strain-induced valley splitting since "twofold degeneracy" was reported.

We have pointed out that the usual interpretation of SdH measurements is not valid when the Landau levels are sharp. As the width of the level grows, the motion of the Fermi level becomes less and less pronounced. If just one electron system has low mobility, the Fermi level will be fairly constant, so that the usual superposition of oscillations can be observed. We believe this is the case in measurements reported on p-channel Si inversion layers[13] and on other semiconductor materials.[14]

The authors thank the Alexander von Humboldt Foundation for an award (to A.W.O.) and the Deutsche Forschungsgemeinschaft for support via Sonderforschungsbereich 128.

[(a)]On leave from Purdue University, West Lafayette, Ind. 47907.

[1]P. Stallhofer, J. P. Kotthaus, and G. Abstreiter, Solid State Commun. 32, 655 (1979).

[2]H. Gesch, G. Dorda, P. Stallhofer, and J. P. Kotthaus, Solid State Commun. 32, 543 (1979).

[3]A. A. Lakhani and P. J. Stiles, Phys. Lett. 51A, 117 (1975).

[4]T. Neugebauer, K. v. Klitzing, G. Landwehr, and G. Dorda, Solid State Commun. 17, 295 (1975).

[5]D. C. Tsui and G. Kaminsky, Solid State Commun. 20, 93 (1976).

[6]G. Dorda, H. Gesch, and I. Eisele, Solid State Commun. 20, 429 (1976).

[7]G. Dorda, I. Eisele, and H. Gesch, Phys. Rev. B 17, 1785 (1978).

[8]F. Stern and W. E. Howard, Phys. Rev. 163, 816 (1967).

[9]Th. Englert, G. Landwehr, K. v. Klitzing, G. Dorda, and H. Gesch, Phys. Rev. B 18, 794 (1978).

[10]D. C. Tsui and G. Kaminsky, Phys. Rev. Lett. 42, 595 (1979).

[11]D. C. Tsui and G. Kaminsky, in Proceedings of the Yamada Conference on the Electronic Properties of Two-Dimensional Systems, Lake Yamanaka, Japan, 1979 (to be published).

[12]F. F. Fang, in Proceedings of the Yamada Conference on the Electronic Properties of Two-Dimensional Systems, Lake Yamanaka, Japan, 1979 (to be published).

[13]K. v. Klitzing, G. Landwehr, and G. Dorda, Solid State Commun. 14, 387 (1974).

[14]F. Koch, in Proceedings of the Yamada Conference on the Electronic Properties of Two-Dimensional Systems, Lake Yamanaka, Japan, 1979 (to be published), and references therein.

Role of Clusters in the Approach to Localization of Josephson-Coupled Granular Lead Films

A. F. Hebard and J. M. Vandenberg
Bell Laboratories, Murray Hill, New Jersey 07974
(Received 13 August 1979)

The superconducting properties of thin granular lead films provide evidence for the existence of clusters which mitigate electrostatic charging effects and allow Josephson coupling to remain favorable in samples with normal-state sheet resistance near the maximum metallic resistivity of 30 000 Ω/sq. The disappearance of zero-resistance transitions together with the precipitous onset of temperature broadening in this range is therefore most likely associated with the localization of electronic states.

The resistive transitions of thin two-dimensional (2D) granular superconducting film typically exhibit a broadening in temperature[1] as the normal-state resistance per square $R_\square^N$ approaches the maximum metallic resistivity[2] of 30 000 Ω/sq. Recent measurements on thin films of Pb, Sn, and Al also confirm that a transtiion to zero resistance will not occur if $R_\square^N$ is greater than 30 000 Ω/sq.[3] In this report we describe experimental observations on thin (~300 Å) granular lead films with $R_\square^N \lesssim 30\,000$ Ω/sq which provide new and convincing evidence for the presence of clusters[4,5] which we shall demonstrate play an important role in maintaining Josephson coupling across regions of the film containing a relatively large number of grains. The appearance on the I-V characteristics of uniformly spaced voltage steps, separated by twice the energy gap (2Δ) of lead, is a direct manifestation of the fact that Josephson phase-slip processes are occurring along the oxide boundaries separating these macroscopic-size clusters of grains. Our data show that the clusters decrease in size with increasing $R_\square^N$ and yet remain large enough to mitigate the effects of electrostatic charging[6] for films with $R_\square^N \simeq 30\,000$ Ω/sq. We argue that the absence of transitions toward the zero-resistance state above $R_\square^N = 30\,000$ Ω/sq together with the precipitous onset of temperature broadening below this value is most likely associated with the presence of localized electronic states on the superconducting clusters and/or grains.

