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CRITICAL STATE OF TYPE-II SUPERCONDUCTORS*

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Flux Motion in the Self-Organized Critical State of Type-II Superconductors

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Abstract

A model based on the concept of self-organized criticality (SOC) is proposed to describe the nonequilibrium flux dynamics of type-II superconductors. The nonlogarithmic magnetic relaxation of high- T_c superconductors is argued as an indication of the SOC. The magnetic relaxation of single crystals of $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$, $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x$ as well as Nb is studied systematically. The relevance to models of the classical Anderson-Kim flux creep, thermally activated flux flow (TAFF), collective creep and vortex-glass is discussed.

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The critical current j (or magnetization) of a hard superconductor decays with time¹, due to the dissipative motion of the Abrikosov flux lines. A mechanism for the flux depinning suggested by Anderson² is the thermally activated flux creep, in which the flux lines or bundles can hop between pinning sites via thermal activation. Earlier experimental verification of the Anderson flux creep theory by Beasley et al.³ was based mainly on the characteristic of logarithmic decay. As pointed out by Campbell and Evetts⁴, a more convincing experimental proof of the Anderson theory should be based on the temperature dependence of the relaxation rate. Such type of experiments were carried out by Mitin⁵, at very low temperatures, at which the observed creep rates were larger than the theoretical predictions. The magnetic relaxation is found to be more significant in the high- T_c superconductors (HTS). The so-called "irreversibility line"⁶⁻⁹ ("onset" of reversibility in the dc magnetization or the loss peak of the ac susceptibility) in the B vs T phase diagram has been the focus of debate about the vortex dynamics of the HTS. The Anderson flux creep theory is widely used to interpret the magnetic relaxation⁷ and the resistive transition^{10,11} of the HTS. The magnetic relaxation is found to depend logarithmically on time at low temperatures or fields (where j is large), but the logarithmic dependence breaks down when approaching the "irreversibility line" (where j is small)^{7,9}. Several models have been proposed to describe these "new" vortex phenomena, which include the thermally activated flux flow (TAFF)¹² and collective creep (CC)¹³. Both TAFF and CC models assumed the nonlogarithmic decay occurs in the regime of $j \ll j_0$, or well below Bean's critical state¹⁴.

In this Letter, we provide experimental evidence which shows that the nonlogarithmic decay occurs near the critical state $j \sim j_0$, not far below j_0 . We propose a new model based on the concept of self-organized criticality (SOC)¹⁵, which could give a natural account for the experiments. We present magnetic relaxation data obtained over wide temperature and field regions and time scales in

both conventional and high- T_c superconductors (Nb, $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ and $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x$ single crystals).

The single crystals of $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) and $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x$ (BSCCO) were grown by a flux method. The YBCO crystal is of size $1 \times 2 \times 0.1 \text{ mm}^3$, and the BSCCO crystal is of $2 \times 1.5 \times 0.1 \text{ mm}^3$. X-ray diffraction showed that the samples are of single phase, and the zero-field-cooled (ZFC) magnetization at 20 G exhibited a sharp transition at 90 K for the YBCO crystal and at 86 K for the BSCCO crystal. The Nb sample was cut from a Nb single crystal and ZFC magnetization at 20 G show a sharp transition at 9.1 K with a width less than 0.05 K. Some structural defects were induced in the Nb crystal by mechanical stress during cutting. This results in an appreciable pinning, but its pinning force is much weaker than that of a cold-worked Nb sample we studied. Magnetization measurements were carried out on a Quantum Design SQUID magnetometer. The magnetic fields were turned on in small steps (less than 0.02 T) to avoid overshoot. The last step of the field increment in a relaxation measurement is typically 0.01 T. The travel length of the sample in each scan was 2 cm for all measurements.

To avoid the controversy of the incomplete field penetration, we only discuss the relaxation measurements performed above the fully penetrating field H^* . Typical relaxation data at different fields (5 kG and 6.5 kG) on the YBCO sample is shown in Fig. 1. The field-ramping rates are the same in the two runs. The magnetization is normalized to the first data point and is plotted against $\ln(t)$. The time of the first data point is the time between the field-ramping stops and the end of the first scan. In Fig. 2, M vs. $\ln(t)$ of the BSCCO sample at 1 kG but at different temperatures (12 K and 30 K) is shown. The inset shows the same data but normalized to the first data point. Similar behaviors are also found in Nb. We studied two Nb samples, one was cut from a single crystal with weak pinning, another was cut from a cold-worked wire (data not shown). Both Nb samples show

