Observation of Quantum Dissipation in the Vortex State of a Highly Disordered Superconducting Thin Film

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We present transport measurements on a highly disordered, amorphous thin film superconductor that shows a crossover from activated to temperature independent resistance at low temperatures in the presence of a magnetic field. We interpret the data in terms of quantum tunneling of vortices, and find that within this interpretation the tunneling objects are not single vortices but rather dislocations and antidislocations of the vortex lattice, which nucleate quantum mechanically as unbound pairs. We note, however, that other interpretations of the data may be possible.

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The properties of vortices in 2D and highly anisotropic 3D superconductors have attracted widespread theoretical and experimental attention in recent years—motivated in part by the advent of high-T_c superconductivity [1]. Theoretical interest has focused upon a variety of intriguing new concepts predicted to govern the behavior of vortices in 2D superconductors as $T \rightarrow 0$, including the superconducting-insulating [2] and vortex-glass [3] transitions, new mechanisms for collective flux creep [4], and macroscopic quantum mechanical tunneling of vortices (MQTV) [1–5]. While MQTV is well established experimentally in single Josephson junctions [6], the experimental evidence in crystalline superconductors (high-T_c) [7], thin films [8], and 2D Josephson junction arrays [9] is uneven at best. The issue is important, since the observation of MQTV questions the existence of a true superconducting state $dV/dI|_{I=0} = 0$ in 2D for $H > 0$, even at $T = 0$.

In this paper we present transport measurements on a highly disordered, amorphous thin film superconductor in an applied magnetic field. The sample resistance shows a clear crossover from activated behavior at high temperatures to temperature independent finite resistance as $T \rightarrow 0$, indicative of a quantum mechanical process. We can explain the high-temperature regime in terms of a model in which thermal processes unbind pairs of dislocations and antidislocations of the locally well-defined vortex lattice; their subsequent independent and relatively unimpeded motion leads to the observed sample resistance. The most natural interpretation of the low-temperature quantum regime is to suppose that quantum rather than thermal fluctuations drive the same dislocation-antidislocation pair nucleation process. Thus in this interpretation, the "particles" that are tunneling are not single vortices moving from one pinning site to the next, as has been proposed previously [1,5]. Instead, each tunneling event involves the collective motion of many vortices and the particles that result do not exist beforehand—they are created in pairs in the tunneling process itself. Within this picture, however, the data imply a nonintuitive field dependence of the tunneling time. Therefore, we do not exclude the possibility of a novel quantum mechanical process or phase lurking behind our lowest-temperature data.

The sample for which we present data in this paper is a 30 Å thick amorphous $Mo_{43}Ge_{57}$ thin film, sandwiched between insulating layers of amorphous Ge on a SiN substrate. The sample has a high sheet resistance $R_\square \approx 1350\ \Omega$ and a superconducting transition temperature $T_c \approx 500$ mK. Previous measurements on films from the same extended family of samples grown in consecutive sputtering runs formed the basis for a detailed examination by us of the field-tuned superconducting-insulating (S-I) transition [10]. Details of the growth and characterization of the films, all of which show behavior characteristic of homogeneous, amorphous samples over all relevant length scales, were published previously [10,11].

We patterned the film into a four-probe structure seven squares long and measured it in our dilution refrigerator inside a shielded room using standard ac techniques. The data were taken at a measurement frequency $f_{ac} = 27.5$ Hz with an applied bias $I_{ac} = 10$ nA. The results, however, are independent of frequency from 2.75 Hz to at least 1 kHz and the ac current bias was always well within the Ohmic regime. The measurement setup employed SHE electronics to regulate the temperature of the mixing chamber and a PAR 124 lock-in amplifier to monitor the sample resistance; the instrumentation was carefully configured to preclude spurious noise effects. The most convincing evidence that heating or extrinsic noise do not affect these measurements, however, is in the systematics of the data itself, to which we now turn.

Figure 1 is an Arrhenius plot of the temperature dependence of the zero bias resistance of the sample R_0 for six values of the applied magnetic field ranging from 1.5 to 8.5 kOe. The data clearly show the expected activated behavior at high temperatures (straight line fits), followed by a leveling off of R_0 at low temperatures, suggestive of temperature independent quantum motion of vortices [12].

