# Ordering of the creeping vortex system in $Bi_2Sr_2CaCu_2O_{8+\delta}$ single crystals at low temperatures

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Zero-field-cooling magnetization relaxation measurements performed on $Bi_2Sr_2CaCu_2O_{8+\delta}$ single crystals in the low-temperature domain, with the external magnetic field oriented parallel to the *c* axis, reveal the ordering of the creeping vortex system due to the macroscopic currents induced in the sample. This explains many apparently conflicting results concerning the vortex phase diagram of disordered high-temperature superconductors, which have led to the dichotomy elastic vortex glass—plastic vortex assembly for the vortex phase at high magnetic fields.

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The static and dynamic properties of the vortex matter in the presence of pinning in the high-temperatures superconductors (HTS's) with pronounced anisotropy, such as $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Bi-2212) single crystals, has attracted much interest. One of the most remarkable aspects of the vortex phase diagram in the conditions of relevant pinning in the low-temperature $T$ domain is the occurrence of a quenched-disorder driven transition between a quasiordered vortex phase in the low magnetic field $H$ range (the Bragg glass, stable against dislocation formation) and a disordered vortex phase at higher $H$.[1,2]

This transition manifests itself by a second peak (SP) on the dc magnetization curves, representing a sudden increase of the irreversible magnetization with increasing $H$,[3] due to a better accommodation of vortices to the pinning centers in the disordered vortex phase.

The presence of the Bragg glass is supported by neutron scattering experiments,[4] but the nature of the vortex phase in the high $H$–low $T$ domain is still not clear. The existence of a three-dimensional (3D) elastic vortex glass[5] in Bi-2212 single crystals at high $H$ was first proposed in Ref. 6, based on the fact that the flux dynamics at low $T$ was found to be elastic (collective). The elastic pinning barriers diverge when the current density $J \rightarrow 0$.[7] For the similar highly anisotropic $Tl_2Ba_2CaCu_2O_8$ single crystals, the dc magnetization relaxation at low $T$ suggests elastic vortex creep, as well.[8] In the case of $Tl_2Ba_2CaCu_2O_8$ thin film specimens, it was reported[9] that a 3D elastic vortex glass persists up to $H \approx 7$ kOe, well above the crossover field $B_{3D\text{-}2D} = \Phi_0/\gamma^2 s^2 \approx 500$ G (where $\Phi_0$ is the magnetic flux quantum, $\gamma$ is the anisotropy parameter, and $s$ is the distance between the superconducting Cu-O layers). At least for HTS's with pronounced anisotropy, the presence of a 3D elastic vortex glass is surprising, since above $B_{3D\text{-}2D}$ the two-dimensional (2D) vortex fluctuations become essential, and the vortex system is highly unstable to defect formation.[10,11]

On the other hand, nondiverging (defect-mediated) plastic vortex pinning barriers have been observed in dc magnetization relaxation experiments at relatively high $T$ (Ref. 12) for $H$ above the SP field (of the order of a few hundreds of Oe), and it was surmised[12,13] that the vortex pinning barriers in Bi-2212 may remain nondiverging in the low-$J$ limit at all temperatures for $H$ above the SP.

At the same time, it is known that a moving vortex system is more ordered than a static one.[14,15] The nonequilibrium states of driven vortices in the presence of random pinning received considerable attention,[16] and new dynamic vortex phases, such as a moving Bragg-glass phase or crystal, and a smectic-flow phase have been predicted. Most of these studies deal, however, with a rapidly driven vortex system, i.e., for $J$ above the critical current density $J_c$, and/or in the vicinity of the thermally induced vortex-lattice melting.

In this work, it is shown that the dichotomy "*elastic vortex glass–plastic vortex assembly*" for the high $H$–low $T$ vortex phase in HTS's arises from *the ordering of the creeping vortex system during experiments*. Here we discuss the zero-field-cooling (zfc) dc magnetization measurements (widely used for the investigation of the vortex phase diagram of HTS's) performed on Bi-2212 single crystals at low $T$ and relatively high $H$. A nonmonotonic $J$ dependence of the normalized vortex-creep activation energy for $H \gg B_{3D\text{-}2D}$ is reported. This indicates the ordering of the creeping vortex system due to the macroscopic currents in the sample, which, at low $T$, considerably reduce the effective pinning energy.

