expected to be of order unity. Also, the SNST could be a quantum phase transition with associated critical point amplitude fluctuation effects. Recent theories (see, for example [7, 8]) predict the growth of fluctuations in the superconducting order parameter amplitude that diverge at the transition. The conductance of a large area tunnel junction measures a density of states averaged over these fluctuations. It should be broader than the density of states for a BCS superconductor with a single gap. The resistively measured transition would be sensitive to the accompanying spatial distribution of T_{c0}'s in the film and would also become broader. In two dimensions, the predicted divergence for a superconductor-to-normal metal transition is logarithmic [8], in accord with the result shown in Fig. 5b.

The above results and discussion suggest a qualitatively different picture of the SNST for homogeneous films than the 'standard' dirty boson models of the SIT (see, for example [9]). In the latter, increasing the normal state sheet resistance decreases the Josephson coupling between individually superconducting islands in the film. The concomitant reduction in the superfluid density or phase modulus leads to the growth of phase fluctuations between the islands. If the phase fluctuations are large enough, all phase coherence is lost, Cooper pairs become localized to their individual islands, and the system becomes an insulator [9, 10]. In contrast, in the homogeneous films considered here, the growing amplitude fluctuations increase the quasi-particle density. This increase comes at the expense of the Cooper pair or superfluid density. At high enough sheet resistances the superfluid density is driven to zero and a nonsuperconducting phase consisting of weakly localized quasi-particles takes over.

5. Conclusions

We have presented electron tunneling and transport measurements on homogeneous films that provide evidence that fluctuations in the superconducting order parameter amplitude grow as the amplitude collapses at the SNST. These fluctuation effects dominate the experimentally accessible portion of the resistive transitions of the films closest to the SNST. Theories suggest that the mesoscopic fluctuation, quasi-particle lifetime, and quantum critical fluctuation effects could each contribute to the growth of the amplitude fluctuations. Taken as a whole, our results suggest that amplitude fluctuations become so strong near $R_n \simeq R_Q$ that they drive the superfluid density to zero to destroy the superconducting state.

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Vortex states at low temperature in disordered thin and thick films of a-Mo_xSi_{1-x}

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Abstract. We have measured the ac complex resistivity in the linear regime, as well as dc resistivity, for thick (100, 300 nm) amorphous (a-)Mo_xSi_{1-x} films at low temperatures ($T > 0.04$ K) in constant fields B. The critical behavior associated with the second-order transition has been observed for both dc and ac resistivities, which is similar to that observed for *granular* indium films. This is the first convincing evidence for the vortex glass transition (VGT) in the *homogeneously* disordered low-T_C superconductors containing microscopic pinning centers. We have found that the VGT persists down to $T \sim 0.1 T_{C0}$ up to $B \sim 0.9 B_{C2}(0)$, where T_{C0} and $B_{C2}(0)$ are the mean-field transition temperature and upper critical field at $T = 0$, respectively. At $T \to 0$ the VGT line $B_g(T)$ extrapolates to a field below $B_{C2}(0)$, indicative of the presence of a $T = 0$ quantum-vortex-liquid phase in the region $B_g(0) < B < B_{C2}(0)$.

For thin (4 nm) films the ($T = 0$) field-driven superconductor–insulator transition takes place at B_C. We have not obtained evidence of the metallic quantum liquid phase below B_C, while in $B > B_C$ an anomalous negative magnetoresistance (MR) suggesting the presence of the localized Cooper pairs has been observed. The negative MR is commonly observed for thin films; however, for thick films the MR is always positive. This means that the two-dimensionality plays an important role in the appearance of the negative MR (or localized Cooper pairs). The negative MR is no longer visible as the field is applied parallel to the film surface, consistent with the view that mobile vortices, as well as localized Cooper pairs, are present in $B > B_C$.

1. Introduction

The superconductor–insulator transition (SIT) in two dimensions (2D) has been thoroughly studied during recent years [1–16]. While most of the studies have focused on the electronic states near the SIT, the vortex states at low temperature have not yet been well clarified. In moderately

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disordered superconductors, the vortex solid state at low temperature is a vortex glass [17]. Experimental evidence for the vortex glass transition has been obtained in several 3D high-T_C superconductors (HTSC's). However, it was not evident until recently that the VGT also exists in low-T_C superconductors (LTSC's) [18, 19]. For 2D superconductors it is more difficult to verify the VGT, since the VGT in 2D is considered to be a quantum phase transition which occurs at zero temperature ($T = 0$). According to Fisher, it is a transition from the VG to Bose glass (BG) phase [20].