the general behaviors of the HTS shown above. The normalized magnetization vs. $\ln t$ of the crystal Nb at the same temperature 5 K and field 3.4 kG, but with different field-ramping sequence, is shown in Fig. 3a. In the no-waiting sequence, the relaxation data was recorded at 3.4 kG during the hysteresis curve measurement. In the with-waiting sequence, all the field-ramping history were the same except a 12 hours pause at 3.3 kG before the last 100 G increment. The long time linear slopes of the two runs are the same, but the early time nonlinearity is apparent in the no-waiting sequence. This field-ramping history dependence of the nonlogarithmic relaxation also appears in some intermediate field region in the HTS (data not shown). In Fig. 3b, we see again the M vs. $\ln(t)$ is linear at 3 kG, but is nonlinear at 3.4 kG. The M vs. $\ln(t)$ becomes completely nonlinear in the whole measurement when the field is very close to the upper critical field H_{c2} ($H_{c2} \sim 5.5$ kG at 5 K). The magnetization of the crystal Nb also becomes reversible at a field slightly below H_{c2} .

From the experimental observations, we find the following general facts: (1) the nonlogarithmic decay occurs in both high- T_c and Nb superconductors; (2) the nonlogarithmic decay occurs in the early time of relaxation; (3) in our time window, the nonlogarithmic decay becomes apparent when j is small (at high field or temperature). These facts imply that, one will only see logarithmic decay (no matter how low the j), if the field-ramping was infinitesimally slow, because all the early-time relaxation events are already finished during the field-ramping. One should see nonlogarithmic decay (no matter how high the j), if the magnetometer can probe the very early time effect after a step-like field-ramping. If one takes the logarithmic decay as an indication of the Anderson-Kim creep, the earlier time nonlogarithmic decay means the system has not yet reached that regime.

Within the classical Anderson-Kim theory, Beasley et al.³ showed that a $1/t$ solution of the diffusion constant of flux creep gives a logarithmic decay of j . A perturbation term of $1/t^2$ in the diffusion constant gives a nonlogarithmic decay. It

was suggested³ that this type of nonlogarithmic decay disappears in too short a period of time to be detected. For the HTS, Kes et al.¹² proposed that the measured critical current density j is well below the Bean critical current j_0 , or $j \ll j_0$, due to thermally activated flux flow (TAFF). From the classical Anderson-Kim formula of flux hopping rate¹⁶ v : $v = v_0 \sinh(jU_0/j_0 kT) \exp(-U_0/kT)$, where v_0 is the attempt frequency, j , the apparent critical current density, j_0 , the Bean critical depinning current density, U_0 , the intrinsic pinning barrier. Kes et al. linearized¹² the $\sinh$ term by its argument, thus $-dj/dt \sim V \sim v \sim j$ and j should decay exponentially in the limit of $j \ll j_0$. Thus the TAFF model interprets the observed nonlogarithmic decay as a crossover from the Anderson-Kim $\ln t$ decay to an exponential TAFF decay, and the crossover time t_{cr} is of $(1/v_0) \exp(U_0/kT)$. Feigel'man et al.¹³ proposed the theory of collective creep (CC), in which the effective barrier U is a nonlinear function of current j , and predict that j decays as an inverse power of $\ln t$, or $j(t) = j_0 (U_0/kT \ln(t/t_0))^{1/\alpha}$, $t_0 \sim 1/v_0$. Fisher¹⁷ described a 3-dimensional vortex-glass phase and predicted a similar decay of j . It should be pointed out that Anderson's assumptions² that $U \gg kT$ and $U = U_0 (1 - j/j_0)$ also imply $j \ll j_0(1 - kT/U_0)$ or $j \ll j_0$. It was not clear, however, in the AK theory, how a pinned vortex state decays from the Bean critical state ($j = j_0$ and $U = 0$) to the Anderson-Kim regime ($j \ll j_0$ and $U \gg kT$).

The fundamental aspect of the thermal activation models is that the flux jumping events are spatially and temporally uncorrelated. Each flux line or bundle can jump out of its potential well with the same probability assisted by thermal energy kT . Here we propose that the new concept of self-organized criticality (SOC), which accounts for the spatial and temporal correlation of depinning flux lines, could shed light on many unsettled issues about the flux dynamics. Bak, Tang and Wiensfeld proposed¹⁵ that a dissipative open system, such as a sandpile or a pinned FLL, evolves into a critical state (as that of Bean's) through self organization.

