Two-dimensional vortex dynamics in decoupled $YBa_2Cu_3O_7/PrBa_2Cu_3O_7$ superlattices

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Two-dimensional (2D) vortex dynamics is studied in $YBa_2Cu_3O_7/PrBa_2Cu_3O_7$ superlattices by measuring the current-voltage (*I-V*) characteristics. In the high current limit, the 2D collective creep is observed with an activation energy $U(j)\propto j^{-\mu}$, with j the current density. The exponent μ is 0.8 at low temperatures and $\mu=0.2$ at high temperatures. A dislocation-mediated vortex melting occurs when the temperature is increased. In the low current limit, the exponential growth of the energy barrier for elastic motion is prohibited by the plastic deformation of vortices. The observed plateau in the resistive transition is attributed to the possible quantum tunneling of vortices.

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I. INTRODUCTION

The unusual resistive and magnetic behavior of high-temperature superconductors has greatly stimulated the efforts to understand the influence of disorder, thermal fluctuations, and dimensionality on the vortex dynamics in those materials.[1] The theory of collective vortex creep describes the thermally assisted motion of vortices in a random potential caused by quenched disorder.[2] Central to this model is the elasticity of the vortex system. Due to the finite values of the compress modulus c_{11} and the shear modulus c_{66}, when moving from one metastable position to another one, the vortices have the tendency to jump in bundles with a radius R_c in order to balance the elastic energy and the energy related to the Lorentz force. The characteristic energy barrier for activation, $U(j)$, increases exponentially with decreasing current because R_c increases. However, this picture breaks down when *plastic motion* of vortices is taken into account.[3] For a two-dimensional (2D) vortex lattice (VL), due to the finite energy for creation of dislocations or dislocation pairs, a plastic motion of vortices is favorable when the elastic energy barrier exceeds a certain threshold. The plastic creep sets a cutoff for the exponential growth of $U(j)$. Linear resistivity will persist in the plastic regime and thus there is no real vortex glass in a 2D VL. The plastic motion of vortices has been studied by a lot of research groups,[4] and recently it was directly observed by Lorentz microscopy in Nb films with artificially introduced pinning centers.[5]

At low temperatures, *quantum tunneling* of vortices is expected to replace thermal activation as the dominant dissipation mechanism.[6] Tunneling of vortices has been observed in magnetization and transport measurements,[7,8] and it was shown that the magnetic relaxation rate did not extrapolate to zero when $T=0$ K.[7] A transition from thermally activated to a temperature independent resistance was observed in $Mo_{43}Ge_{57}$ ultrathin films and was attributed to the quantum creep of vortices.[8]

Plasticity of the vortex lattice and tunneling of vortices are two possible contributions prohibiting the existence of a vortex glass state, which is important for future applications. In this paper, we report on the vortex dynamics in a $YBa_2Cu_3O_7/PrBa_2Cu_3O_7$ (YBCO/PBCO) superlattice at different temperatures, magnetic fields, and transport currents. We observed a crossover from collective (elastic) creep to plastic creep of vortices when the elastic limit of the VL was reached. A dislocation-mediated melting of the VL was observed in the high current limit. And the possibility of quantum creep of vortices at low temperature is discussed. Our results provide a comprehensive picture for the vortex dynamics within a 2D VL in artificial superconducting superlattices.

II. EXPERIMENTAL

The *c*-axis oriented [YBCO(12 Å)/PBCO(144 Å)]$_{25}$ superlattices were fabricated by *in situ* magnetron sputtering.[9] The x-ray diffraction showed satellite peaks up to the third order, indicating the high layering quality of the sample. High-resolution transmission electron microscopy revealed the sharp interface between YBCO and PBCO layers.[10] The zero-field transition temperature T_{c0} of the sample we reported here was about 31.5 K. The superlattice was photolithographically patterned into a strip with a width of 0.1 mm and a length of 1 mm; dc *I-V* characteristics were recorded with voltage as a function of current. The resistive transition $\rho_{ac}(T)$ was measured by a four-terminal ac locking-in technique with an excitation current of 1 μA at a frequency of 17 Hz. Resistivity and current density were calculated with the total thickness of YBCO layers. The magnetic field was generated by a 15-T Oxford superconducting magnet. During the measurements, the field was kept perpendicular to the film plane and the temperature stability was better than 10 mK.