For LTSC's compelling evidence for the VGT has been first reported in the ac complex resistivity of thick granular In films with mesoscopic ($\sim$ 10 nm) grain sizes [19]. In the meanwhile, in *homogeneously* disordered thick amorphous (a-)Mo_3Si films, whose transition temperature and critical field are of comparable order to those for the In film, absence of the critical dynamics of the VGT has been reported from the ac resistivity measurements [21]. Therefore, it has been suggested that the structure of the pinning centers may play a crucial role in the appearance of the VGT, even though they are of mesoscopic sizes. Theoretically, influences of disorder on the VGT have not yet been fully clarified.

Another interesting problem to be explored is whether the VGT is actually observed down to low enough temperatures. This problem has not yet been studied experimentally, since most of the experiments have been performed using HTSC's whose upper critical field B_{C2} at low temperature is extremely high. Although several theories [22, 23] treat quantum melting of the vortex lattice in clean systems, it has not been studied how quantum fluctuations affect the VG state in disordered systems.

In this paper, we present convincing evidence for the VGT in uniformly disordered thick a-Mo_xSi_{1-x} films [24], which is contrary to the previous report [21]. The VGT persists down to low enough temperatures $T \ll T_{C0}$ and up to high fields near $B_{C2}(0)$, where T_{C0} and $B_{C2}(0)$ are the mean-field transition temperature and upper critical field at $T = 0$, respectively. At $T \to 0$ the VGT line $B_g(T)$ extrapolates to a value below $B_{C2}(0)$, indicative of the presence of a $T = 0$ quantum liquid phase in the regime $B_g(0) < B < B_{C2}(0)$. We also present the measurements of the dc resistance at low T for a series of thin (2D) films with various disorder. Based on the data, we construct the $T = 0$ phase diagram for the field-driven and zero-field SIT's in 2D. Possible vortex states at $T = 0$ will be discussed in comparison with those in thick films.

2. Experimental

All of the films used in this study were prepared by co-evaporation of pure Mo and Si in vacuum better than 10^{-8} Torr [11, 25]. The structure of the film was confirmed to be amorphous by means of transmission electron microscopy (TEM). It is known that Mo_xSi_{1-x} films are amorphous as far as $x < 0.75$, but crystallization of Mo is expected at $x > 0.75$ [26]. The values of x of our films stay in the range 30 – 61 at.%, which are well below 75 at.%. The resistance was measured by four-terminal method. The ac transport data, the amplitude ρ_{ac} and phase φ of the ac resistivity, were taken in the linear regime as a function of the temperature and frequency f [19, 24]. The ac voltages induced across the sample were measured using the LCR meter after being enhanced with a low-noise preamplifier. We evaluated a frequency-dependent gain and/or phase delay of the amplifier by measuring the standard noninductive load resistor which was connected in place of the sample. Thus we obtained the frequency-dependent resistivity in the frequency range $f \sim 200$ kHz – 6 MHz from the measured voltages.

3. Results

There is a large difference between the transport properties of the thick films and thin films in the presence of the field. Upon cooling, the resistivity ρ for the thick films goes to zero at finite nonzero temperature $T_C(B)$. In contrast, for the thin films Arrhenius type of the resistance R is observed, indicating that $R(T)$ goes to zero only at $T = 0$ for any nonzero field studied. These results are consistent with the prediction of the VG theory [17, 20] that the VG phase is present at finite T in 3D but only at $T = 0$ in 2D. In order to study the difference in the vortex dynamics between 3D and 2D vortex systems, we have measured the current-induced voltage noise spectral density S_V [27]. The results show that the vortex dynamics probed by S_V strikingly depends on thickness of the film [25]. In Figure 1 we plot the field dependence of S_V at $f = 900$ Hz for the thick (100 nm) and thin (6 nm) films measured at constant voltage V and $T = 1.8$ K ($< T_{C0}$). In the case of the thick film, a clear peak in S_V is observed over the broad field region lower than B_g, which originates from a plastic-flow motion of the vortex glass. This result is contrasted with that for clean 2H-$NbSe_2$ [28] where noise is localized in a narrow range of B prior to first-order melting B_m. This difference is due to the different pinning strengths between two systems. In our system the concentration of pinning is high enough for the highly disordered plastic flow to be observed even at fields far below B_g. The field dependence of S_V for the thin film is markedly different from that for the thick film. By applying a small field, S_V decreases abruptly below the background level