We first examine the high-temperature activated regime. Figure 2 shows the activation energies extracted from

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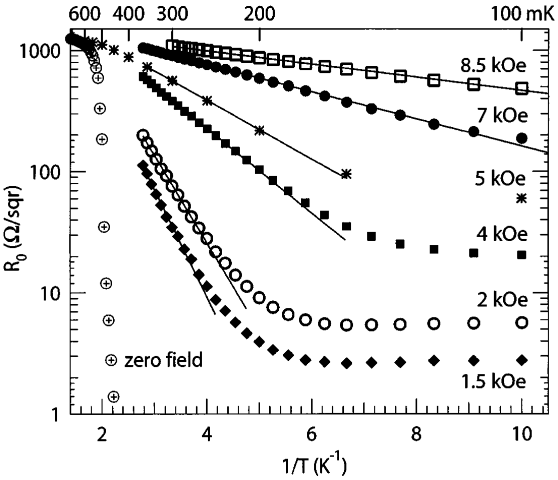

FIG. 1. Arrhenius plot of the zero bias resistance of the sample R_0 for six values of the applied magnetic field. The straight solid lines demonstrate that the behavior is activated in the high-temperature limit and then saturates at lower temperatures, likely indicating MQTV. The apparent slight minimum in the two lowest curves is real [12].

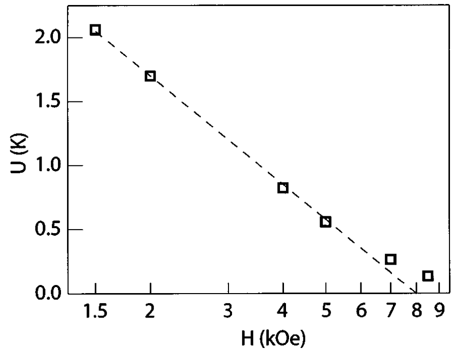

FIG. 2. Activation energies derived from the slopes of the solid lines in Fig. 1 shown on a semilogarithmic plot as a function of magnetic field. The dashed line is a fit by Eq. (1) and indicates that the sample resistance is governed by the formation of free vortex lattice dislocations.

the straight line fits in Fig. 1. The dependence of the activation energy $U(H)$ on the applied magnetic field is consistent with the form

$$U(H) = U_0 \ln(H_0/H), \quad \text{with } H_0 \approx H_{c2} \quad \text{and } U_0 = \varphi_0^2 d/256\pi^3\lambda^2, \tag{1}$$

extensively observed previously in lower resistance films [10,13] and expected in the collective creep regime [4]. In this pinning regime, the motion of dislocations of the 2D flux lattice dominates the sample resistance. Equation (1) gives the energy necessary to nucleate a free dislocation-antidislocation pair [14]. The creation of the pairs is the rate limiting step: Once nucleated, the dislocations move relatively freely since the energy required to overcome the pinning of the unbound dislocations is substantially less than that required to create the pairs initially [4]. The dashed line in Fig. 2 is the best fit by Eq. (1) for the first five points. The resulting best fit value for U_0 is within a factor of 2 of the predicted value—within the uncertainty of the penetration depth of the sample λ. The best fit value for H_0 is 8 kOe, somewhat below the critical field for the S-I field-tuned transition for the sample $H_c = 12.45$ kOe and the mean-field critical field H_{c2} estimated [10] to be about 14 kOe. The theory, however, is heuristic and neglects small factors of order unity in the derivation of the argument of the logarithm and the prefactor in Eq. (1), so the best fit value for H_0 is entirely reasonable. The breakdown of the physical justification for Eq. (1) as the argument approaches unity together with the proximity to the S-I transition presumably account for the deviation above 7 kOe.

Equation (1) should be valid when the effective size of a free dislocation $R_d \equiv a_0^2/\xi$ exceeds R_c, the Larkin correlation length, where a_0 is the intervortex spacing ($a_0 = 1000$ Å at $H = 2$ kOe) and $\xi \approx 150$ Å [10] is the vortex core size. R_c is the size of a region for which the root mean square displacement of vortices from a locally defined ideal lattice due to random pinning forces is less than ξ, and is the length scale which determines the critical current in the elastic collective pinning model [15]. Our sample is in the limit of strong disorder, and we expect $R_c \approx (1\text{–}2)a_0$ [13]. The justification for Eq. (1) in the limit when R_c is small is simple: The strain field induced by a dislocation is not sufficiently strong to displace vortices at a distance $r > R_d$ from the center of the dislocation into a different local minimum of the disorder potential, i.e., $\delta u(r > R_d) < \xi$. Thus, beyond this distance, the intrinsic disorder dominates the configuration of the flux lines and cuts off the long-range mutual interaction of a dislocation pair at the finite value $U \approx 2U_0 \ln(R_d/a_0)$. Substituting $a_0/\xi \approx \sqrt{H_{c2}/H}$ into the definition of R_d yields Eq. (1) [4].

