where v_0 is the maximum velocity for the flux motion and j_s is the screening current density associated with the electric field E at a certain temperature T. However, in a practical experiment, j_s varies very slowly with E, especially in the low-temperature regime, within the experimentally accessible window for E; e.g., from 10^{-8} to 10 V/m, j_s varies only about one order of magnitude or less. Therefore, it is impossible to use such a short $U(j_s,T)$ segment for a comparison with theory and to determine the general pinning properties of the sample. By measuring the magnetization relaxation $M(t)$ of bulk samples at different temperatures and then scaling them to one curve, Maley *et al.*[6] first proposed a way to determine the $U(j)$ dependence at $T=0$ K in a wide current-region, which was followed by many other contributors.[7–13] In these scaling procedures, $c=\ln(v_0B)$ was treated as a fitting parameter, which assumes quite different values in different papers. In addition, two other unknown functions, $j_c(T)$ and $U_c(T)$, have to be assumed for the scaling, which makes the obtained relation $U(j,T=0\text{ K})$ somewhat arbitrary.

In this paper, we present E-j_s curves measured with the magnetic sweeping method at temperatures ranging from 10 to 85 K. The low-temperature ($T<10$ K) data and analysis will be omitted because of the presence of flux quantum tunneling effects.[14] We will show that the c value can be predetermined, and the unknown functions $j_c(T)$ and $U_c(T)$ can be obtained by treating the data via the so-called generalized inversion-scheme (GIS).[15] Finally, a smooth $U(j,T=0\text{ K})$ curve is obtained in a very wide current regime by scaling all the E-j_s curves of different temperatures onto one master curve. A physical interpretation of the resulting $U(j,T=0\text{ K})$ relation is discussed.

II. EXPERIMENTAL

A highly sensitive torque magnetometer [10^{-10}–10^{-11} N m (Refs. 16 and 17)] was used to measure the magnetic torque acting on the sample. The field was applied 45° to the c-axis of the film. A high-quality $YBa_2Cu_3O_{7-\delta}$ thin film made with the molecular-beam-epitaxy technique was used as the sample. It has a square shape with dimensions of 4 mm×4 mm×150 nm and with $T_{c0}=90.8$ K and $\Delta T_c=0.5$ K. X-ray diffraction showed that only (00l) peaks were observable and the full width at half maximum of the (005) peak was about 0.28°, which indicated that the film had very good crystallinity. An Oxford Instruments cryogenic system providing a field up to 8 T was used for the measurements. The temperature is stabilized to better than ± 0.1 K in the sweeping process. The fields sweeping rate was varied from 0.001 to 40 mT/s, corresponding to electric fields from 10^{-9} to 4×10^{-5} V/m.

III. RESULTS

By sweeping the magnetic field up and down at the same rate around a certain field, we measured the width of the irreversible magnetic moment ΔM. In the field sweeping process, the electric field induced at the perimeter is $E(R)=(R/2)(dB/dt)$, where R is the radius of the specimen. By using the basic assumption of the Bean critical state model,[18] i.e., that the current density is uniform within the sample; $j_s(R)$ can be determined via $j_s(R)=3\Delta M/2\pi R^3 d$, where d is the thickness of the film. Therefore the experimental curves ΔM vs dB/dt can be converted into $E(j_s)$ curves at different temperatures. In Fig. 1, we present the $E(j_s)$ curves obtained in this way at temperatures ranging from 10 to 85 K. In Fig. 2, the experimental critical current density j_s is plotted logarithmically as a function of temperature (shown by open circles). It is clear that in the low-temperature region, $\ln(j_s)$ drops linearly with T, which was attributed to the collective-pinning effect[14,19] and we will use this linear part to determine the c value. The slopes of $\ln(j_s)$ vs $\ln(E)$ at $E_c=4\times10^{-5}$ V/m were also plotted as a function of T in Fig. 2 (filled circles). As has been extensively discussed in Refs. 20 and 21, $d\ln j_s/d\ln E$ determined in a magnetic sweep process is equivalent to the normalized relaxation rate $s=-d\ln M(t)/d\ln t$ measured in a conventional magnetization relaxation experiment.

