Double reentrant superconductor-insulator transition in thin TiN films

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We studied the superconductor-insulator transition in very thin films of TiN (and NbN) at very low temperature ($T=35$ mK) as a function of sheet resistance and magnetic field. The critical temperature dependence versus sheet resistance follows the Finkel'stein theory for homogeneous films. A TiN sample very close to the critical point exhibits a double reentrant behavior below $T=150$ mK as the magnetic field increases. At large magnetic field and very low temperature a negative magnetoresistance is observed.

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I. INTRODUCTION

The nature of the superconductor-insulator transition (SIT) in thin disordered films of superconductors is a long debated topic. The question is the suppression of the superconductivity by the Coulomb interaction in the presence of disorder. Near the superconducting transition the conductivity has two components, one coming from fluctuating Cooper pairs (boson contribution) and the other from the conductivity of quasiparticles. One can neglect the quasiparticle contribution in artificial networks of Josephson junctions. Due to the fixed charge on each island by Coulomb blockade and the charge-phase duality, the phase of Cooper pairs strongly fluctuates from island to island and the critical temperature decreases down to zero.[1] This reflects the competition between Coulomb interaction and Josephson coupling in the context of pure granular systems. This suppression of the critical temperature is referred to as the boson scenario in the literature.

But the Coulomb interaction could also localize quasiparticles in thin homogeneously disordered films due to vanishing dynamical screening as the disorder increases. In homogeneous superconducting films, the effects of Coulomb interaction and superconducting attraction are entangled:[2] the critical superconducting temperature is a function of the normal sheet resistance of the film, through the parameter $\gamma=\ln(\hbar/k_BT_{c0}\tau)$, where τ is the elastic mean free time and T_{c0} the bulk BCS critical temperature. For values of γ larger than 5 (large disorder, small bulk T_{c0}), the boson contribution is unimportant.[3] In granular films, depending on parameters such as the grain size, diffusion constant, and intergrain transfer probability, the boson contribution could be restored.[3] A full theoretical analysis of the competition between Coulomb repulsion and superconducting attraction embracing the homogeneous diffusive case and the pure granular case is still lacking, despite many recent developments.[3–7]

From an experimental point of view, superconducting films, which are partly homogeneously disordered and partly granular, should display interesting properties. Few generic resistance versus temperature behaviors have been identified in the literature: in films with strong granularity, as one increases the normal-state resistance (usually by decreasing film thickness), the anomaly of the resistance at the bulk critical superconducting temperature is unaffected, and the low-temperature resistance first drops and then develops an upturn, exhibiting a type of reentrant behavior. On the contrary, in more homogeneous films, the critical temperature is a monotonous decreasing function of the sheet resistance and no reentrant behavior is seen. In addition, at low temperature the resistance can saturate in some films; this has been attributed to dissipative coupling of vortices with the environment.[8] An enormous number of experimental reports indicate that interpretation of the SIT in terms of either bosons or a quasiparticle mechanism cannot be the full story (see, for instance, Refs. 9–12). According to most authors, more experimental work is needed to clarify the problem.

In this work we studied thin superconducting films of TiN and NbN by resistance measurements down to very low temperature and with high magnetic fields; for the sample which is the closest to SIT's in zero magnetic field, we observed a double reentrant behavior when a magnetic field is applied to drive the SIT. This observation is similar to the behavior reported in arrays of Josephson junctions, a canonical case for phase-only driven SIT's (Ref. 1) but has not been reported previously in superconducting films. For samples farther from the critical point, a more standard behavior is obtained for the resistance versus temperature: either a saturation or a single reentrant behavior at low temperature. The double reentrant behavior suggests either the existence of multiscale structural heterogeneities or the existence of a double transition near the critical point. The similarity with artificial systems supports the second hypothesis.

II. SAMPLE FABRICATION AND CHARACTERIZATION

TiN and NbN films has been deposited by reactive dc-magnetron sputtering (argon and nitrogen plasma) from pure targets of Ti and Nb.[13,14] We have used Si/SiO_2 substrates at 350 °C for TiN deposition, and *R*-plane Al_2O_3 substrates at 600 °C for NbN. X-ray diffraction measurements have confirmed that the films are microcrystallized in the cfc cubic phase with lattice parameter of ~0.424 nm for TiN and in epitaxially grown (100) orientation for NbN.