A reactive ion-beam–sputtering technique,[7] in which the partial pressure of oxygen could be precisely controlled, was used to obtain stable films with $R_\square^N$ in the range of 100 to 10^8 Ω/sq. Stripes 2 mm in length by 20 μm wide (100 squares) and 200 μm wide (10 squares) were delineated using conventional photoresist and etching procedures. Successively higher normal-state resistances on the *same* film could be obtained by heating the film under an infrared lamp to approximately 60°C, thereby causing an irreversible increase in the oxide thickness surrounding each grain.

Figure 1 represents a family of curves of the resistance transitions of the *same* film (250 Å thick × 200 μm wide) with $R_\square^N$ ranging from 9.2

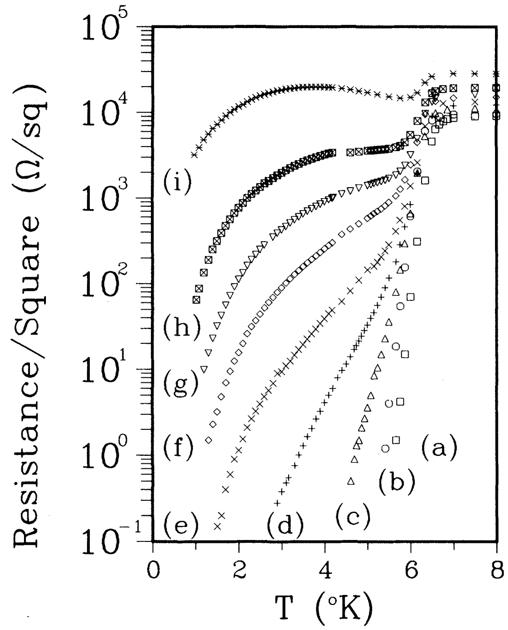

FIG. 1. Family of resistive transitions for a film 250 Å thick by 200 μm wide showing the resistance per square as a function of temperature. The monotonically increasing normal-state resistance per square [$R_\square^N$ =9.2, 9.7, 11.3, 12.3, 13.3, 15.1, 16.3, 19.5, and 28.3 kΩ/sq for (*a*) – (*i*), respectively] were obtained by successive anneals as described in the text. Excitation currents as low as 10 nA rms at 25 Hz were used to ensure that the measurements were made in regions in which the voltage scaled linearly with current. These levels were well below those currents at which voltage steps appeared.

(curve *a*) to 28.3 kΩ/sq (curve *i*). A typical resistance transition exhibits a pronounced decrease in resistance ΔR at $T \simeq 6.5$ °K, slightly less than the bulk transition temperature T_{cg} of the grains, followed by a long resistance tail heading towards a second transition to zero resistance at $T = T_{2D}$, the transition temperature of the two-dimensional film. Because the resistivity of the isolated metallic grains is insignificant compared to the resistivity of the film (> 10^4 μΩ cm), it can be argued that the first transition occurs which the grains are strongly coupled over large enough regions to short out significant portions of the sample.[8]

The distinguishing feature of the family of curves in Fig. 1 is the linear decrease in $\Delta R/R_\square^N$ with increasing $R_\square^N$. We interpret this trend as arising from the fact that the average cluster radius r_c decreases as the grains within a cluster become more decoupled with increasing oxidation. Smaller clusters are less effective in shorting out significant portions of the sample and one would expect $\Delta R/R_\square^N$ to be negligibly small when

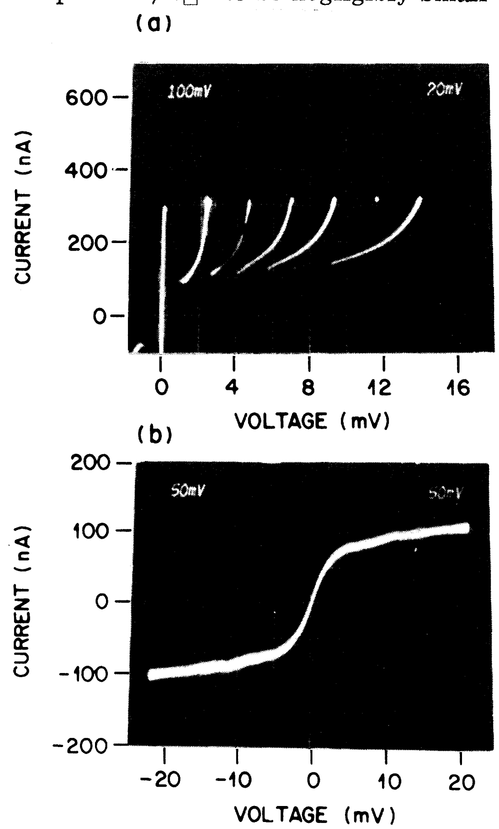

FIG. 2. Current-voltage characteristics at 1.4 K corresponding to (a) a film 300 Å thick by 20 μm wide by 2 mm long with $R_\square^N = 3.1$ kΩ/sq and (b) the 200-μm-wide film with resistive transition curve (*i*) of Fig. 1 where $R_\square^N = 28.3$ kΩ/sq.