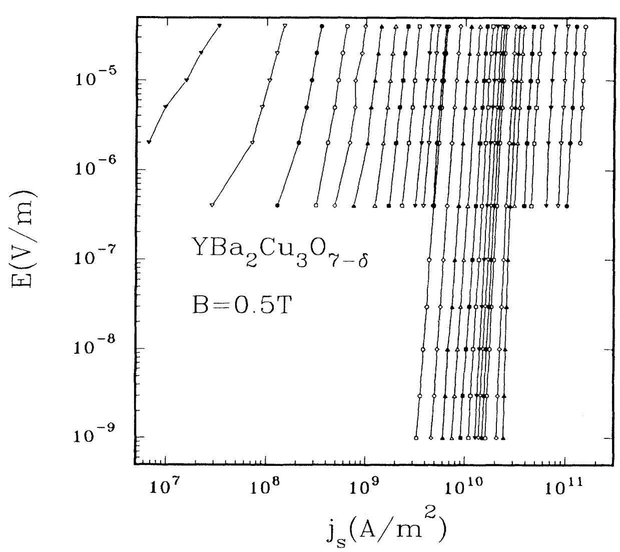

FIG. 1. E-j_s curves at temperatures ranging from 10 to 85 K. The corresponding temperatures from right to left are 10, 15, 20, 25, 32, 40, 44, 49, 51, 53, 55, 57, 58, 60, 62, 64, 68, and 72 K, and from 73 to 85 K with increments of 1 K.

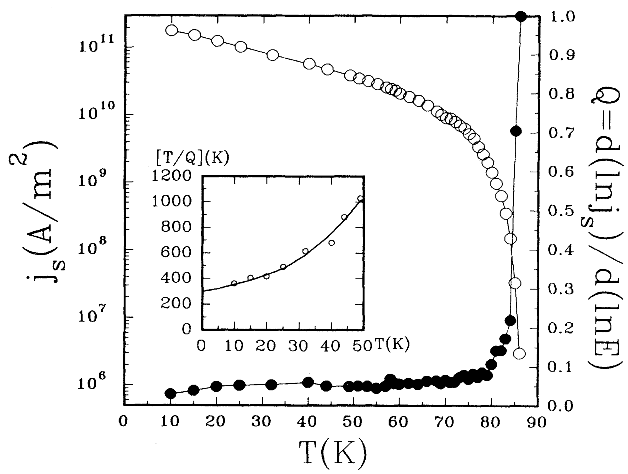

FIG. 2. Temperature dependence of the induced semiconducting current density j_s (open circles) and the dynamical relaxation rate $Q=d\ln j_s/d\ln E$ (filled circles) at the sweeping rate of 40 mT/s (or$E=4\times10^{-5}$ V/m). The inset shows the correlation between T/Q and T at 0.5 T. The value of T/Q in the limit $T\rightarrow 0$ K is about 300 K.

IV. DISCUSSION

A. Determination of $U(j,T=0\text{ K})$ relation from the $E(j_s,T)$ curves

In the presence of a magnetic field B, when a macroscopic current is flowing perpendicular to the field direction, a Lorentz driving force F_l exerted on the vortices will drive them to jump over the pinning barriers with the help of thermal activation. The electric field E induced by the flux motion, can be written as

$$E(j_s,T)=v_0 B\exp[-U(j_s,T)/k_B T]\ , \tag{1}$$

where v_0 is the maximum velocity for flux motion and j_s is the screening current density which is a function of temperature and field sweeping rate. Equation (1) can be rewritten as

$$U(j_s,T)=k_B T[c-\ln E(j_s,T)]\ , \tag{2}$$

with

$$c=\ln(v_0 B)\ . \tag{3}$$

Equation (2) indicates that the E-j_s curve at temperature T can be used to construct a short portion of $U(j_s,T)$. But, unfortunately, the current density j_s varies very slowly with the electric field (or the sweeping rate). In order to determine the $U(j)$ relation in a wide current region, we assume that the activation energy $U(j_s,T)$ can be written as