R_c, which is small for our samples, does not determine the range of applicability of the elastic flux line lattice model. The relevant *elastic* correlation length of the flux lattice is the size of a region in which the cumulative root mean square displacement of vortex lines from their ideal lattice positions is less than the intervortex spacing $\delta u_\Sigma \leq a_0$ and is given by $R_{el} \approx R_c(a_0/\xi)^p$ with $p > 1$ [16]. Note that $R_{el} \geq R_d > R_c$. Thus the description of the vortex configuration in terms of a locally well-defined lattice with dislocations is meaningful here.

The dislocation-antidislocation nucleation process described above is well established. Our group has observed this same behavior in numerous (less resistive) samples and found excellent agreement between the theory outlined above and the experimental results, which have appeared in print previously [10,13]. The point of the above discussion is to show that the high-temperature behavior is consistent with a simple theoretical model that has successfully explained the vortex properties of a wide range of highly disordered superconducting thin films [17]. The difference between the present data and our earlier results is the flattening of the resistance at low temperatures, which we address next.

Another way to express the tendency of the resistance to flatten is to determine the "effective" temperature, as a function of the actual temperature, at which the sample would have to be in order to manifest the observed resistance at low temperatures by the activated process extrapolated from higher temperatures. That is, T_{eff} is defined implicitly as a function of T by

$$R_0(T,H) = \tilde{R}_0(H)\exp\{-U(H)/k_B T_{\mathrm{eff}}(T)\}, \qquad (2)$$

where the temperature independent values $\tilde{R}_0(H)$ and $U(H)$ are determined unambiguously by the straight line fits in the high-temperature activated region in Fig. 1. Figure 3 shows that the limiting value of T_{eff} depends on the field, demonstrating that the behavior is intrinsic to the sample and not the result of inadequate cooling or thermal anchoring. Furthermore, the deviation from activated behavior occurs at *lower* temperatures in the presence of higher magnetic fields, where the sample resistance and hence the power dissipated and the heating in the sample are greatest. This circumstance precludes any simple explanation of the data in terms of local heating effects. Figure 4 tallies the low-temperature limiting values of the resistance (R_0^{QL}) and the effective temperature (T_{eff}^{QL}), gathered from Figs. 1 and 3, respectively, as a function of the applied field.

We do not know of any calculation of the quantum tunneling rate of vortices that explicitly accounts for multivortex processes. We nevertheless still expect the low-temperature limit of the sample resistance, which is proportional to the rate of quantum nucleation of free dislocations, to have the usual general form in the limit of strong dissipation [5],

$$R_0^{QL} = \tilde{R}_0^{QL}(H)e^{-S_E/\hbar}, \quad \text{with } \frac{S_E}{\hbar} \propto \frac{\hbar/e^2}{R_\Box}S(H). \qquad (3)$$

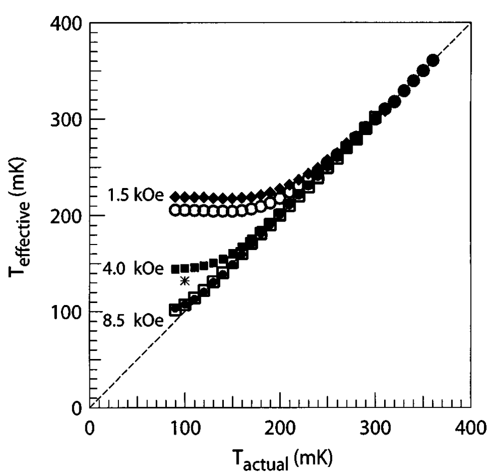

FIG. 3. The effective sample temperature necessary to explain the data of Fig. 1 in terms of thermally activated behavior extrapolated to include all of the data. The symbols refer to the same field values as in Fig. 1.