Cubic nitrides of transition metals are interstitial compounds that are stable on a large range of compositions. The TiN_x atomic ratio has been found in the range $1\leqslant x\leqslant 1.3$.[15,16] NbN_x is more likely understoichiometric in nitrogen $0.8\leqslant x\leqslant 1$. TiN films sputtered at higher nitrogen pressures have more defects (interstitials and voids) due to

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higher relative amounts of nitrogen; such films are much more resistive: for TiN, 100 nm films have resistivities at 10 K ranging from 80–1100 $\mu\Omega$ cm when the total sputtering pressure is varied from 1 to 6 mTorr by increasing the nitrogen gas flow rate in the deposition chamber.

TiN and NbN exhibit the typical electronic properties of dirty metals: low relative resistive ratio (RRR), relatively high residual resistivities, and short electronic mean free paths (ranging from 20 to 0.7 nm for TiN films). A type-2 superconductivity is observed with rather large penetration lengths around 250 nm (measured by single-coil mutual inductance technique[13]) and very short coherence lengths (values computed from the dependance of H_{c2} with temperature are 6.5 nm for TiN and 5 nm for NbN).[17] The bulk superconducting transition temperature T_{c_0} is around 4.7 K for TiN and 16 K for NbN; T_{c_0} is rather insensitive to disorder for thick films: for 100 nm TiN films, the critical temperature only decreases from 4.7 to 4.4 K when the resistivity changes by one order of magnitude from 80–1100 $\mu\Omega$ cm.

Films are prepared for electronic transport measurements by sputtering titanium-gold contacts on diced samples of area about 0.2 cm^2. TiN and NbN films are found to be very stable in air at room temperature, showing only a slight aging due to residual oxidation; the surface layer of native oxide is reduced compared to simple metals as niobium.

The four-probe resistance measurements are performed in the I-V linear regime even at very low temperature. Samples are connected by 2.5 m of Phillips thermocoax lossy coaxial cables for rf filtering. Below 1 K, an homemade dilution refrigerator with a plastic mixing chamber is used in conjonction with a 16 T Oxford magnet. An RuO_2 sensor based thermometry with proper calibration in magnetic field is used. All our samples have been measured using the same setup and experimental procedure, such that a direct comparison of their low-temperature behavior makes sense.

III. RESISTIVITY DEPENDENCE OF THE CRITICAL TEMPERATURE

The set of resistance curves obtained for various thicknesses is characteristic of homogeneous disordered films, with a continuous decrease of the critical temperature without significant broadening of the transition and no reentrance of the resistance at low temperature (see Fig. 1, at zero magnetic field).

Very thin films are found to be insulating (see Fig. 1); TiN films in this insulating regime show resistance obeying a simple Arrhenius activation law at low temperature, with an activation temperature of 770 mK for the sample closest to the SIT. This could be due to Coulomb blockade on grains with Coulomb charging energy $E_c \simeq k_B \times 0.77$ K (see Fig. 2).

The continuous decrease of the critical temperature with the increase of resistance is well fitted within the Finkel'stein model,[2] which applies for homogeneously disordered superconducting films:

$$\frac{T_c}{T_{c0}} = e^{\gamma}\left(\frac{1/\gamma - \sqrt{t/2} + t/4}{1/\gamma + \sqrt{t/2} + t/4}\right)^{1/\sqrt{2t}} \quad (1)$$

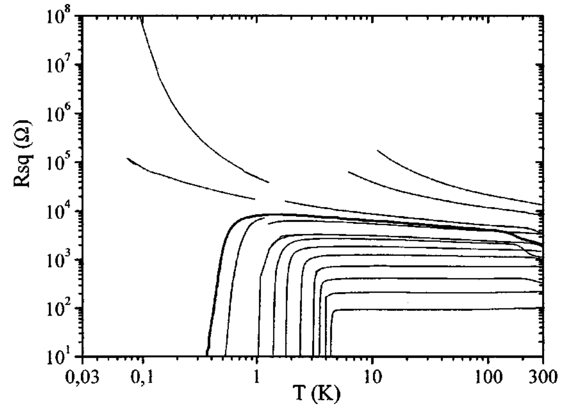

FIG. 1. Resistance per square versus temperature in TiN films of various thicknesses. The thick line is the sample NH63, the closest one to the SIT in zero magnetic field.

with $t = R^*_{sq}/\pi R_Q$ and $\gamma = \ln(\hbar/k_B T_{c0}\tau) \geqslant 0$.