$$U(j_s,T)=U_c(0)g(T/T_c)f[j_s(T)/j_c(T)]\ . \tag{4}$$

Actually the above assumption is a general case which covers all the models proposed to date.[2–5,22] Combining Eqs. (2) and (4) leads to

$$f[j_s(T)/j_c(T)]=k_B T[c-\ln E(j_s,T)]/[U_c(0)g(T/T_c)]\ . \tag{5}$$

Since $U_c(T)=U_c(0)g(T/T_c)$ is related to the temperature through parameters such as ξ, H_c, and λ, which depend on temperature very weakly near 0 K, therefore at $T=0$ K, it is easy to find that $g(0)=1$ and $U(j,0)=U_c(0)f[j/j_c(0)]$. Thus at $T=0$ K, with a current density $j(T=0\ \mathrm{K})/j_c(0)=j_s(T)/j_c(T)$, we have $f[j(T=0\ \mathrm{K})/j_c(0)]=U(j,0)/U_c(0)=f[j_s(T)/j_c(T)]$; the activation energy $U(j,T=0\ \mathrm{K})$ is

$$U(j,T=0\ \mathrm{K})=[k_BT/g(T/T_c)][c-\ln E(T,j_s)]\ . \tag{6}$$

The above equation shows that, if we have the temperature dependence of j_c and U_c, we can easily determine the $U(j,T=0\ \mathrm{K})$ relation by scaling the E-j_s curves of different temperatures onto one master curve. Principally, c is a function of T through v_0, but as shown by Eq. (3) it will vary logarithmically with v_0, which indicates that c is a very weak temperature-dependent function. In addition, in our measurement, $|\ln E|$ is about 11 to 20, which is much larger than the c value (see below). Therefore, we can safely take c as a constant throughout the whole temperature region.

B. Determination of the c value

During a field sweep at constant rate, the electric field $E=(\mu_0R/2)dH/dt$ is a constant. According to Eq. (2), we have

$$d[\ln U(j_s,T)]/d[\ln T]|_E=1\ . \tag{7}$$

However, from Eq. (4), we have

$$\left.\frac{d\ln U(j_s,T)}{d\ln T}\right|_E=\frac{\partial \ln f}{\partial \ln j_s}\left[\frac{d\ln j_s}{d\ln T}\right]_E+\frac{\partial \ln f}{\partial \ln j_c}\frac{d\ln j_c}{d\ln T}+\frac{d\ln g(T/T_c)}{d\ln T}=\frac{\partial \ln f}{\partial \ln j_s}\left[\frac{d\ln(j_s/j_c)}{d\ln T}\right]_E+\frac{d\ln g(T/T_c)}{d\ln T}\ . \tag{8}$$

When T approaches 0, $d[\ln g(T/T_c)]/d[\ln T]=0$; thus equating Eq. (7) and Eq. (8) gives

$$\frac{\partial \ln f}{\partial \ln j_s}\left[\frac{d\ln(j_s/j_c)}{d\ln T}\right]_E=1\ . \tag{9}$$

On the other hand, from Eqs. (2) and (4), we have

$$\left.\frac{\partial \ln U}{\partial \ln j_s}\right|_T=\left.\frac{\partial \ln f}{\partial \ln j_s}\right|_T=\left.\frac{\partial \ln U}{\partial \ln \dot B}\frac{d\ln \dot B}{d\ln j_s}\right|_T=-\left.\frac{1}{Q(T)}\frac{1}{c-\ln E}\right|_T\ , \tag{10}$$

where $Q=d\ln j_s/d\ln(dB/dt)$. When $T\to 0$, combining Eqs. (9) and (10) leads to

$$c=-\left.\frac{T}{Q(T)}\frac{d\ln(j_s/j_c)}{dT}\right|_{T\to 0}+\ln E_c\ , \tag{11}$$