At this self-organized critical state, the sandpile (or FLL) is at a minimally stable state with a scale-invariant structure, in which the spatial correlation length between the particles (flux lines) diverges. The spatial correlation governs the dynamics of the sand (flux) flow. By imposing the scaling picture to the critical state, Tang and Bak predicted¹⁸ that, below the critical point the sand-flow decays with time as $t^{\phi-1} \exp[-(t/t_{co})^{D(\tau-2)/z}]$, where ϕ , D , τ , z are dynamical exponents, t_{co} is a cutoff time scale which is of the longest "avalanche" time. In relaxation measurements, by applying this SOC picture to the decay of Bean's critical state, j ($\sim$ slope of sandpile) should decay as a stretched exponential function of time¹⁹, if the system is not far from the criticality (where a scaling approach is valid), or $j \sim j_0$,

$$j(t) = j_0 \exp(-(t/\tau)^\beta), \quad 0 < \beta < 1 \quad (1)$$

Hagen et al.²⁰ have shown that their torque relaxation of their HTS sample obeys Eq. (1). We find that this unusual form of relaxation not only exists in the high- T_c superconductors, it also exists in the conventional superconductors. When the magnetization ($\sim j$) is very small (at either high fields or temperatures) and M vs. $\ln t$ becomes completely nonlinear. Eq. (1) gives the best fit to the early-time data shown in Figs. 1, 2, 3 (fitting curve not shown). The fitting parameter τ is about 10^{-10} seconds, β is about $1/3$, and varies slightly from sample to sample. To further analyze the data and compare the different characteristics in various models, we introduce a new approach. For simplicity, we only discuss the temporal effect of the diffusion constant $D(t) \sim -dj/dt$. We define a quantity S_D , the decay rate of diffusion constant,

$$S_D = d \ln D / d \ln t. \quad (2)$$

One finds that S_D is very sensitive to the dynamics used in different models. In the following, the explicit expression of S_D for different models (including the SOC) can be easily derived:

Anderson-Kim (AK):	$S_D = -1,$
TAFF:	$S_D = -t/t_{cr},$
Collective creep (CC):	$S_D = -1 - t_0 (\alpha^{-1} + 1)/\ln(t/t_0),$
Self-Organized Criticality (SOC):	$S_D = -1 + \beta (1 - (t/\tau)^\beta).$

A schematic of S_D as a function of time for different models is shown in Fig. 4. To compare the time dependence of S_D , be aware of the time constant τ which has different scales and physical meanings in different models. In the TAFF model, $\tau = t_{cr}$ varies from the order of 10 to 10^{100} seconds, depending on the value of U_0 and kT are chosen. In the CC model, $\tau = t_0$ is of the order of 10^{-10} seconds. In the SOC model, τ is of the order of 10^{-10} seconds as extracted from the data. In the AK model, S_D is a constant of -1. For the TAFF model, S_D is a linear function of $-t$, with $S_D = -1$ at t_{cr} . For the CC model, S_D approaches -1 for large t (the divergence of S_D at t_0 is not important since it is assumed $t \gg t_0$). For the SOC model, S_D decays with a weak power of time. We can see that only the SOC model gives a general agreement with the data (as shown in Fig. 5).

Since S_D is the tangential slope of the plot $\ln(-dj/dt)$ vs. $\ln t$, a plot of $\ln(-dM/dt)$ vs. $\ln t$ should be straight forward for clarifying the dynamics of flux creep. A typical example is shown in Fig. 5, in which the time dependent magnetization of the BSCCO crystal is converted to $\ln(-dM/dt)$ vs. $\ln t$. The noisy behavior at long times is an artifact from digitally differentiating the M vs. t data. For the data of 12 K, at which M vs. $\ln t$ was nearly linear (shown in Fig. 2), S_D is about -1. At higher temperature, the negative tangential slope of the plot $\ln(-dM/dt)$ vs. $\ln t$ is slowly increasing with time. The change of the time-dependent slope is more significant at higher temperatures or magnetic fields at which the usual plot M vs. $\ln t$ becomes

more nonlinear. This behavior (shown by 30 K data) is clearly not the kind predicted by the AK, TAFF, and CC models, but can be fitted by the S_D formula (above) of the SOC with fitting parameters $\tau \sim 500$ sec. and $\beta \sim 1/3$. While it was difficult to distinguish the time-dependent behavior suggested by different models in the M vs. $\ln t$ plot, here we have developed a unique approach to analyze the relaxation data.

The universal feature of our relaxation data viz. that, for fixed temperature (or field), the magnetization relaxes logarithmically for large magnetization (hence large j_0 , at lower field or temperature), but follows a stretched exponential when j_0 is small (at larger field or higher temperature), can be easily understood in the picture of the SOC, and can be visualized as in a sandpile model. For two sandpiles with different critical slopes $\theta_{1,2}$ (due to different friction constants), for the same travel length, the ratio of the times for a particle to slide, will be $t_1/t_2 \sim (\sin\theta_2/\sin\theta_1)^{1/2}$. That is, larger the slope, the faster the system leaves criticality. The same physics should also apply to the Bean critical state. For a superconductor, if j_0 is small, it takes longer for the system to decay. If j_0 is large, the system decays faster. When the vortex state decays away from the Bean criticality j_0 far enough, the spatial correlation between the depinning flux lines may be reduced significantly, so that the scaling picture of the SOC becomes inappropriate. We are not certain, however, whether the logarithmic decay (for large j) in the measurement time window represents the long time behavior of Eq. (1), or indicates the Anderson-Kim dynamics.