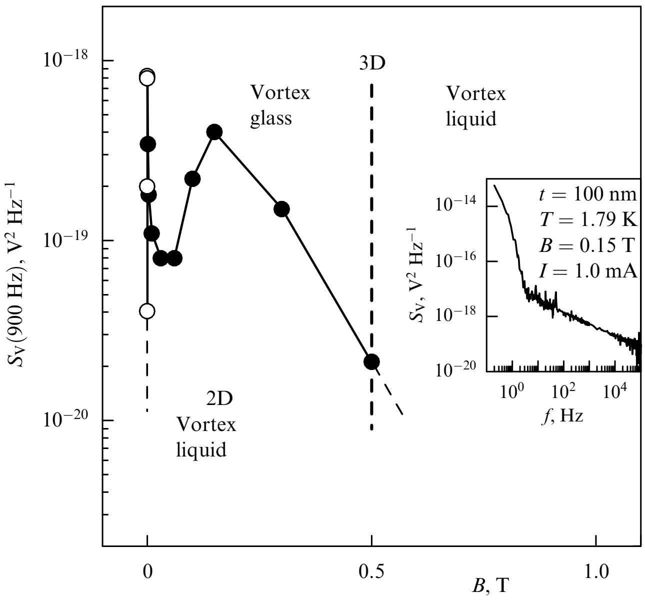

Figure 1. S_V vs. B for the thick film (•) and the thin film (○) at $T = 1.8$ K. The background level is $\sim 10^{-20}$ V^2 Hz^{-1}. For the thick film a broad peak, as well as $1/f^{\beta}$ ($\beta < 0.6$)-like noise spectra, originating from a plastic-flow motion of the vortex glass is visible in fields lower than $B_g \approx 0.5$ T. Inset: $S_V(f)$ for the thick film at $B = 0.15$ T [25].