the grains become sufficiently decoupled so that a cluster becomes equivalent to a few grains. At low $R_\square^N$, as in curve *a* of Fig. 1, a giant cluster forms at $T \lesssim T_{cg}$ which is as large as the sample and gives rise to a sudden transition to the zero-resistance state with $\Delta R \approx R_\square^N$.

The fact that Josephson tunneling occurs between clusters of grains which can be as large as the film width is shown convincingly in the *I-V* characteristics of Fig. 2(a) for a film one hundred squares in length (20 μm wide by 2 mm long) with $R_\square^N = 3.1$ kΩ/sq. This behavior is indistinguishable from that which would arise by series connection of discrete Josephson tunnel junctions of uniform area and nominally uniform critical-current densities.[9] An average tunneling resistance $\langle R \rangle$ between adjacent clusters is obtained by dividing the total sample resistance (88 kΩ) just below the ΔR feature by the total number (108) of observed critical currents. The zero-temperature critical current[10] $I_c(0) = \pi\Delta(0)/2\langle R \rangle = 2.7$ μA is somewhat higher than the range of currents (0.3–1.0 μA) actually observed. This difference could well be accounted for by the presence of thermal and external noise which could be suppressing the observed critical currents.

With the wider (200-μm) films, in which the average cluster size is small compared to the film dimensions, we usually observed a series-parallel behavior[11] characterized by a relatively small number of steps superimposed on a rising conductance background. For both the 200-μm and the 20-μm films, the number of steps decreases and becomes washed out as the critical currents decrease with increasing $R_\square^N$ and the clusters become small compared to the film dimensions. All evidence of steps has disappeared for the highest-resistance data of Fig. 1 (curve *i*) as shown in the *I-V* characteristic of Fig. 2(b) taken at 1.4 °K.

For those samples with $R_\square^N > 30\,000$ Ω/sq we always observed activated behavior with an increase in activation energy at a temperature $T = T_{cg}$ where the onset of superconductivity occurs. Similar behavior has been obtained in quench-condensed superconducting films[3] and in films of $(SN)_x$ and In (Ref. 12) and is clear evidence that the lead grains remain superconducting in the strongly activated regime. It is the *superconducting coupling between* the clusters (and grains) and not the *superconductivity of* the grains which is quenched at $R_\square^N = 30\,000$ Ω/sq.

Abeles[6] has used an earlier analysis by Anderson[13] to argue that Josephson coupling cannot be responsible for the superconductivity observed in granular metals with isolated grains because of electrostatic charging effects. The reasoning is based on the observation that, for small enough capacitance C between grains, the electrostatic charging energy $E_c = e^2/2C$ required to transfer an electron between two neutral grains can exceed, by a large factor, the Josephson coupling energy E_J between grains. More detailed analysis[14] shows that Josephson coupling between grains with n nearest neighbors is quenched when a decoupling condition $E_c > nE_J$ is satisfied, where n is the number of nearest neighbors. This decoupling condition at zero temperature can be written equivalently as

$$R_\square^N > R_\square^J = n\pi\Delta(0)CR_0/2e^2, \qquad (1)$$

where $R_0 = \hbar/e^2 = 4114$ Ω and $R_\square^J$ can be thought of as a *maximum Josephson coupling resistivity* for 2D Josephson arrays. We see from this expression that for strong enough coupling [as measured by $\Delta(0)$] and large enough C (as allowed for by the existence of clusters) $R_\square^J$ can be greater than the maximum metallic resistivity of 30 000 Ω/sq for 2D films. Clearly $R_\square^J$ depends on the microstructure of the system and is a fundamental limit to superconducting coupling *only* when $R_\square^J$ is less than approximately 30 000 Ω/sq. If we write $C \simeq \epsilon_0 \epsilon r_c d_F/d_{0x}$ for the capacitance between clusters in a 2D array of clusters of radius r_c, then we can estimate the minimum size cluster for $R_\square^J = R_\square^N = 30\,000$ Ω/sq which will give rise to resistive transitions towards zero resistance at $T = 0$. For the film in Fig. 1 and the parameters d_{0x}/ϵ 1 Å, $d_F = 250$ Å, $\Delta(0) = 1.4$ mV, and $n = 4$ we find $r_c \simeq 1000$ Å which is greater than the 250 Å and smaller diameter grains we would expect to find for this film.[15]