The investigated specimens are overdoped Bi-2212 single crystals grown by the traveling solvent floating zone method,[17] with the zero-field critical temperature $T_c = 87$ K, and a transition width of $\approx 1$ K. The characteristic sample dimensions were $\sim 1 \times 0.5 \times 0.04$ mm$^3$. The magnetization $M$ (identified with the irreversible magnetization) was measured in zfc conditions with $H$ parallel to the $c$ axis, using a commercial Quantum Design (SQUID) magnetometer in the RSO mode, with the amplitude of 1 cm and the frequency of 1.5 Hz. The results obtained by decreasing the amplitude to 0.5 cm and/or increasing the frequency to 4 Hz are similar.

For the measured Bi-2212 samples, the SP on the dc magnetization curves $M(H)$ appears in the $T$ interval between $T_0 \approx 20$ K and $\approx 35$ K, with the onset field $H_{on}$ increasing when approaching $T_0$ from above.[18] The magnetization relaxation $M(t)$ was measured after increasing $H$. The relaxation time $t$ was considered to be zero when the magnet charging was finished, and the first data point was taken at $t = t_1 \approx 100$ s.

Figure 1 (main panel) shows the $M(H)$ curves at $T = 5$ and 10 K (where the SP is absent), and the inset to Fig. 1 illustrates $M(t)$ in log-log scales for $H = 50$ kOe and $T$ between 3

 

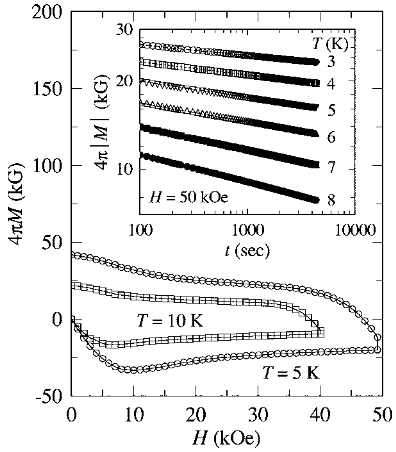


FIG. 1. Zero-field-cooling dc magnetization curves $M(H)$ for overdoped Bi-2212 single crystals, in $H$ parallel to the $c$ axis, at $T=5$ and 10 K, where the second magnetization peak is absent. The inset illustrates the magnetization relaxation $M(t)$ after increasing $H$ to 50 kOe, for $T$ between 3 and 8 K (step 1 K, log-log plot).

and 8 K. The $M(t)$ data at low $T$ are often used to "construct" the $J$ dependence of the actual vortex-creep activation energy $U$ by plotting the quantity $-T[\ln(dM/dt)-AT]\propto U$ vs $J\propto|M|$, with $A=$constant.[19] The pinning potential is practically $T$ independent at sufficiently low $T$, where the main role of an increasing $T$ is to supply lower $J$ values (due to an enhanced overall relaxation in the $t$ interval up to $t_1$).

Figure 2 (main panel) illustrates the attempt to put the relaxation data sets from the inset to Fig. 1 on the same curve, in the above scales. As can be seen, the data cannot be simultaneously aligned. The misalignment is not large (and is sometimes neglected), since $U(J)$ is dominated by the reduction of the pinning barriers with increasing driving force (lowering $T$), and may be influenced by extrinsic effects, such as the barrier distribution.[20]

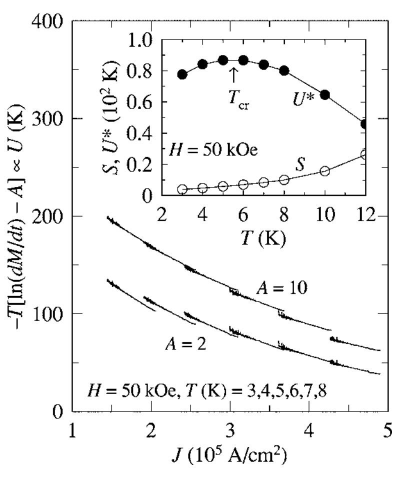


FIG. 2. Main panel, an attempt to construct the $J$ dependence of the actual vortex-creep activation energy $U(J)$ with the relaxation data $M(t)$ from the inset to Fig. 1 following Ref. 19, by plotting the quantity $-T[\ln(dM/dt)-AT]\propto U$ vs $J\propto|M|$, with $A=$constant. As exemplified for two $A$ values, the data sets cannot be simultaneously aligned, suggesting a change in the vortex dynamics between 5 and 6 K. The inset shows the $T$ dependence of the averaged normalized magnetization relaxation rate $S=-\Delta\ln(|M|)/\Delta\ln(t)$, and that of the related averaged normalized vortex-creep activation energy $U^*=T/S$ for $H=50$ kOe. A maximum in $U^*(T)$ appears at $T=T_{cr}\approx 5.5$ K (indicated by an arrow).