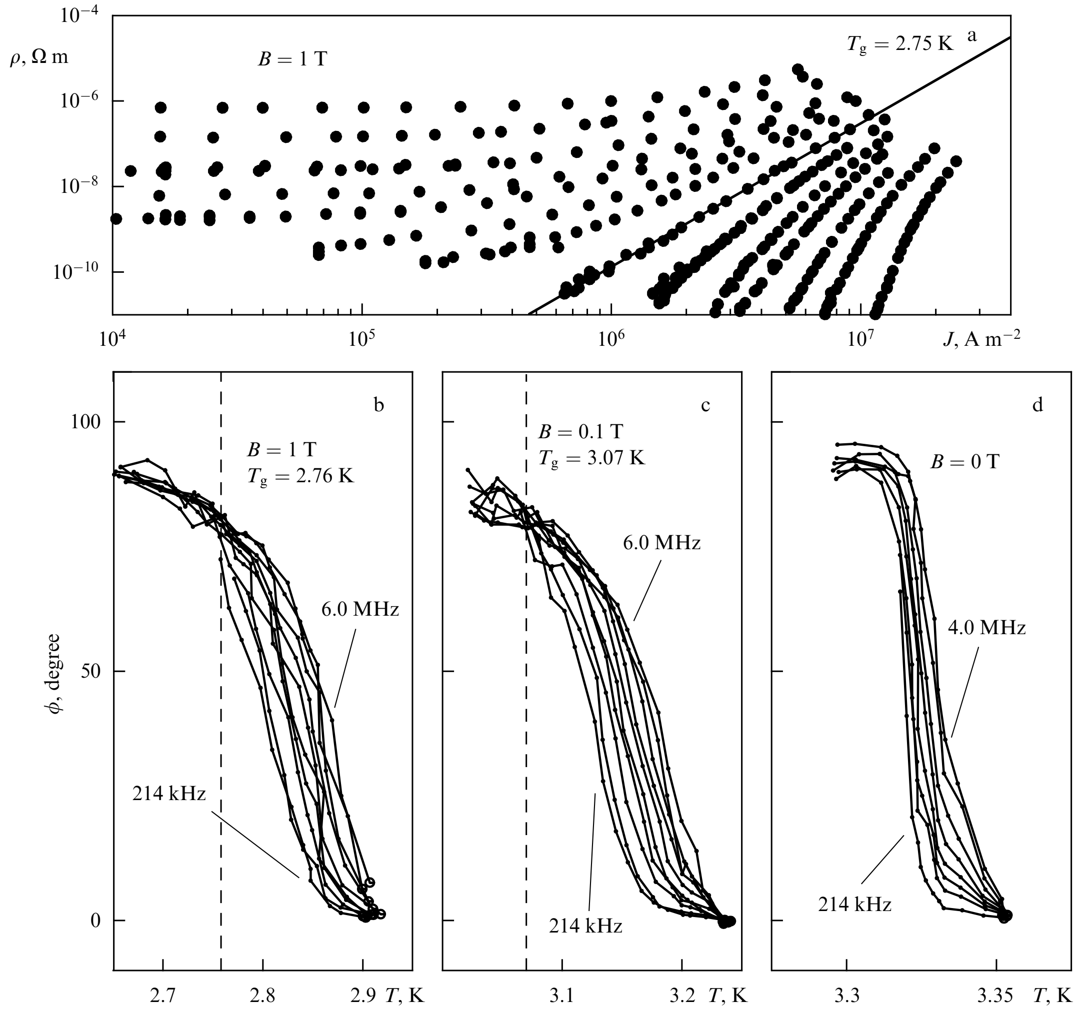

Figure 2. DC $J-\rho$ isotherms in $B = 1$T for the 300-nm thick film. With decreasing T, the isotherms (2.31 – 2.98 K) shift to the right. The straight line indicates the VGT ($T_g \approx 2.75$ K) (a). Temperature dependence of the phase $\phi(T,f)$ of the ac resistivity for different f in $B = 1$ T (b), $B = 0.1$ T (c), and $B = 0$ (d).

without showing a peak. This result is again consistent with the picture of the VG theory: in 2D there is no vortex-glass state at $T > 0$ and hence, there is no plastic-flow motion which yields large noise.

We first present the measurements of dc and ac complex resistivities for the thick (3D) films. Shown in Fig. 2a are the dc current density J vs resistivity ρ isotherms in $B = 1$ T for the 300-nm thick film with $T_{C0} = 3.35$ K and $B_{C2}(0) \approx 7.9$ T. With decreasing T, curvature changes from positive to negative at a certain temperature, suggesting that the VGT occurs at this temperature T_g. From the slope of the critical isotherm, the dynamical exponent z is obtained to be $z \approx 7-8$ [24]. Figures 2b, c, d depict the temperature dependence of the phase φ of the ac resistivity for different f in $B = 1, 0.1$, and 0 T, respectively. In $B = 1$ and 0.1 T we find a particular point $[T_g(B)]$ at which $\varphi(T,f)$ at different f merges to the same value (φ_g) lower than 90°. However, such a point is not visible in $B = 0$, indicating that the peculiar behavior of φ shown in Figs 2b, c indeed originates from field-induced vortices. The phase transition temperature T_g and critical exponents (z, ν) are consistently determined from both dc and ac resistivities [24]. These results provide convincing evidence for the VGT in *homogeneously* disordered LTSC's containing microscopic pinning centers.

Critical behavior associated with the VGT is also observed for the 100-nm thick film with $T_{C0} = 2.4$ K and $B_{C2}(0) \approx 5.7$ T. Using this film, we have explored the existence of VGT in the low-T high-B region. In fields B lower than 5.3 T, the dc resistivity ρ goes to zero at finite $T_g(B)$, where $\rho(T)$ follows a power-law formula

$$\rho(T) \sim \left(\frac{T}{T_g} - 1\right)^{\nu(z-1)}$$

predicted by the VG theory (Fig. 3). In contrast, in fields higher than $B_0 = 5.31$ T, $\rho(T)$ shows upward curvature at low T and tends to the finite resistivity at $T \to 0$, as shown in the inset of Fig. 3. Thus, B_0 is a critical field for the superconductor – metal transition at $T = 0$. We notice that B_0 almost coincides with the field above which crossing behavior of the phase disappears. In the limit of zero temperature, the VGT line $B_g(T)$ extrapolates to a field close to $B_0 = 5.31$ T, that is clearly lower than $B_{C2}(0) \approx 5.7$ T. Here, $B_{C2}(T)$ is defined as a field at which ρ decreases to 90% of the normal-state resistivity ρ_n and $B_{C2}(0)$ is determined by a smooth extrapolation of $B_{C2}(T)$ to $T = 0$. The implication of this result is that the $(T = 0)$ metallic quantum-vortex-liquid (QVL) phase is present in the field region $B_0 < B < B_{C2}(0)$.