R^*_{sq} is the Boltzmann sheet resistance of the film (without the localization correction) and $R_Q = h/e^2$. We choose to estimate the critical temperature as the temperature at which the resistance has decreased by 90%. On Fig. 3 the evolutions of the critical temperature with the bare normal resistance (solid symbols) and the estimated Boltzmann resistance (open symbols) are plotted, together with the fit by Eq. (1). The free parameter γ is 6.8, which indicates a strong disorder case. γ is mainly determined by the initial slope at small resistance. A good accordance with the theory is observed all over the range of resistance if the Boltzmann resistance is considered. Using $T_{c0} \simeq 4.6$ K, this gives $\tau \simeq 1.8 \times 10^{-15}$ s. For the NbN series (not shown) we deduce approximately the same Finkel'stein parameter.

The value obtained for the elastic mean free time in TiN is realistic and comparable to previous independent reports.[15,16,18] Using this mean free time and the measured resistivity of 1.1 mΩ cm for a 100 nm thick film, we estimate a free electron density of about 2×10^{21} cm^{-3}, a diffusion constant of $D \simeq 1 \times 10^{-4}$ m^2 s^{-1}, and a mean free path

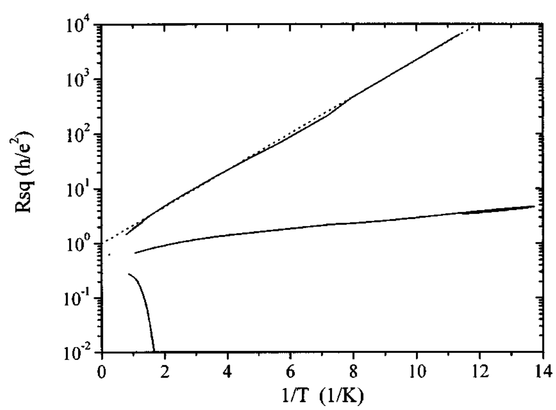

FIG. 2. Resistance per square versus inverse of temperature for three TiN films on the insulating side of the transition. These samples, shown here below $T = 1$ K, are the three most resistive samples in Fig. 1. The dashed line is an Arrhenius law with an activation energy of $k_B \times 0.77$ K.

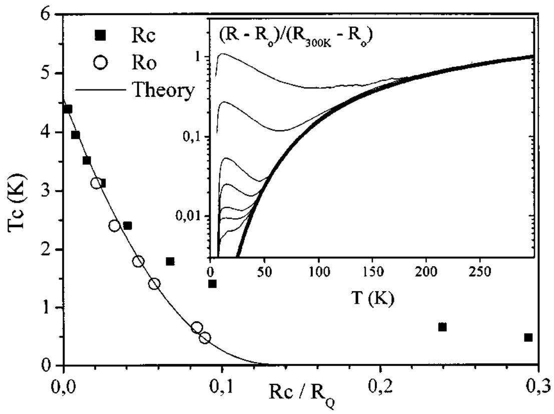

FIG. 3. Superconducting critical temperature versus resistance per square (normalized to the quantum of resistance h/e^2) for the TiN samples shown on Fig. 1. The solid symbols are for the bare resistance measured above the critical temperature, the open symbols are for the Boltzmann resistance after substraction of the logarithmic localization-interaction correction (see text). The solid line is the theoretical fit of Ref. 2 with the adjustable parameter $\ln(k_B T_c \tau/\hbar) \simeq 6.8$. Inset: The normalized resistance variation versus temperature for the most conductive films. The thick line is the Debye phonon contribution with a Debye temperature of 540 K. At low temperature the localization-interaction correction overwhelms the phonon contribution to the resistance. This correction is more and more prominent as the resistance per square increases. From bottom to top: $R_{sq} = 7.45, 8.19, 15.0, 46.7, 98.3, 220, 421\ \Omega$.

of $l \simeq 0.7$ nm. This corresponds to a $k_F l = 3Dm/\hbar$ parameter of the order of 2.6, characterizing a strongly disordered metal. Nevertheless, Eq. (1) has been obtained in the two-dimensional case, i.e., for an elastic mean free path larger than the film thickness W. This hypothesis is not fullfilled with this value of the elastic mean free path. Only the thermal length $L_T = \sqrt{\hbar D / 2\pi k_B T}$ is larger than the thickness W at low temperature. In that case the elastic mean free time τ in Eq. (1) should be replaced by the diffusion time through the thickness: $\tau^* = (W/l)^2 \tau = W^2/3D$. Although we have no way to precisely characterize the thickness of our films below 10 nm, we estimate their thickness around 5 nm, by scaling down the deposition parameters (time and power of sputtering). $\tau^* \simeq 1.8 \times 10^{-15}$ s and $W \simeq 5$ nm would give $D \simeq 5 \times 10^{-3}\ \mathrm{m^2\,s^{-1}}$, a too large value as compared to the literature, even larger than measured for thin films of pure metals. Therefore we eliminate the $l \geqslant W$ hypothesis.