Figure 1 shows representative *I-V* characteristics measured at temperatures ranging from 2 to 32 K at an interval of 2 K in the magnetic field $H=2$ T. Each *I-V* curve was measured by averaging individual current and voltage characteristics four times at each temperature. Two features are clearly visible: (i) *in the high current limit* ($j>5\times10^3$ A/cm^2), the *I-V* curves show a downward curvature at the lowest tem-

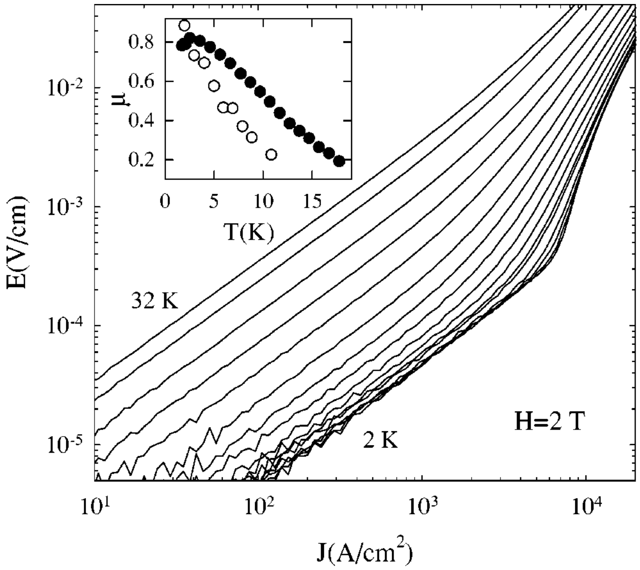

FIG. 1. Selected I-V curves at different temperatures ranging from 2 to 32 K with an interval of 2 K. Inset: temperature dependence of μ at $H=0.5$ T (filled circles) and 2 T (open circles), respectively.

perature; (ii) *in the low current limit*, linear I-V curves were observed for the whole temperature region. From the linear parts in the I-V curves, we can extract the linear resistivity ρ_L which is shown in Fig. 2 together with the measured $\rho_{ac}(T)$ in the Arrhenius plot. We find that ρ_L coincides with ρ_{ac} quite well. In the high-temperature regime, a thermally activated flux flow (TAFF) behavior is clearly visible in Fig. 2. The activation energies are extracted from the linear parts in the Arrhenius plot, and shown in the inset in Fig. 2. At lower temperatures, the resistivity deviates from the TAFF behavior and saturates.

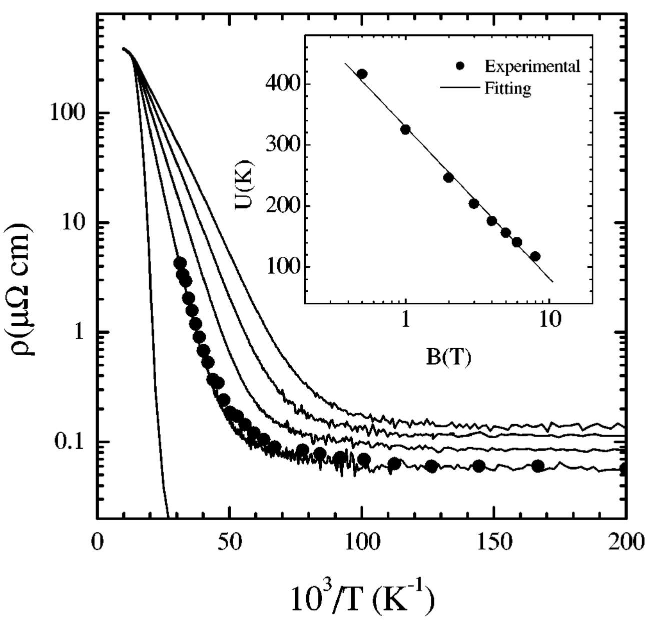

FIG. 2. Temperature dependence of ρ_{ac} at difference applied magnetic field. The solid lines are the experimental data. From left to right: $H=0$, 2, 4, 6, 8 T. Fitted circles are data points of ρ_L extracted from the linear parts in Fig. 1. Inset: Activation energies from a TAFF fit to $\rho_{ac}(T)$, filled circles are the obtained data, solid line is a fitting by $U=329.5-247.9\ln B$.