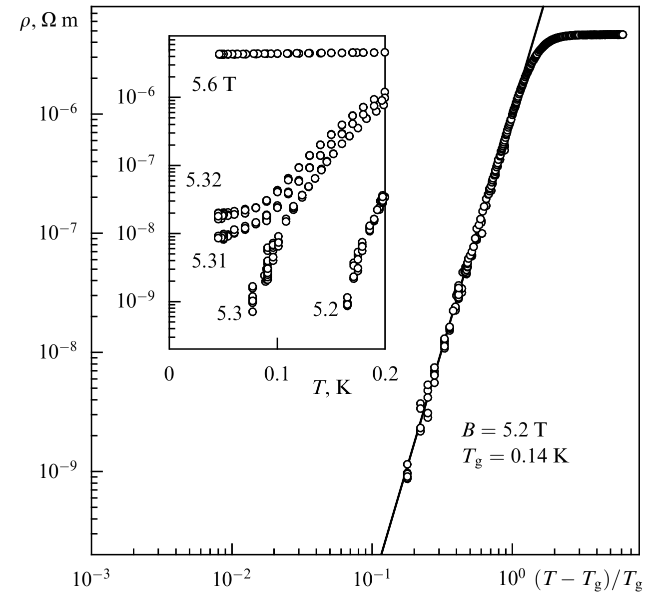

Figure 3. DC resistivity ρ of the thick (100 nm) film in $B = 5.2$T vs $(T - T_g)/T_g$ in a log – log plot. The full line represents the fit of the data to the functional form predicted by the VG theory. Inset: temperature dependence of ρ in different B.

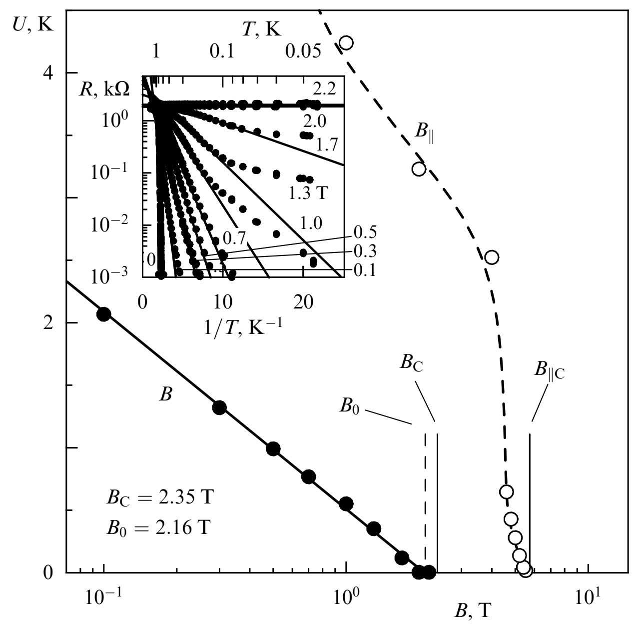

Figure 4. Activation energies U (•) of the thin (4 nm) film extracted from the Arrhenius plot of $R(T)$ in perpendicular B plotted against log B. Slope (○) of the Arrhenius plot of $R(T)$ in parallel $B_\parallel$ is also plotted. Inset: Arrhenius plot of $R(T)$ in various perpendicular B lower than B_C [29].

Existence of the metallic QVL phase has been reported in several thin-film (2D) superconductors [8, 9]. It has been claimed based on the fact that with increasing B, the activation energy $U(B)$ extracted from the activated resistance $R(T)$ decreases and extrapolates to zero at a field B_0 substantially lower than the critical field B_C. Thus, B_0 is a critical field for the $T = 0$ superconductor – metal transition in 2D. Furthermore, for some systems the apparent temperature-independent resistance, so-called flat tail, is observed at low T and it is attributed to quantum tunneling of vortices [8, 10]. These results are important because the picture of the 2D SIT might be questioned.

Let us focus on the results for the thin (4 nm) films [12, 29]. An Arrhenius type of the resistance is commonly observed in finite (perpendicular) fields B below the critical field B_C for the field-driven SIT. Shown in the inset of Fig. 4 is an example for the relatively disordered film with $B_C = 2.35$ T, which lies in the vicinity of the zero-field SIT. In $B < 1$ T the activated resistance is observed down to $R \to 0$. However, in $B > 1$ T the slope of the straight line (i.e., activation energy U) shows a discontinuous drop below about 0.1 K. Such a drop in U is no longer visible for the less resistive films, suggesting that disorder may play a role in the reduction of U at very low temperatures. The origin of it is not clear at present. However, we can say that it is not related to flux motion, because the decrease in the slope is also seen for parallel fields $B_\parallel$. We suggest that in order to obtain further proof for the metallic QVL, it is important to perform the transport measurements in parallel fields for other superconductors in which the flattening of the resistance is observed in perpendicular fields.