We have to estimate the localization correction to the Boltzmann resistance in the normal state. The procedure is the following: first we deduce the phonon contribution which is only significant for thick films (see Fig. 3 inset) by using the empirical equation

$$\rho_{\mathrm{ph}} \propto (n-1) T_D (T/T_D)^n \int_0^{T_D/T} \frac{x^n dx}{(e^x - 1)(1 - e^{-x})} \qquad (2)$$

with $n = 3$ and the Debye temperature $T_D = 540$ K. Then we observe that the residual resistance follows a typical logarithmic correction in temperature (see the most resistive films in the inset of Fig. 3), that we attribute to the 2D localization-interaction correction (here 2D means $L_T \geqslant W$): $\delta\sigma = -(C/\pi)\ln(\hbar/k_B T\tau)$, where σ is the sheet conductance (counted in units of e^2/h). Our best fit gives the prefactor $C = 3.46$. We attribute $C \geqslant 1$ to an additional contribution of electron-electron interaction given by $\delta\sigma = -(1/\pi)(1 + \frac{3}{4}\lambda_\sigma^{j=1})\ln(\hbar/k_B T\tau)$ for $d = 2$, where $\lambda_\sigma^{j=1}$ is a phenomenological parameter related to the direct term of the Coulomb repulsion in the Hartree-Fock approximation.[19,20]

We substract this component to deduce the Boltzmann resistance. In thick films the Boltzmann resistance is close to the resistance in the normal state just above the critical temperature, because the localization correction is small. But in thinner films the correction is large, up to $\delta R_{sq} = 0.2h/e^2$ for NH63 (see Fig. 3, main panel, most right point).

In addition to this theory for homogeneous disordered 2D superconducting films, a logarithmic correction in temperature has recently been calculated in the case of granular metals:[5,21] $\delta\sigma/\sigma = -(2\pi g_{\mathrm{tun}} d)^{-1} \ln(g_{\mathrm{tun}} E_C / k_B T)$ where g_{tun} is the dimensionless tunneling conductance between grains and E_c is the Coulomb energy on each grain (in our films E_c is maximized by $k_B \times 0.77$ K, as deduced from the temperature dependence in one insulating sample). The correction is logarithmic whatever the dimension and does not depend on the mean free path. For $d = 2$, $\sigma = 2e^2/hg_{\mathrm{tun}}$, and we get $\delta\sigma = -(1/2\pi)\ln(g_{\mathrm{tun}} E_C / k_B T)$. We note that the prefactor of the logarithmic dependence of $\delta\sigma$ (counted in units of e^2/h) versus $\ln(T)$ is close to $-\pi^{-1}$, as for the weak localization and interaction correction in homogeneous disordered 2D films.

To summarize this section, we estimate that our films are homogeneous with respect to evolution of the critical temperature with sheet resistance. A good accordance is found with the Finkel'stein theory, with a large γ parameter (films with strong homogeneous disorder). The logarithmic T-dependent correction to the Boltzmann conductance is larger by a factor of 3 as compared to the weak localization, interaction, and granularity corrections taken independently. Despite the observed evolution of the critical temperature, which characterizes an homogeneously disordered film, we now show that our NH63 sample exhibits a double reentrant behavior in the magnetic field which has only been reported up to now in artificial networks of Josephson junctions.

IV. MAGNETIC FIELD DRIVEN SUPERCONDUCTOR-INSULATOR TRANSITION AND THE DOUBLE REENTRANT BEHAVIOUR

Figure 4 shows the evolution of the resistance with temperature for various magnetic fields (applied perpendicular to the films) in different samples: the top panel is for the TiN sample very close to the zero magnetic field SIT (NH63) and the bottom pannel is for the second closest to SIT TiN sample (NH57). The top panel in Fig. 5 is a blow-up of the top panel of Fig. 4, which reveals that the critical TiN film exhibits a double reentrant behavior below $T \simeq 200$ mK in a certain range of magnetic field around $H = 1.635$ T. By this we mean that $\delta R/\delta T$ is negative above $T \simeq 600$ mK, then is positive between $T \simeq 600$ and 150 mK, then is negative again