From the $M(t)$ data in double logarithmic scales an averaged normalized magnetization relaxation rate[21] $S=-\Delta\ln(|M|)/\Delta\ln(t)$ and the related averaged normalized vortex-creep activation energy $U^*=T/S$ were determined. Shown in the inset to Fig. 2 are the $S(T)$ dependence and the resulting $U^*(T)$ variation for $H=50$ kOe. The striking feature is the occurrence of a maximum in $U^*(T)$ at high $H$. The temperature $T_{cr}$ for the $U^*(T)$ maximum corresponds to the $T$ value where the misalignment in the plot from the main panel of Fig. 2 appears, indicating a change in the vortex dynamics. The analysis and discussion below will be focused on the nonmonotonic variation of $U^*$ and the meaning of the $T_{cr}(H)$ line in the $H$-$T$ diagram.

Following Ref. 22, the activation energy $U$ was parametrized in Ref. 9 as $U(T,H,J)=(U_c/p)[(J_c/J)^p-1]$, where $U_c$ is the characteristic pinning energy, whereas $p(T,H,J)$ was identified with the collective pinning exponent $\mu>0$ in the case of elastic creep, and $p<0$ for plastic vortex creep. From the general vortex-creep equation[23] $[U=T\ln(t/t_0)$, with the macroscopic time scale for creep[11] $t_0\propto T/(H|\partial U/\partial J|)]$, and $U^*(T,H,J)=-Td\ln(t)/d\ln(|M|)$, one obtains $U^*(T,H)=U_c+p(T,H)T\ln(t_w/t_0)$, where $t_w$ is the time window of the experiment. This can explain the nonmonotonic $U^*(T)$ variation from the inset to Fig. 2, as a crossover between plastic creep at high $T$ and elastic creep in the low-$T$ range.[9] Since the above relation for $U^*$ does not contain $J$ as explicit variable, the behavior of the $U^*(T)$ maximum with increasing $H$ was interpreted in Ref. 9 as a field induced suppression of the vortex glass temperature in highly anisotropic $Tl_2Ba_2CaCu_2O_8$ thin films.

Alternatively, it is argued below that the observed nonmonotonic $U^*(T)$ is essentially the result of a nonmonotonic $U^*(J)$ variation. As pointed out in Ref. 24, $U^*(J)$ is very sensitive to the change of the intrinsic (model-dependent) $J$ variation of the pinning barriers.

Figure 3 (main panel) illustrates the $U^*(J)$ dependence obtained with the $M(t)$ data for $3\text{ K}\leq T\leq 10\text{ K}$ and $H=50$ kOe, and $J$ extracted with the Bean model.[11] The new feature is the maximum in $U^*(J)$ at $J(T_{cr})$, separating the low-$T$ elastic creep regime and the plastic creep at high $T$. Note the $U^*(J)$ variation inside the $M(t)$ data set (where $H$ and $T$ are constant). The $M(t)$ data sets at $H=30$, 15, 5, and 1.5 kOe (not shown) give similar results, and $T_{cr}$ increases with decreasing $H$. It is the $U^*(J)$ maximum which generates the nonmonotonic $U^*(T)$ variation at low $T$, where the pinning potential does not depend on $T$. With the same $U(T,H,J)$ relation, but keeping $J$ as explicit variable, for the elastic creep domain it results $U^*(J)=U_c(J_c/J)^{\mu}$. The exponent $\mu$ was extracted using the relaxation data at $T=3$ and 4 K, as shown in the main panel of Fig. 3. The value $\mu\approx 0.5$ determined this way for $H=50$ and 30 kOe indicates 2D collective pinning (large vortex bundle[25]), i.e., the independent ordering of pancake vortices in the Cu-O layers with increasing $J$ (decreasing $T$ in standard zfc magnetization