In Figure 4 we plot the slope of the $\log R$ vs. $1/T$ plots as a function of the field both for perpendicular and parallel orientations. The slope (i.e., U) for perpendicular fields B follows the $\log B$ formula predicted by the dislocation model [30], whereas for parallel fields $B_\parallel$ it never obeys this formula, indicating that the resistance in the presence of perpendicular B is indeed dominated by flux motion. A characteristic field B_0 at which U extrapolates to zero is slightly lower than (or very close to) B_C. Thus, in our 2D system the metallic QVL phase below B_C is not so distinct, most likely absent [29].

We next consider the magnetoresistance (MR) for the thin films at low temperatures [12]. For insulating films the MR is always monotonic and positive, and reproduced by a $\log B$ formula at fields higher than ~1 T. This is explained by the 2D weak localization theory for fermions in the presence of the strong spin-orbit scattering, which originates from Mo atoms. On the other hand, for superconducting films the anomalous peak and subsequent decrease in MR are observed at low temperature in fields higher than B_C. Since superconducting films are in the weakly localized regime for fermions and strong spin-orbit scattering is present due to higher Mo concentration than in insulating films, the MR originating from unpaired electrons should be positive. Therefore, the origin of the decrease in MR is not attributed to delocalization effects of unpaired electrons with increasing B. It is sought for the localized Cooper pairs, which are present even on the insulating side of the field-driven SIT [12].

For parallel fields $B_\parallel$, the critical field $B_{\parallel C}$ is much larger than B_C for perpendicular fields (e.g., $B_{\parallel C} = 5.7$ T and $B_C = 2.35$ T), where $B_{\parallel C}$ is defined as a field at which R at low temperature is temperature-independent. The most striking difference between the perpendicular and parallel MR is that the peak is not visible with parallel fields, as shown in the inset of Fig. 5. This result looks consistent with the view that mobile vortices, as well as the localized Cooper pairs, are present in the insulating region where $R(B)$ shows a decrease. We have also found that for thicker (100, 300 nm) films the anomalous peak in the MR is not observable down to the lowest temperatures irrespective of disorder, indicating that the 2D (two-dimensionality) plays an important role in the appearance of the anomalous peak in $R(B)$ [29]. This result is again consistent with the 2D VG theory by Fisher.

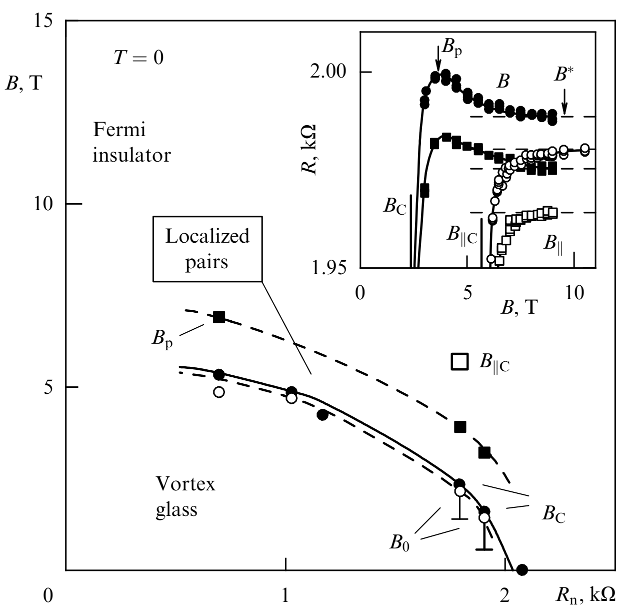

Figure 5. $B - R_n$ phase diagram for the zero-field and field-driven SIT's in 2D at $T = 0$. According to the Fisher's model, B_C is interpreted as a phase boundary separating the VG and BG phases: B_{C2} around which the localized Cooper pairs (and mobile vortices) disappear is close to B_p (or $B_{\parallel C}$) where the peak in $R(B)$ occurs. Inset: high R regions of MR for the thin (4 nm) film at $T = 0.05$ K (circles) and 0.1 K (squares) with perpendicular (solid symbols) and parallel (open symbols) orientations.