III. DISCUSSIONS

It was found by Brunner *et al.* that 96 Å of PBCO is enough to decouple the YBCO layers and that L_c of YBCO is of the order of hundreds Å.[11] In our case, due to the large PBCO layer thickness (144 Å), the YBCO layers are essentially decoupled and consist of only one unit cell (12 Å). For such a small superconducting layer thickness, L_c will be reduced to the thickness of the YBCO layer thickness, therefore, we can analyze the vortex dynamics in the 2D regime. The following 2D vortex dynamic scenarios can be used to describe the observed I-V characteristics.

It is expected that in a 2D disordered superconductor, collective creep of vortices plays an important role in the dissipation process. The 2D collective creep theory is applicable when the vortex bundle size $R_c=a_0(\varepsilon_0 d/U_{pc})(\xi/2a_0)^2$ is larger than $a_0\approx(\Phi_0/B)^{1/2}$, the vortex spacing. Here, $\varepsilon_0=\Phi_0^2/16\pi^2\lambda^2$ is the energy for a pancake vortex per unit length, d the length of the vortex along the field direction, ξ the coherence length λ the penetration depth and $U_{pc}\simeq(\gamma\xi^2 d)^{1/2}$ the collective pinning energy for a single pancake vortex.[3,12] Due to a lack of sufficient experimental data on such a system, an exact estimation of R_c is hard to make. However, we can still make a rough estimation by using the parameters reported for YBCO films. For a two-dimensional vortex system, U_{pc} can be expressed as $U_{pc}\simeq T_c(j_{pc}/j_0)(1-t)/G_i^{2D}$, where $G_i^{2D}=T_c/\sqrt{2}\varepsilon_0 d$ is the 2D Ginzburg number, j_{pc} and $j_o=cH_c/3\sqrt{6}\pi\lambda$ being the pancake depinning current density and depairing current density, respectively.[1] Substituting the expressions of U_{pc} and G_i^{2D} to that of R_c, we have $R_c=a_0(j_0/\sqrt{2}(1-t)j_{pc})(\xi/2a_0)^2$. With $j_0\approx 2\times 10^8$ A cm^{-2}, $j_{pc}\approx 10^6$ A cm^{-2}, $\xi\approx 15$ Å, and $a_0\simeq 320$ Å for a field of 2 T, we have $R_c\simeq 10a_0$.[1] Therefore, we will assume that the 2D collective creep theory is applicable to the system we studied here. When vortices undergo collective creep, the current density dependent energy barrier is $U(j)\propto(j/j_0)^{-\mu}$, where j_0 is a characteristic current density. Within the TAFF approximation, the I-V curves due to $U(j)$ is[2]

$$E=E_0 e^{-(j/j_0)^{-\mu}}. \qquad (1)$$

Following the methods used by Dekker *et al.*,[13] we reformulate Eq. (1) by taking its logarithm and fitting the I-V curves with downward curvature to it, with $\ln E_0$, j_0^{μ}, and μ as fitting parameters. The exponent μ as a function of temperature is shown in the inset in Fig. 1 and changes gradually from $\mu=0.8$ to $\mu=0.2$ as the temperature increases. This result agrees well with the observation of Dekker *et al.*[14] The $\mu(T)$ variation is closely related to the bundle size of the vortices, as pointed out by Vinokur *et al.*,[12] depending on the current density, temperature, and magnetic field. For medium-sized vortex bundles [defined by $j_{lb}\approx j_c(R_c^3/\Lambda a_0^2)^{8/15}<j<j_{mb}\approx j_c(R_c/\Lambda)^{8/15}$, where

$\Lambda = \lambda\ \coth(d/2\lambda)$], an exponent $\mu = 13/16 \approx 0.8$ is expected. When $j < j_{lb}$, the bundle size increases and the exponent $\mu \approx 0.5$. For a constant magnetic field the bundle size is $R_c \propto \xi^2/\lambda^2 U_{pc}$. At temperatures not too close to T_c, where ξ and λ are nearly constant, U_p decreases with increasing temperature and therefore, R_c will increase. As a result, the bundle size changes from medium to large, thus decreasing μ value, and explaining the decrease of μ.