It is straightforward to show that the decoupling condition can also be written in the form $R_\square^N > (\frac{1}{2}n\beta)^{1/2}R_0$ where the damping parameter[16] is $\beta = 2\pi I_c R^2 C/\varphi_0$. Thus the Josephson coupling can exist at $R_\square^N = 30\,000$ Ω/sq only if the junctions are operating well into the undamped regime ($\beta \gg 1$). For an array of clusters the product RC, and hence β, is to first order independent of r_c. We conclude from these arguments that the undamped hysteretic behavior observed so clearly at lower $R_\square^N$, and in agreement with estimates of $\beta > 100$ ($R_\square^N > 58\,000$ Ω/sq for $n = 4$) for the film with *I-V* characteristics in Fig. 2(a), is preserved as the large clusters break up into smaller clusters as $R_\square^N$ approaches and passes through 30 000 Ω/sq.

In Fig. 3 we illustrate, for the data of Fig. 1,

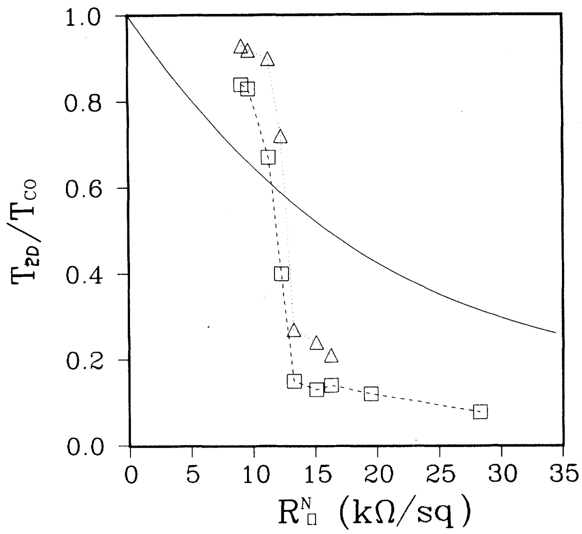

FIG. 3. Plot of the temperature ratio $T_{2\mathrm{D}}/T_{c0}$ vs $R_{\square}^{N}$ for the sequence of resistance transitions depicted in Fig. 1. Estimates of $T_{2\mathrm{D}}$ were made using a 1% of $R_{\square}^{N}$ criterion (triangles) and extrapolations with power-law fits (squares). Comparison with the Kosterlitz-Thouless prediction (solid line) is made by assuming the mean-field transition temperature $T_{c0} = T_{cg}$.

the rapid onset of temperature broadening as $R_{\square}^{N}$ increases towards 30 000 Ω/sq by plotting the temperature ratio $T_{2\mathrm{D}}/T_{cg}$ vs $R_{\square}^{N}$. As the zero-resistance transition temperature $T_{2\mathrm{D}}$ is well below the 1-K limit of our pumped helium cryostat we found it necessary to estimate $T_{2\mathrm{D}}$ using a 1% of $R_{\square}^{N}$ criterion[1] (triangles of Fig. 3) and by extrapolations based on power-law fits[8] of the form $R = R_0(T/T_{2\mathrm{D}} - 1)^s$ to the low-temperature resistive tails (squares of Fig. 3). Irrespective of the method we use to estimate $T_{2\mathrm{D}}$, the data of Fig. 3 reveal a precipitous onset in temperature broadening at $R_{\square}^{N} \simeq 12$ kΩ/sq. This behavior is in striking contrast to the smooth variation (solid line of Fig. 3) predicted by applying the Kosterlitz-Thouless notions of a vortex-antivortex dissociation transition to this thin-film system.[1]