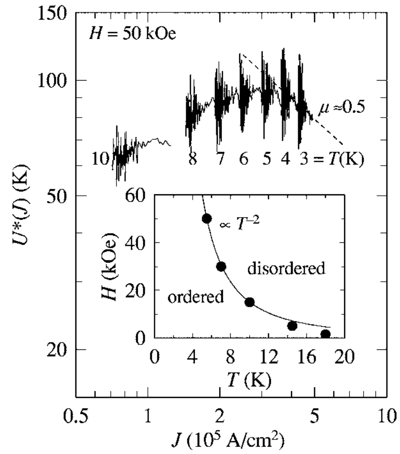


FIG. 3. Main panel, current density $J(\propto|M|)$ variation of the normalized vortex-creep activation energy $U^*(J) = -Td\ln(t)/d\ln(|M|)$ (log-log scales) obtained from the $M(t)$ data sets for $H=50$ kOe at $T$ between 3 and 8 K (step 1 K), and $T=10$ K. Note the maximum in $U^*(J)$ at $J(T=T_{cr}\approx 5.5$ K), separating the elastic-creep regime at low $T$, and the plastic creep at high $T$. The collective pinning exponent $\mu\approx 0.5$ was determined using the relaxation data at 3 and 4 K. The inset illustrates the $T_{cr}(H)$ values, describing the dynamic order-disorder separation line in the $H$-$T$ diagram. For our dc magnetic measurements, the separation line at low $T$ is close to $H\propto T^{-2}$.

measurements). The exponent seems to increase with decreasing $H$,[8] which would indicate a possible dynamic recoupling of the pancake vortices in adjacent layers.

The decrease of $U^*$ with decreasing $J$ in the plastic creep domain (see the main panel of Fig. 3) arises mainly from the definition of $U^*$, which implies that for every small $J$ segment the (nondiverging) $U(J)$ variation in the plastic creep region[26] is approximated by $U(J)=U_0\ln(J_c/J)$, where $U_0=U^*=U(J=J_c/e)$. It is easy to show, graphically, that if one uses this approximation for a nondiverging $U(J)$ the resulting $U_0$ (i.e., $U^*$) decreases with decreasing $J$.

The determined $T_{cr}(H)$ values plotted in a $H$-$T$ diagram describe a line, delimiting the creeping ordered and the disordered vortex phases, as shown in the inset to Fig. 3. For our dc magnetic measurements, the separation line at low $T$ is close to

$$H = aT^{-2}, \tag{1}$$

where $a=1.5\times 10^3$ kOe K$^2$. This relation can be derived by considering that for $J=J(T_{cr})$, corresponding to the $U^*(J)$ maximum (Fig. 3, main panel), the actual pinning energy equals the elastic energy $E_{el}\propto H^{-1/2}$.[11] This is similar to the energy balance equation for the static quenched-disorder driven order-disorder transition,[2] with the difference that the static pinning energy is substituted by the $J$-dependent pinning energy $U_p(H,T,J)\propto U(H,T,J)$. For a limited relaxation time window, the vortex-creep activation energy $U(H,T,J)\propto T$ (neglecting the variation of $t_0$), which immediately leads to the above relation. The deviation at high $T$ (see the inset to Fig. 3) may be due to the decrease of $t_0$, caused by the increase of $|\partial U/\partial J|$ with decreasing $J$ (increasing $T$), as well as to the fact that $U$ involves the intervortex interactions. It is worthy to note that the dynamic order-disorder line should depend on the measuring technique, at least through $t_w$. In the case of our magnetization measurements, the order-disorder line crosses the point of coordinates $B_{3D\text{-}2D}$ and $T_0\approx 20$ K, merging with the $H_{on}(T)$ variation observed in Ref. 18.

In summary, the nonmonotonic $U^*(J)$ dependence determined in zfc magnetization relaxation measurements for Bi-2212 single crystals indicates the current-induced ordering of the vortex system, which at high $H$ behaves like independently creeping 2D Bragg-glass phases in the Cu-O layers. This phenomenon appears at low $T$, where the macroscopic currents induced in the sample considerably reduce the pinning energy. In these conditions, the intervortex interactions become important,[27] and for a limited voltage sensitivity can mimic a transition toward a static elastic vortex-glass state.[9]

The dissipation process in the high-$H$ statically disordered vortex phase in the presence of a transport current starts with defect-mediated (plastic) vortex creep at low $J$,[28] followed by elastic creep at higher $J$. The ordering of the creeping vortex system explains the observation of diverginglike pinning barriers well above the crossover field $B_{3D\text{-}2D}$,[8,9,29] or above the SP field (for static conditions),[30] as well as the disappearance of the SP on the zfc dc magnetization curves of highly anisotropic HTS's at low $T$.[18] The ordering of the creeping vortex system is strongly supported by the "anomalous" shift of the onset field for the SP to higher $H$ values with decreasing $T$ in the low-$T$ range.[18,31]

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