As the temperature increases, dislocation-mediated vortex melting occurs when[15,16]

$$A c_{66} a_0^2 d = 4\pi k T_m^{2D}, \tag{2}$$

where $A \sim 0.4–0.7$ is a parameter due to the renormalization of c_{66} from the defects in the VL and the nonlinear lattice vibration.[16] Here, we identify the melting temperature T_m^{2D} as the one at which the *I-V* curves at high currents change from positive to negative curvature. We obtain $T_m^{2D} = 12$ K for $H = 2$ T. Shear modulus c_{66} is given by[17]

$$c_{66} = B\Phi_0/(16\pi\lambda^2\mu_0). \tag{3}$$

Inserting Eq. (3) into Eq. (2), with the typical values for parameters of YBCO, we obtain $A \approx 0.43$, which agrees well with the previously obtained value.[4,18]

Above the melting temperature, the vortices are in a liquid state pinned by available pinning centers. The *I-V* characteristics can thus be described by the classic Anderson-Kim model with $E = E_0 \sinh(j/j_0)$.

The collective pinning theory depends critically on the elasticity of the vortex system. Due to the finite energy for the generation of dislocations or dislocation pairs in a 2D VL, an infinite energy barrier in the elastic limit is not expected. When the energy barrier reaches the characteristic energy for a small dislocation pair, it will be cut off and a crossover from elastic to plastic motion should occur.[3] The motion of dislocation pairs can be well described by TAFF with an activation energy nearly independent of the current density. Therefore, linear *I-V* curves are expected in the plastic region. As a result, the linear resistivity does not vanish and a 2D vortex glass *does not exist* at any finite temperature. This is what we observed in our *I-V* curves. When below a current density $j \sim 5 \times 10^3$ A/cm^2, the *I-V* curves become linear over more than two orders of magnitude in the current and voltage. The typical energy for a dislocation pair made up of two edge dislocations separated by a_0 is[3]

$$U_e = \frac{\Phi_0^2 d}{16\pi^3\lambda^2(T)} \ln(B_0/B). \tag{4}$$

With $\lambda \approx 1400$ Å, we have $U_e \approx 500$ K. This value is roughly consistent with the activation energy we extracted from the linear parts [the region with $\rho(H,T)$ between 1 and 90 $\mu\Omega$cm] of the $\rho_{ac}(T)$ in the Arrhenius plot in Fig. 2. The explanation based on plastic motion mediated by the generation of dislocation pairs is further supported by the lnB dependence of the activation energies as shown in the inset in Fig. 2. Such a lnB dependence of activation energy has been previously reported by several groups.[11,19] It is noticed that there is an upward curvature at around 3 T, which could possibly be due to the increased interaction between vortices when the field is increased. So, we conclude that in the low current limit, the dissipation is governed by plastic motion of dislocation in the 2D VL.

The above picture is consistent with Monte Carlo simulations of the VL subject to random disorders under a driving force.[20] It is found that at large driving force, the vortex motion is elastic. As the driving force decreases, the VL becomes defective and the dissipation is dominated by plastic motion.

Measurements performed in ultrathin YBCO films reported *I-V* curves with upward curvature, while the glass correlation length ξ_{VG} diverged at $T = 0$ K.[14] The authors concluded that a 2D vortex glass could not exist at any temperature. This is consistent with our results. However, in our case, we observed, in addition, a crossover from downward to upward curvatures. This discrepancy may be due to the different length scale involved in the experiments. It is well known that from a topological point of view, by applying a current density j, a vortex pair of a characteristic size $L_j = (ckT/j\Phi_0)^{1/2}$ is excited.[21] The current probes the vortex dynamics on a length scale of L_j. With $j \sim 5 \times 10^3$ A/cm^2, $T \sim 10$ K, we obtain $L_j \sim 1000$ Å $< R_c$. It is possible that we are probing the collective creep behavior in the high current limit at low temperature. In the case of YBCO ultrathin film, the length scale was probably larger than R_c and thus, liquidlike behavior was detected.