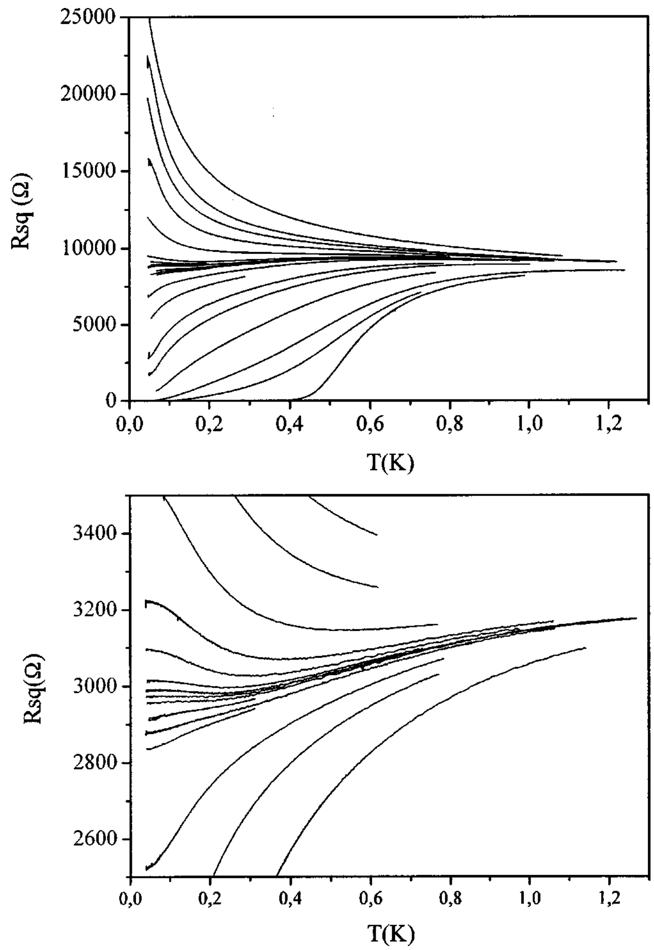

FIG. 4. Resistance per square versus temperature for various perpendicular magnetic fields. Top: TiN NH63 sample. From bottom to top the magnetic field is 0, 0.3, 0.6, 1.0, 1.3, 1.4, 1.5, 1.57, 1.60, 1.615, 1.625, 1.635, 1.64, 1.645, 1.65, 1.7, 1.8, 1.9, 2.0, 2.4 T. Bottom: TiN NH57 sample. From bottom to top the magnetic field is 2.8, 2.9, 2.95, 2.985, 2.99, 2.995, 3.0, 3.002, 3.005, 3.01, 3.02, 3.1, 3.25, 3.5 T.

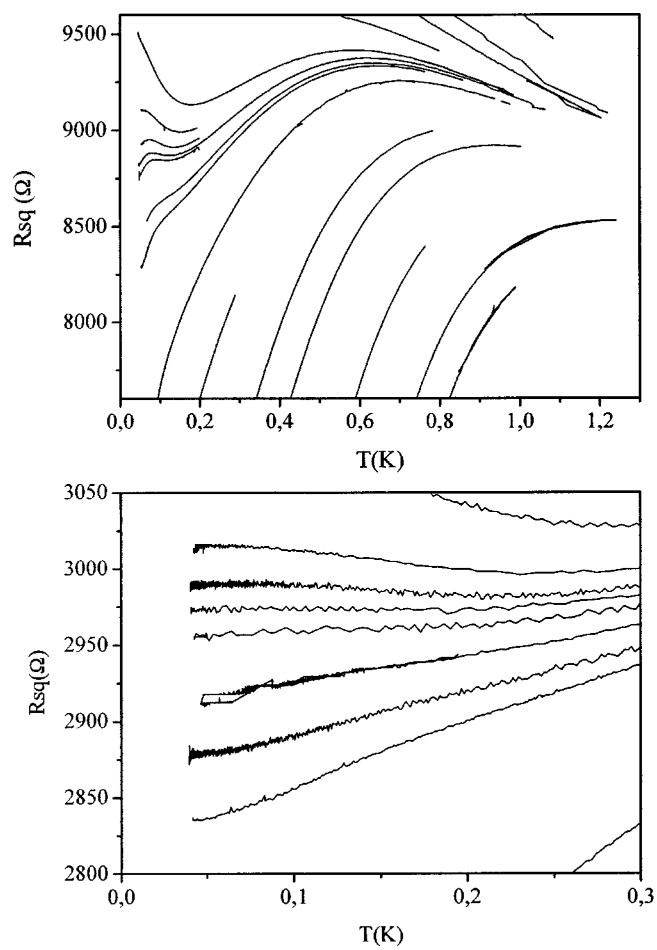

FIG. 5. Top: Resistance per square versus temperature for various perpendicular magnetic fields for the TiN NH63 sample (detail of Fig. 4, top). From bottom to top the magnetic field is 0, 0.6, 1.0, 1.3, 1.4, 1.5, 1.57, 1.60, 1.615, 1.625, 1.635, 1.64, 1.645, 1.65, 1.7, 1.8, 1.9, 2.4 T. Note the double reentrant behavior below $T = 0.2$ K. Bottom: NH57 TiN sample. From bottom to top the magnetic field is 2.95, 2.985, 2.99, 2.995, 3.0, 3.002, 3.005, 3.01, 3.02 T. Note the different behaviors for both samples, measured in the same setup: double reentrant behavior (NH63) and saturation (NH57).