where E_c is the electric field established in the sample during a sweep at a rate of 40 mT/s. Since j_c is a function of ξ, H_c, and λ, which are weak temperature-dependent functions in the low-temperature region, with $T\to 0$ K, $d\ln j_c/dT\approx 0$. Therefore, we have

$$c=-[(T/Q(T))d\ln j_s/dT]|_{T\to 0}+\ln E_c\ . \tag{12}$$

From Fig. 2 we know that $\ln j_s$ linearly decreases with temperature; therefore, the slope $d\ln j_s/dT|_{T\to 0}$ can be easily determined. In the inset of Fig. 2, the relation of $T/Q(T)$ vs T is presented. Although the curve is not straight when T approaches 0 K, there is a clear finite intercept of T/Q at $T=0$ K which is about 300 K. Using this and $E_c=4\times 10^{-5}$ V/m into Eq. (12), c is found to be 2.45; therefore, the maximum velocity for flux motion is about 20 m/s.

C. Determination of the $j_c(T)$ and $U_c(T)$ relations

As soon as the c value is known, with proper choices for $j_c(T)$ and $U_c(T)$ we should be able to determine the $U(j,T=0\ \mathrm{K})$ relation, by means of Eq. (6), and $j=j_s(T)j_c(0)/j_c(T)$. However, to the best of our knowledge, there is no *a priori* unique way to predefine the temperature dependence of j_c and U_c. As shown by Schnack *et al.*,[15] this problem can be solved by means of GIS. Basic assumptions for the GIS are Eq. (4) and

$$g(T,B)=[j_c(T,B)/j_c(0,B)]^p[G(T)]\ , \tag{13}$$

where p and $G(T)$ depend on the dimensionality of the vortex system and the regime of vortex pinning.[23,24] For example, for a collectively pinned single vortex, the pinning potential $U_c=j_c\phi_0L_cr_p$, where L_c is the collective-pinning length which depends strongly on the quenched disorder parameter j_0/j_c, ϕ_0 is the flux quantum, r_p is the pinning range, typically $r_p=\xi$. For randomly distributed weak point pinning centers, Blatter *et al.*[25] derived the collective-pinning length $L_c\propto\xi(j_0/j_c)^{1/2}$, where j_0 is the depairing current density ($\propto H_c/\lambda$). Putting $\xi\propto(1+t^2/1-t^2)^{1/2}$, $H_c\propto 1-t^2$, and $\lambda\propto(1/1-t^4)^{1/2}$ into the above expression for U_c, one can easily find that for a collectively pinned single vortex, $p=0.5$ and $G(T)=(1-t^2)^{-1/4}(1+t^2)^{5/4}$. From Eqs. (8), (10), and (13) and in combination with p and $G(T)$, we have then

$$j_c(T)=j_c(0)\exp\left[\int_0^T\frac{(c-\ln E_c)Q(T')\{1-[d\ln G(T'/T_c)/d\ln T']\}+[d\ln j_s(T')/d\ln T']}{1+pQ(T')(c-\ln E_c)}\frac{dT'}{T'}\right]\ . \tag{14}$$

The above equation shows that we can determine the temperature dependence of j_c by a series of integrals step by step from low temperature. In Fig. 3, we present the relations of $j_c(T)$ (open circles) and $j_s(T)$ (filled circles). It is clear that with $T \geq 20$ K, $j_c(T) \geq 2j_s(T)$, which indicates that the relaxation due to thermally activated flux creep is already very strong. The $g(T)$ relation can be easily obtained through Eq. (13).

In the above discussion, we took only the case for single vortex pinning into account. This is because the field in our experiment (0.5 T) is much lower than B_{sb}, the minimum field for the vortices to be pinned collectively in the form of small bundles. According to Ref. 25, $B_{sb} = \beta_{sb} B_{c2}(j_c/j_0)$, with $\beta_{sb}=5$, where j_c is the true critical current density corresponding to $U=0$, and j_0 is the depairing current density which is the order of $H_c/\lambda \approx 10^{13}$ A/m^2. In our experiment, j_c is about 10^{11} A/m^2, with $B_{c2} \approx 140$ T; we have $B_{sb} \approx 7$ T. In Fig. 4, we present the $U(j, T=0$ K) relation obtained by using $p=\frac{1}{2}$ and $G(T)=(1+t^2)^{5/4}(1-t^2)^{-1/4}$. It is evident that all 31 E-j_s curves are smoothly scaled onto one master curve, which may provide an evidence for the single vortex collective pinning. Actually, with the choices for p and $G(T)$ corresponding to small bundles or large bundles collective pinning,[24] the scaling looks terrible and the data points scatter over a very wide band.