4. Discussion

Based on the data, we construct the field-disorder ($B - R_n$) phase diagram for the zero-field and field-driven SIT's in 2D at $T = 0$ [29] in Fig. 5. One can see from the figure that the metallic QVL regime ($B_0 < B < B_C$) is very narrow for any R_n, while there is the broad anomalous insulating regime suggesting the presence of the localized Cooper pairs just above B_C. It is difficult to determine experimentally the upper boundary of this regime (or B_{C2}), around which the localized Cooper pairs disappear. This field is nearly identified with the $T = 0$ characteristic field B^* above which the decrease in $R(B)$ ceases. From the temperature dependence of the peak in $R(B)$, we have known that $B^*(T)$ in the limit of $T = 0$ is close to B_p (or $B_{\parallel C}$) where the peak in $R(B)$ occurs [12].

According to the theory by Fisher [20], at B_C there is a Bose-glass (BG) insulator with localized Cooper pairs and mobile vortices. Either at B_p or $B_{\parallel C}(\sim B_{C2})$, the Bose insulator goes to a Fermi insulator. Thus there are the finite amplitude of the order parameter and strong phase fluctuations in the field region $B_C < B < B_{C2}$. On the other hand, with parallel-field orientation, a field $B_{\parallel}$ couples to (only) the order parameter amplitude but not affects phase fluctuations. Above $B_{\parallel C}$ there is no order parameter amplitude and hence no Bose insulator. Thus $R(B_{\parallel})$ is a monotonic function of field. All of the data seem to be consistently explained with this scenario [29]. We consider, however, that to prove the existence of the BG phase more convincingly, direct detection of flux motion is necessary.

We discuss here alternative interpretations of the data. The peak in the MR has been also reported for granular films [1, 3]. This is explained as follows: as the field B increases up to a certain field B_C, the superconducting clusters are uncoupled and global superconductivity disappears, where the transport is dominated by single particle excitations. The resistance is high because of the energy gap Δ in the excitation spectrum of the quasi-particles in the clusters. At some field that is higher, the superconductivity (Δ) within the clusters is destroyed, resulting in a decrease in the resistance. On the other hand, with the parallel-field orientation, decoupling of superconducting clusters does not occur as far as the film is thin and the onset of the resistance coincides with the complete destruction of the superconductivity at $B_{\parallel C}$. This picture seems to be in accordance with the observations of MR presented here. We note, however, that for granular films the peak is clearly visible even in insulating films [1, 3], which is contrasted with the present results. We do not have any experimental proof suggesting the presence of superconducting clusters in our system, while we cannot completely exclude the possibility of the presence of the small superconducting clusters which are not detected within our resolutions of TEM. The interpretation mentioned above is associated with the morphology of the films. On the other hand, there are some models [31, 32] which have addressed the formation of the superconducting clusters ('physical clusters') near the SIT even in a homogeneous system. Using these models, the similar interpretation could be derived from intrinsic effects.

The anomalous peak in the MR has been also reported in thin superconducting films of a-InO_x [2]. Gantmakher et al. [13, 15] have found on the basis of the studies on 20-nm thick a-InO_x films that the anomalous peak in the MR appears not only for perpendicular fields but also for parallel fields. This result indicates that the origin of the MR peak is not related to flux motion nor the particular model by Fisher. As explained by the authors, it is attributed to dissociation of electron pairs with a distribution of the energy gap Δ. They have also found that the peak in the MR occurs even for the insulating films [16]. Thus, it is interesting to ask whether the InO_x system contains actual superconducting clusters or physical ones.

Also, the peak in the MR has been observed even for the thick (170 nm) a-Mo_3Si film with $T_{C0} \sim 8$ K [33]. This result implies that 2D does not play a role in the appearance of the MR peak, which is in contrast with our results. It is questionable for us, however, whether the Mo_3Si films are actually amorphous, because concentration of Mo is close to the threshold above which formation of superconducting Mo clusters might be expected [26]. The different critical dynamics of the 3D VGT observed between Mo_3Si and our films may be attributed to different film properties [24].