down to $T \simeq 70$ mK and finally is positive down to our lowest temperature of about $T \simeq 40$ mK. This behavior is only seen close to the critical magnetic field. Note that no saturation is observed at low temperature in this sample (NH63) in contrast to NH57 (see Fig. 5, bottom pannel). Because both samples are measured in the same setup and have approximately the same impedance and shape, the saturation seen below $T = 100$ mK for NH57 cannot be attributed to heating by rf pickup.

We also note that above $T = 1$ K, the magnetic field driven SIT for NH63 and NH57 samples follows the ordinary lines expected for homogeneous superconducting films. Only very low-temperature measurements discriminate the drastically different behaviors, for instance, at $H = 1.6$ T and $H = 1.64$ T (NH63). This means that the involved energy scales for this double reentrant behavior are very small, less than $k_B \times 1$ K. By contrast, the NH57 sample exhibits a simple reentrant behavior as has been observed in granular films or in artificial networks of Josephson junctions at zero magnetic field.[1,22] In a 2D network of Josephson junctions[1] the magnetic field driven transition also exhibits a double reentrant behavior for samples in the immediate vicinity of the zero field SIT [samples S4 and T2 (Fig. 8) in Ref. 1]. For these samples the ratio between Coulomb energy and Josephson energy is in the range $E_C/E_J = 1.2 - 1.7$. The nature of this double reentrant behavior is not known (see, for instance, the references quoted in Ref. 1), but our results as well as those of Ref. 1 indicates that it happens close to the critical point for the SIT in zero magnetic field. We also note that the value of the sheet resistance at which we observe the SIT in the NH63 sample is rather similar to previous reports. The value of the critical resistance is debatable as compared to the universal $h/4e^2 \simeq 6450\ \Omega$, initially predicted by Fisher[23] using charge-vortex duality arguments. For the NH63 sample, we find a larger critical value of approximately 8800 Ω, but the extrapolation down to zero temperature is hazardous, considering our observation of a double

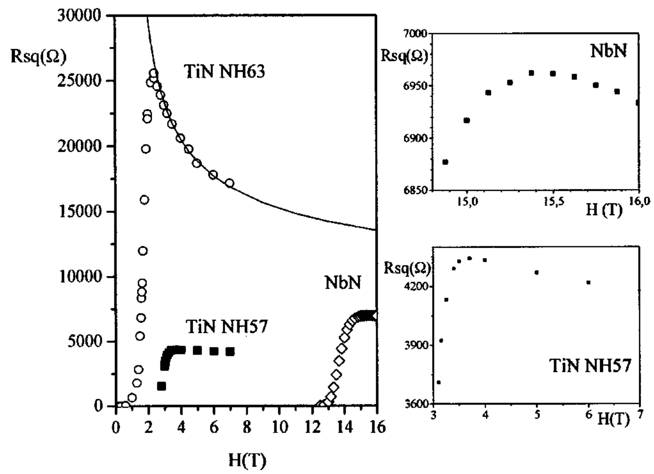

FIG. 6. Resistance per square versus perpendicular magnetic fields at the lowest temperature ($T \simeq 40$ mK). Left panel, open circles: TiN NH63 sample, solid squares: TiN NH57 sample, open diamonds: NbN sample. Top right panel: Blow-up for the NbN sample showing the negative magnetoresistance at very large field. Bottom right panel: Blow-up for the TiN NH57 sample showing the negative magnetoresistance at moderate field. The solid line is a phenomenological fit $G_{sq}(H) \times h/e^2 = 0.6 - (4/3\pi)\ln[H_{c2}/(H - H_{c2})]$ with $H_{c2} = 0.7$ T (see text).