D. Discussion of the resulting $U(j, T=0$ K) relation

Here we compare the resulting $U(j, T=0$ K) relation with theoretical models. To do so, we need to know the value of $j_c(0)$. Since at $T=0$ K, there is no thermal activation effect on the flux motion, therefore, as determined from Fig. 2, $j_c(0)=j_s(T)|_{T\to 0\,\mathrm{K}}=2.8\times 10^{11}$ A/m^2. We first tried to fit the data to the logarithmic model $U(j,0)=U_c(0)\ln[j_c(0)/j]$. Unfortunately, no good fit could be found in this way by taking whatever value for $U_c(0)$. We then fitted the data to the collective-pinning and creep models, $U(j,T=0) \propto [j_c(0)/j]^\mu$, as first proposed by Feigel'man *et al.*[4] and the modified form, $U(j,T=0) = [U_c(0)/\mu]\{[j_c(0)/j]^\mu - 1\}$, later given by Malozemoff,[26] where $U_c(0)$ was shown[24] to be roughly the value of T/Q at $T=0$ K. From the inset of Fig. 2, one can easily find that $U_c(0) \approx [T/Q]_{T\to 0\,\mathrm{K}} \approx 300$ K. Putting the values for $U_c(0)$ and $j_c(0)$ into the above equation, we will find the best fit by adjusting the μ value. As shown in Fig. 4 by the solid lines, a good fit can really be obtained for the data of $T \leq 79$ K where μ is chosen as 0.79. The $U(j, T=0$ K) relation based on the data at temperatures $T > 79$ K shows, however, disagreement with the above interpretation. It is interesting to note that 79 K corresponds to the threshold for the rapid dropping of j_s and the drastic increase of Q as shown by open and filled circles, respectively, in Fig. 2, and is also equivalent to the so-called glass transition temperature T_g [as defined by the crossover from negative curvature to positive curvature for the ln(j) vs ln(E) curve]. This transition may be attributed to the occurrence of reverse hopping of flux lines,[27] since in the whole treatment done above, we took only the forward hopping into account. As soon as the reverse hopping is concerned, one of our basic assumptions, i.e., the down-hill flux creep [shown by Eq. (1)] breaks down.

In obtaining the final smooth curve $U(j, T=0$ K), we have made three basic assumptions: (1) the TAFM model $U(j_s,T)=k_B T \ln[v_0 B/E(j_s,T)]$; (2) $U(j_s,T) = [U_c(T)] f[j_s(T)/j_c(T)]$; and (3) $U_c(T,B) \propto [j_c(T,B)/j_c(0,B)]^p [G(T)]$. The first assumption means that the

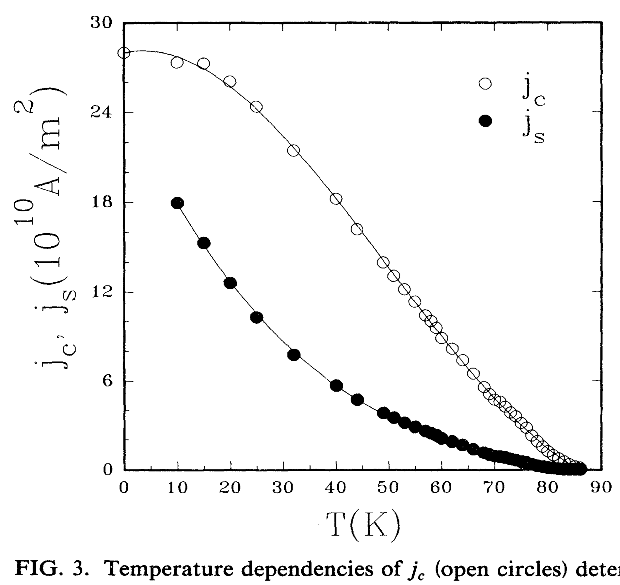

FIG. 3. Temperature dependencies of j_c (open circles) determined by means of GIS and j_s (solid circles). It is clear that $j_c(T)/j_s(T) \geq 2$ when $T \geq 20$ K.