Finally, we compare the possible vortex states at $T = 0$ for the thin films with those for the thick films. For simplicity we compare the results for two films whose $B_{C2}(0)$'s (~ 5.7 T) are almost identical to each other. For the thick (100 nm) film the VG phase persists up to $B_0 = 5.3$ T, that is close to $B_{C2}(0)$. Above B_0 there is the *metallic* QVL phase. On the other hand, for the thin (4 nm) film the upper boundary of the VG phase is $B_C = 2.35$ T, that is far below $B_{C2}(0)$. The metallic QVL phase below B_C is not evident ($B_C \approx B_0$), most likely absent, while there is the unusual insulating regime $B_C < B < B_{C2}(0)$ suggesting the presence of the localized Cooper pairs. Within the Fisher's picture, this regime is interpreted as the *insulating* QVL phase. We do not know appropriate theories which take account of the effects of quantum fluctuations on the VGT or which deal with the quantum liquid of vortices in the disordered systems in which the VGT is observed. Therefore, we attempt to analyze our results in terms of the ($T = 0$) first-order melting theory of Blatter et al., which takes account of quantum fluctuations [22]. Here, we assume that a $T = 0$ melting field B_m in the theory corresponds to $B_g(T \to 0)$ in the present system: i.e., $B_m \approx B_0$ and $B_m \approx B_C(\approx B_0)$ for the thick and thin films, respectively. Using the experimental values of $B_{C2}(0) \approx 5.7$ T and ρ_n (or R_n), we heuristically evaluate the Lindemann numbers c_L for the melting transition in the theory. There is no theoretical justification for making such an analysis, however, interestingly, extracted values of c_L for both films stay around 0.1 – 0.2, which are of reasonable magnitude.

In summary, we present convincing evidence for the VGT in thick a-Mo_xSi_{1-x} films, which persists down to low enough temperatures $T \ll T_{C0}$ and up to high fields near $B_{C2}(0)$. At $T \to 0$ the VGT line $B_g(T)$ extrapolates to a value below $B_{C2}(0)$, indicative of the presence of the metallic quantum liquid in the regime $B_g(0) < B < B_{C2}(0)$. For thin films we have not obtained evidence for the metallic quantum liquid phase, while we have observed an anomalous peak in the magnetoresistance (MR) suggesting the presence of the localized Cooper pairs on the insulating side of the field-driven SIT. The peak in MR is no longer visible with parallel fields or for thicker films. All these data seem to support the picture of the 2D VG – BG transition. To demonstrate the existence of the BG phase more convincingly, however, further experiments including direct detection of flux motion are necessary.

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Metallic single-electron transistor without traditional tunnel barriers

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Abstract. We report a new type of single-electron transistor (SET) comprising two highly resistive Cr thin-film strips (~ 1 μm long) connecting a 1 μm-long Al island to two Al outer electrodes. These resistors replace small-area oxide tunnel junctions of traditional SETs. Our transistor with a total asymptotic resistance of 110 kΩ showed a very sharp Coulomb blockade and reproducible, deep and strictly e-periodic gate modulation in wide ranges of bias currents I and gate voltages V_g. In the Coulomb blockade region ($|V| \leq$ about 0.5 mV), we observed a strong suppression of the co-tunneling current allowing appreciable modulation curves $V(V_g)$ to be measured at currents I as low as 100 fA. The noise figure of our SET was found to be similar to that of typical $Al/AlO_x/Al$ single-electron transistors, viz. $\delta Q \approx 5 \times 10^{-4} e/\sqrt{\mathrm{Hz}}$ at 10 Hz.

1. Introduction

The SET is a system of two ultra-small metal – insulator – metal tunnel junctions attached to a small island which is capacitively coupled to a gate electrode. Due to their considerable resistance, $R \gg R_Q \equiv h/4e^2 \cong 6.5$ kΩ, the tunnel junctions ensure charge quantization on the island. On the other hand, the junctions still enable the (correlated) charging and discharging of the island by individual electrons when the temperature is sufficiently low, $k_B T \ll E_c$. Here $E_c = e^2/2C_\Sigma$ is the charging energy and $C_\Sigma = C_1 + C_2 + C_0 + C_g$ is the total capacitance of the island, which includes the capacitances of the junctions, $C_{1,2}$, the self-capacitance of the island, C_0, and the capacitance between the island and gate electrode, C_g. Transport of electrons is controlled by the transistor gate polarizing the island and therefore changing the Coulomb blockade threshold. Increase in the gate voltage V_g causes a stepping increment of the number of electrons on the island and this leads to e-periodic dependence of the $I – V$ characteristic on V_g. Due to this effect, the transistor provides means for measuring the polarization charge on its island with sub-electron accuracy. This property of SET was successfully exploited in many experiments to measure and monitor sub-electron quantities of charge in mesoscopic systems (see some examples in Refs [1 – 5]). Different materials and methods have been used for the fabrication of

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