S_E is the Euclidean action, which has an unknown dependence on H embedded in the function $S(H)$. We write the resistance in this general form to call attention to the strong dependence on the damping level, proportional to $R_\Box^{-1}$, which is a straightforward consequence of the displacement of the vortices in the presence of a current. This strong dependence is evidently the reason that we did not see convincing evidence of similar low-temperature saturation behavior in samples studied previously. The sample discussed here is at the extreme limit both of a large body of well-characterized samples and of the capabilities of our thin film technology. All of the other samples presently available have a lower sheet resistance, or a composition such that T_c and the other relevant energy scales are lower. Some of these samples do show incipient saturation of the resistance, but below 90 mK, where we cannot confidently exclude the possibility of inadequate sample cooling. Further quantitative investigation of the quantum tunneling effects reported here will require a new set of samples, ideally with higher sheet resistances but not significantly reduced T_c's. Growing and characterizing *homogeneous* thin films with these characteristics is a difficult task, which we intend to undertake soon.

In the absence of a theory describing the action S_E for our system, we simply observe that the low-temperature saturation value of the resistance R_0^{QL} obeys the empirical form $R_0^{QL} = A\exp(H/H')$ quite well with $A = 1.05\ \Omega$ and $H' = 1.4$ kOe, shown as the solid line in Fig. 4. Numerous alternative forms will also fit the data with varying degrees of success, of course, so this result should be considered solely empirical until a proper theory is available.

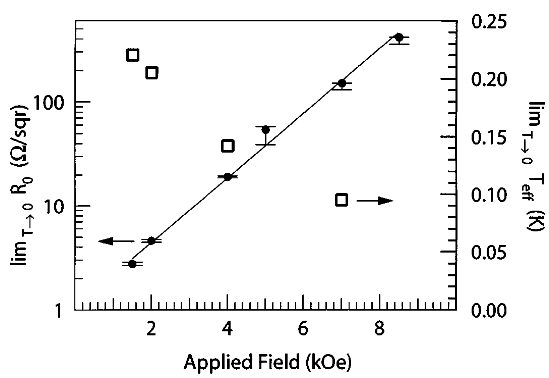

FIG. 4. Low-temperature saturation values of the resistance (solid circles) and of the effective temperature (open squares). The solid line is the best fit by the empirical form $R_0^{QL} = A\exp(H/H')$. The two lowest circles represent the average of the minimum resistance (lower bound) and the resistance at the lowest measured temperature (upper bound). All the other points represent the lowest measured values. The remaining lower bounds are extrapolations obtained by assuming that the shape of the 1.5 kOe effective temperature curve describes each of the curves in turn with the offset that best aligns them.

Finally, we point out that the simplest quantum tunneling description of our data—developed below—is not self-consistent; a more sophisticated approach is required. In the usual first cut approach to calculating the quantum tunneling rate of a simple particle through a simple barrier, the action is written $S_E = U\tau$, where τ is the bounce time between the turning points of the inverted potential [5]. If we model the dislocation nucleation process in this way and assume that the nature of the barrier and its height U remain the same as observed in the thermally activated regime, then we can associate $\tau = \hbar/k_B T_{\mathrm{eff}}^{QL}$. The open squares corresponding to the right coordinate axis in Fig. 4 then imply, following this reasoning, that τ increases monotonically with increasing magnetic field. However, we expect the opposite dependence of τ on H, since the size of dislocations—and thus the width of the barrier and the effective tunneling mass—and also the barrier height all *decrease* with increasing field. (The tunneling rate, proportional to R_0^{QL}, behaves in the intuitive way. It is the crossover temperature that does not.) Following this line of reasoning leads to the conclusion that a simple model that treats the dislocation nucleation event as a point particle tunneling through a simple fixed barrier cannot self-consistently explain our data. The nonintuitive field dependence may arise from the complexities of the intermediate vortex lattice configurations during the quantum nucleation event. Another possibility is that the level of dissipation within the pair nucleation process depends in an unexpected way on the density of normal electrons in the system, and consequently on H. Alternatively, the field dependence may be a clue that a novel quantum process is present. We leave this point as an open question.

In conclusion, we have observed behavior strongly suggestive of quantum mechanical tunneling of vortices in a highly disordered thin film superconductor—most likely in the form of quantum nucleation of unbound vortex lattice dislocation-antidislocation pairs. Independent of our interpretation, the natural extrapolation of our observation to $T = 0$ implies that no true superconducting state exists in 2D for $H > 0$. This conclusion contradicts the prediction of a transition to a superconducting vortex glass state at $T = 0$ [3] and would seem to imply that the superconducting-insulating transition is actually a metal-insulator transition. Somewhat paradoxically, we have previously reported [10] and subsequently confirmed that the S-I scaling theory [2] describes our system extremely well near the S-I critical field H_c in the finite temperature critical regime currently accessible to us experimentally.

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