As shown in Fig. 2, the resistivity gradually deviates from the activated behavior below $T \approx 10$ K and levels off toward a constant value.[22] The resistivity at $T < 10$ K increases with increasing magnetic field. To explain such a plateau within the framework of TAFF requires an activation energy U_0 which increases linearly with T, i.e., $U_0 \propto \alpha T$. We are not aware of any mechanism that provides such a kind of activation energy. We noted that similar resistive plateau in the ρ-T curves was observed by Ephron *et al.* in $Mo_{43}Ge_{57}$ thin films.[8]

Generally, it is believed that quantum tunneling dominates the dissipation at temperatures below several Kelvin and that the resistivity from a quantum creep is nearly independent on temperature.[23–27] We are unable to explain why quantum tunneling is observed at such a high temperature where the thermal energy of vortices is quite high. One possibility is a large normal resistivity ρ_n and a small coherence length ξ. The tunneling rate γ is determined by the effective Euclidean action S_E^{eff} for the tunneling process, $\gamma \propto \exp(-S_E^{eff}/\hbar)$, and $S_E^{eff}/\hbar$ in turn is proportional to $(\hbar/e^2)\xi_c/\rho_n$.[23] Therefore, tunneling is favored by a small coherence length ξ_c and a large normal state resistivity ρ_n. It was mentioned by Blatter and Geshkenbein[24] that superconductors with $\rho_n/\xi_c \gtrsim 1$ kΩ are good candidates for the observation of quantum creeep. In our case, with $\rho_n \sim 400$ $\mu\Omega$ cm and $\xi(c) \sim 3$ Å, our sample would have a ρ_n/ξ_c value of 10 kΩ. Ephron *et al.*[8] reported that they were only able to observe such a plateau in thin films with large sheet resistance, which is consistent with the theoretical consideration.

The predictions about the temperature T_0 at which the dissipation shows a crossover from thermally activated to a

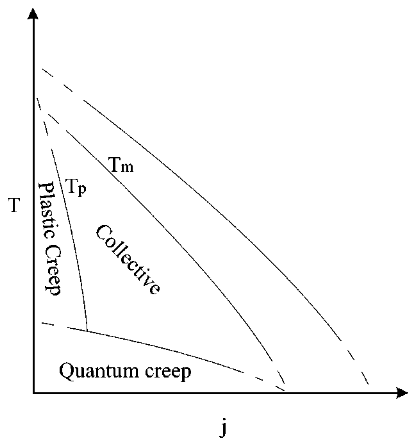

FIG. 3. A schematic j-T phase diagram for a 2D VL at a certain magnetic field. The line T_m is determined by Eq. (2), and line T_p by Eq. (4). In the quantum creep and plastic creep regions, the linear resistivity does not vanish. The collective creep and Anderson-Kim creep regions are characterized by I-V curves with $E \propto e^{j^{-\mu}}$ and $E \propto \sinh(j/j_0)$, respectively.

quantum tunneling range from hundreds of mK to tens of Kelvin. Ivlev *et al.*[25] showed that, for the YBCO compound, a crossover from thermally activated to quantum tunneling dissipation happens at $T_0 \approx 50(j_c - j)^{1/2}/j_c^{1/2}$ K. Meanwhile, Stephen[26] found that $T_0 \approx 2$ K. Experimentally, a crossover temperature ranging from 3.5 K to 10 K was reported for TlBaCaCuO single crystals.[28] We note that Blatter *et al.*[23] found a suppression of the tunneling rate for increasing fields. In this case the interaction between vortices increases, the viscosity grows and more vortices are participating in a single tunneling event. This is consistent with our observations.

IV. CONCLUSION

Our main conclusions are summarized in Fig. 3 and show that the vortex dynamics is essentially dependent on the length scale. In the high current limit, 2D collective creep is observed which is elastic in nature and characterized by an energy barrier with an exponential growth with current density. The 2D VL undergoes a dislocation mediated melting at a temperature T_m^{2D}. In the low current limit, the exponential growth of energy barrier is cut off by the generation of dislocation pairs in the VL and their subsequent motion or quantum tunneling. The linear resistivity at low temperatures ($\lesssim 10$ K) is possibly governed by quantum tunneling of pancake vortices. With increasing Lorentz force, the driven VL recovers its elasticity at high current densities.

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