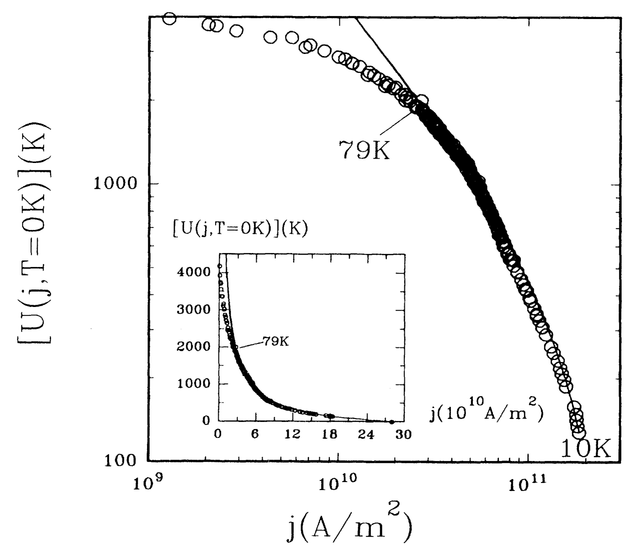

FIG. 4. The symbols show the $U(j, T=0$ K) relation obtained by scaling the E-j_s curves of different temperatures to one master line. The solid lines represent the fits to the collective-pinning expression $U(j,T=0) = [U_c(0)/\mu][(j_c(0)/j)^\mu - 1]$ with $j_c(0)=2.8\times 10^{11}$ A/m^2, $U_c(0)=300$ K, and $\mu=0.79$.

energy dissipation is mainly due to the thermally activated flux motion; the second one indicates that the activation energy can be separated into two parts, one relating to the pinning potential, another one relating to the current density; and the third one describes the relation between $U_c(T)$ and $j_c(T)$; for single vortex collective pinning, it was shown that $p=\frac{1}{2}$ and $G(T)=(1+t^2)^{5/4}(1-t^2)^{-1/4}$. Following these assumptions, we really obtain a smooth $U(j,T=0\ \mathrm{K})$ curve by scaling all the $E(j_s,T)$ curves of different temperatures to one master line. This can be considered as *a posteriori* justification of the validity of the three basic assumptions made in our approach. From the above analysis, although we cannot exclude the possibility of the existence of a vortex glass state at temperature $T\leq 79$ K, the very specific form for $G(T)$ and p values used in our approach was derived only for the collective pinning of a single vortex, which may manifest that the interaction between vortices is actually very weak. Therefore, the short-range translational order for the vortex glass state may not be formed. Finally, we must mention that the μ value (0.79) obtained here is, however, rather different from the theoretical value ($\mu=\frac{1}{7}$) for the single vortex creep.[4] Future experiments at higher magnetic fields may help to find the cause of this discrepancy.

V. CONCLUSIONS

By taking a general assumption for the activation energy $U(j_s,T)=U_c(0)g(t)f[j_s/j_c(T)]$, we obtained the $U(j,T=0\ \mathrm{K})$ relation by scaling all the E-j_s curves of different temperatures to one master curve, where the important parameter $c=\ln(v_0B)$ was predetermined and $j_c(T)$ and $U_c(T)$ were calculated from the experimental data by the so-called GIS method. Finally, the resulting $U(j,T=0)$ relation was compared with several theoretical models. It seems that the logarithmic model for $U(j)$ is inapplicable to our data, while the collective-pinning model interprets the data quite well with the exception that the μ value found here is different from the theoretically predicted values.

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