Another characteristic associated with the scaling form of $j(t)$ is the flow rate $D(j)$, which manifests as a voltage in the transport measurements. Equation (1) implies the following D - j (or V - I) relation

$$D(j) \sim j/(\ln(j_0/j))^{1/\beta-1} \quad (3)$$

In the limit of $j \sim 0$, the resistivity $dV/dj \sim dD/dj \sim 0$ or a zero resistance state. In the intermediate region $j \gg 0$, but $j \ll j_0$, the V - j curve is linear. At $j = j_0$, $D(j)$ diverges, which implies²¹ the SOC. A diverging D at $j = j_0$ does not mean a diverging voltage. One expects the resistivity to saturate at the Bardeen-Stephen flux-flow resistivity²². An interesting feature of Eq. (3) is that it can also reproduce the I - V characteristics of a single Josephson junction near T_c observed by Anderson and Goldman²³. The earlier interpretation of the V - I curve was based on the thermally activated phase slippage ($d\theta/dt$) across the junction, though the measured voltage was much larger than the prediction of the Ambegaokar-Halperin theory²⁴. Tinkham¹⁰ extended this thermally activated phase slippage physics as a microscopic basis for the "giant flux creep"⁷, and interpreted the resistive transition of high- T_c superconductors. The resistance was argued as a large rate of phase slippage due to thermally activated flux motion. In the SOC model, however, the voltage is due to the noise-driven (not necessarily thermal) flux depinning (or phase slipping in a Josephson junction), and is with a driving force j too close to the criticality j_0 . We expect a j -dependent voltage noise measurement in the resistive state of the HTS could provide more insights to the resistive transition. The SOC model, in which the vortex motion is spatially and temporally correlated, gives a natural account¹⁵ for the flicker ($1/f$) noise²⁵ in the flux-flow state. The spatial uncorrelated thermal activated flux flow would probably result in a "white" noise spectrum.

In conclusion, we have applied the concept of self-organized criticality to the Bean critical state of the type-II superconductors. The unusual behavior (nonlogarithmic decay, irreversibility line, etc.) of the vortex state of the high- T_c superconductors is found also exist in the conventional type-II superconductor Nb. In the light of the SOC model, the thermal activation may not be a serious limitation on the capability of carrying a large current by a high- T_c superconductor.

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¹⁹In the eq. (10) of ref. 18, the flow rate $-dj/dt$ is of a stretched exponential times a power t^α , $\alpha < 1$. Integrating over that equation and neglecting the power term in the final expression gives eq. (1).

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Figure Captions

Fig. 1 Normalized magnetization vs. $\ln(t)$ of the YBCO crystal, $T = 85$ K, $H = 5.0$ and 6.5 kG with the same field-ramping rate. Fields are directed normal to the basal plane of the crystal. The straight lines indicate the $\ln t$ dependence.

Fig. 2 M vs. $\ln(t)$ of the BSCCO crystal, $T = 12$ K and 30 K, $H = 1$ kG // c . Inset: Normalized magnetization vs. $\ln(t)$ of the same data.

Fig. 3 (a) Normalized magnetization vs. $\ln(t)$ of the Nb crystal, $T = 5$ K, $H = 3.4$ kG with different field-ramping histories. The upper curve is taken with the waiting sequence, the lower curve is by the no-waiting sequence (see text).

(b) Normalized magnetization vs. $\ln(t)$ of the Nb crystal, $T = 5$ K, $H = 3.0$ and 3.4 kG. The field-ramping rates are the same for the two curves.

Fig. 4 A schematic of S_D vs. time for different theoretical models. The time is normalized to the characteristic time constant of each model. A-K: Anderson-Kim flux creep; CC: Collective Creep; SOC: Self-Organized Criticality; TAFF: Thermally Activated Flux Flow.

Fig. 5 $\ln(-dM/dt)$ vs. $\ln(t)$ of the BSCCO crystal, $T = 12, 20, 30$ K, $H = 1$ kG // c . The raw data of 12 K and 30 K are shown in Fig. 2 (above).

Fig 1

Fig 2

Fig 3 a

Fig 3 b

Fig 4

Fig 5

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