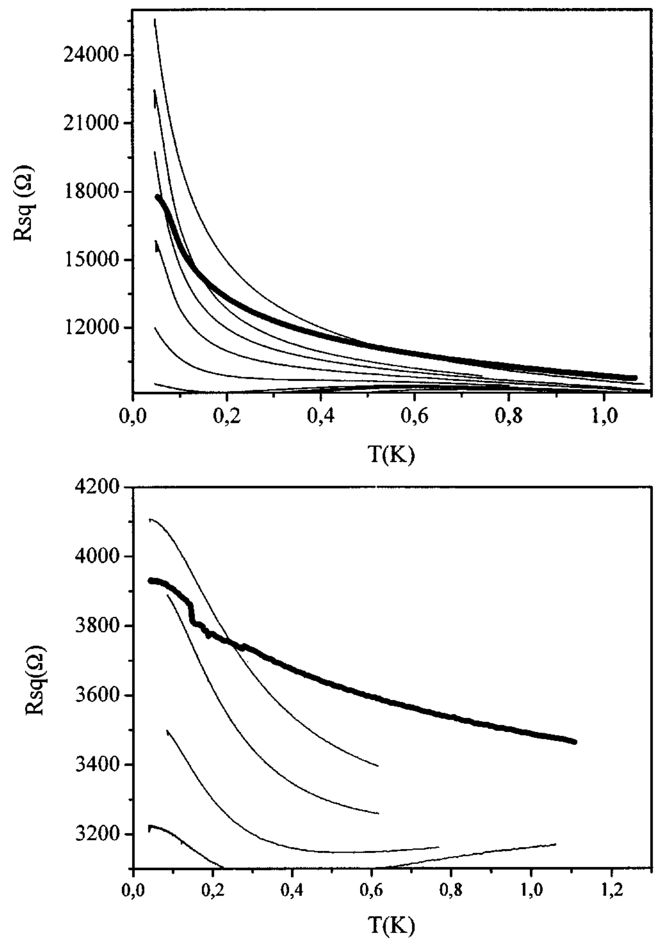

FIG. 7. Resistance per square versus temperature for various perpendicular magnetic fields. Top: NH63 TiN sample. From bottom to top the magnetic field is 1.65, 1.70, 1.80, 1.90, 2.0, 2.4 T (thin lines). The thick line is for $H = 6.0$ T. Bottom: NH57 TiN sample. From bottom to top the magnetic field is 3.04, 3.10, 3.25, 3.5 T (thin lines). The thick line is for $H = 7.0$ T. Note the negative magnetoresistance at high magnetic field and low temperature.

reentrant behavior. The critical resistance and the critical magnetic field strongly depend on the considered sample: for the NH57 sample, we found approximately 2900 Ω.

Tentatively, for NH63, we can attribute the first extremum of resistance around $T \simeq 600$ mK to the competition between the Coulomb and pairing interactions, along the lines proposed by Finkel'stein, and the second extremum around $T \simeq 70$ mK to the competition between a bulk superconducting state and a Bose insulator (i.e., a state where the Cooper pair amplitude persists everywhere, but the phase is fluctuating too much for establishing a bulk supercurrent).[11,12] Evidently more theoretical work is needed.

V. MAGNETORESISTANCE

It is interesting to study the magnetoresistance (MR) for magnetic fields larger than those needed to drive the SIT. Near the transition to superconductivity, several corrections to the Boltzmann conductivity should be considered: the paraconductivity due to the contribution of Cooper pair fluctuations (Maki-Thomson and Aslamasov-Larkin contributions), the weak localization correction which gives an effective reduction for the quasiparticle diffusion, and corrections due to the decrease of the density of states for quasiparticles which result either from Coulomb interactions (Altshuler-Aronov correction[19]) or from a fluctuating superconducting gap.[6,7]

We consider only the low-temperature situation at large magnetic field, such that the Maki-Thomson (MT) and Aslamasov-Larkin (AL) paraconductivity contributions could be neglected. These contributions are dominant near the transition and at nonzero temperature, and their suppression by a magnetic field gives a positive magnetoresistance.[6]

Figure 6 shows the MR at the lowest temperature of our experiment ($T \simeq 40$ mK) in three samples. The main observation—valid for TiN and NbN samples—is a negative MR at high perpendicular magnetic fields, as has been observed in granular aluminum and InO_x films.[24,25] This negative MR disappears at higher temperature, as shown in Fig. 7: At $T \simeq 1$ K, for instance, the MR is positive whatever the magnetic field, reflecting the dominance of the abovementioned MT and AL paramagnetic contributions, close to the critical temperature. The negative MR is much more pronounced for the NH63 sample, which is close to the SIT: for this sample the maximum sheet resistance (25 200 Ω for $H = 2.4$ T) is very close to the quantum $h/e^2 \simeq 25\,800$ Ω, and the resistance then decreases of about 8100 Ω for $H = 7$ T ($\delta R/R_{max} \simeq 0.32$) (see Fig. 6).

To interpret the results of Refs. 24, 25, Beloborodov and Efetov[4] consider the negative MR in granular films which results from the enhancement of the quasiparticle tunneling between grains when the fluctuating pseudogap in grains is reduced by magnetic field. At low temperature, the negative

MR could be as large as $\delta\rho/\rho=\delta/\Delta_0$, where δ is the mean level spacing on grains and Δ_0 is the bulk superconducting gap. This negative MR is calculated for small tunneling between grains such that the weak localization contribution is negligible. The predicted negative MR persists to much larger fields than those which destroy superconductivity inside grains and those which destroy weak localization for intergrain diffusion. Our data (see Fig. 6) show that the negative MR persists to large field. To compare with the model, one should again estimate the (residual) granularity in our films. From the Arrhenius dependence of the resistance for the TiN sample on the insulating side of the zero field SIT (see Fig. 2), we can give an overestimation for the charging energy in less resistive samples (NH63, NH57): $E_c \leqslant k_B \times 0.77$ K (for single electrons; for Cooper pairs this should be multiplied by 4). The single mean level spacing δ is also less than E_C because we do not see variable range hopping or the Efros-Schklovskii exponent for the temperature dependence in the insulating regime. Such behaviors are expected if $\delta \geqslant E_C$, i.e., if the grains could be charged.[26] Δ_0 is about $k_B \times 4.7$ K such that the maximum negative magnetoresistance according to Ref. 4 is about $\delta\rho/\rho=\delta/\Delta_0 \leqslant 0.77/4.7 = 0.16$, which is less than 0.32 observed in NH63. Therefore, we believe that negative MR for NH63 is not fully captured by the granular model.[4] This is not surprising since predictions are valid only in the limit where the dimensionless conductance is much larger than 1 (in quantum units),[6,27] a condition which should be relaxed for quantitative comparison with our experiments.

Alternatively, the case of homogeneous disordered superconducting films—without Coulomb repulsion—has been considered in Ref. 7. A negative logarithmic in magnetic field MR is predicted at very low temperature and for field $H-H_{c2} \ll H_{c2}$ in the dirty limit [$\gamma=\ln(\hbar/k_B T_{c0}\tau) \geqslant 0$, as in our films, see Sec. II]: $\delta\sigma(H)=-(4/3\pi)(e^2/h)\ln[H_{c2}/(H-H_{c2})]$. This effect includes both the variation of density of states and the (anomalous) MT term.[7] The overall picture given in Ref. 7 shows stricking similarities with our data. As an example the negative MR observed in sample NH63 above $H \geqslant 1.62$ T shows a logarithmic dependence on magnetic field, which we phenomenologically fit by $G_{sq}(H) \times h/e^2 = 0.6-(4/3\pi)\ln[H_{c2}/(H-H_{c2})]$ with $H_{c2}=0.7$ T. Here H_{c2} is a purely empirical parameter deduced from the fit. Different sets of parameters could be used such that the test of the results in Ref. 7 cannot be performed at the moment. H_{c2} is formally given by $4eDH_{c2}=(2\pi/\gamma)k_B T_c$, where $\gamma=1.781$ is the Euler constant.[7] Taking $T_c \simeq 0.5$ K from Fig. 1 and $D \simeq 1\times10^{-4}$ m^2 s^{-1}, we calculate $H_{c2} \simeq 0.4$ T.

Finally, we note that the negative magnetoresistance component is of the same amplitude as the localization-interaction-granularity logarithmic correction estimated from the temperature dependence for sample NH63 (see Sec. II and Fig. 3). At the moment, we cannot discriminate between granular and homogenous models from our magnetoresistance data, and we believe that our films present a marginal granularity.

VI. CONCLUSIONS

We studied very thin homogeneous films of nitride superconductors, across the SIT. If we restrict the analysis of the resistance versus temperature and magnetic field at temperature above $T=1$ K, a good accordance with theory for homogenously disordered films is found. Nevertheless, at lower temperature the magnetic field driven SIT for films close to the zero-field SIT exhibits a reentrance of the normal resistance. This reentrance is even a double reentrance in the immediate vicinity of the SIT, near the critical point. The reentrance has been theoretically and experimentally considered for granular systems where Coulomb energy and Josephson energies compete to localize Cooper pairs. The double reentrance is not yet understood, but our results show that its observation is not restricted to artificial Josephson arrays, where the effect has already been reported,[1] but exists as well in disordered superconducting films.

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