In conclusion, it is reasonable to assume that the precipitous onset in the resistive transition broadening at $R_{\square}^{N} \simeq 12$ kΩ/sq is due to the appearance of localized states which limit the range of superconducting coupling to scale lengths, or cluster sizes, less than the sample dimensions. Accordingly, we would expect that in the range determined by $12 \lesssim R_{\square}^{N} \lesssim 30$ kΩ/sq transitions to zero resistance do not occur and that the resistance below $T \simeq 1$ K would flatten out[17] and perhaps even start to increase. For films with $R_{\square}^{N} > 30$ kΩ/sq, we have measured over thirty samples and never seen any evidence of a broad resistive transition heading towards zero resistance. It is clear that the Josephson coupling between grains breaks down at $R_{\square}^{N} = 30\,000$ Ω/sq independently of the value of $R_{\square}^{J}$ (which we have shown is greater than 30 000 Ω/sq for our ion-beam-deposited granular lead films) and is therefore concomitant with the localization of the single-particle electronic wave functions at each tunneling site.[18] In retrospect this connection should not seem surprising when one realizes that the pair and single-particle tunneling probabilities through the same oxide barrier are essentially equivalent.[10] We suspect that the distribution of grain sizes and integrain coupling strengths which favor the formation of superconducting clusters in this inhomogeneous system must be taken into account in order to understand in detail the mechanisms by which the strength and range of superconducting coupling are affected by the increasing localization of the electronic wave functions as $R_{\square}^{N}$ increases through the maximum metallic resistivity of 30 000 Ω/sq.

The authors acknowledge useful discussions with R. N. Castellano, G. Deutscher, G. J. Dolan, D. S. Fisher, P. A. Lee, S. Nakahara, and S. A. Wolf during the course of this work. We also appreciate the technical assistance of R. H. Eick and R. A. Hamm.

[1]M. R. Beasley, J. E. Mooij, and T. P. Orlando, Phys. Rev. Lett. 42, 1165 (1979).

[2]D. C. Licciardello and D. J. Thouless, Phys. Rev. Lett. 35, 1475 (1975), and J. Phys. C 8, 4157 (1975).

[3]R. C. Dynes, J. P. Garno, and J. M. Rowell, Phys. Rev. Lett. 40, 479 (1978).

[4]W. L. McClean, T. Worthington, and M. Gershenson, Bull. Am. Phys. Soc. 22, 404 (1977).

[5]E. Simanek, Phys. Rev. Lett. 38, 1161 (1977).

[6]B. Abeles, Phys. Rev. B 15, 2828 (1977).

[7]R. N. Castellano, Thin Solid Films 46, 213 (1977).

[8]G. Deutscher and M. Rappaport, J. Phys. (Paris) 39, C6-581 (1978).

[9]A. F. Hebard and R. H. Eick, J. Appl. Phys. 49, 338 (1978).

[10]V. Ambegaokar and A. Baratoff, Phys. Rev. Lett. 10, 486 (1963), and Phys. Rev. Lett. 11, 104(E) (1963).

[11]P. K. Hansma and J. R. Kirtley, J. Appl. Phys. 45, 4016 (1974).

[12]M. W. Young, J. M. D. Thomas, C. J. Adkins, and J. W. Tate, J. Phys. (Paris) 39, C6-448 (1978).

[13]P. W. Anderson, in *Lectures on the Many-Body Problem*, Vol. 2, edited by E. R. Caianiello (Academic, New York, 1964), Vol. 2, p. 127.

[14]T. Worthington, P. Lindenfeld, and G. Deutscher, Phys. Rev. Lett. 41, 316 (1978), and references, in particular, to unpublished work by de Gennes, cited therein. A recent calculation in three dimensions by W. L. McLean and M. J. Stephen [Phys. Rev. B 19, 5925 (1979)] gives similar results.
[15]S. Nakahara, results of transmission electron microscopy analysis, to be published.
[16]W. C. Stewart, Appl. Phys. Lett. 12, 277 (1968); D. E. McCumber, J. Appl. Phys. 39, 3113 (1968).
[17]Preliminary measurements at temperatures down to 50 mK, done in collaboration with D. J. Bishop, confirm this behavior.
[18]C. J. Adkins, Philos. Mag. 36, 1285 (1977).

ERRATA

GIANT DIPOLE RESONANCE IN ^{4}He WITH NONCENTRAL FORCES AND TARGET RECOIL CORRECTIONS. Dean Halderson and R. J. Philpott [Phys. Rev. Lett. 42, 36 (1979)].

References to "Solution II" and "Solution I" of Werntz and Meyerhof[1] should be interchanged. We thank Carl Werntz for pointing out the transposition. The theoretical curves on Fig. 1 should be multiplied by a factor of 1.44.

[1]C. Werntz and W. E. Meyerhof, Nucl. Phys. A121, 38 (1968).

LOW-ENERGY PION PRODUCTION AT 0° WITH HEAVY IONS FROM 125 to 400 MeV/NUCLEON. W. Benenson, G. Bertsch, G. M. Crawley, E. Kashy, J. A. Nolen, Jr., H. Bowman, J. G. Ingersoll, J. O. Rasmussen, J. Sullivan, M. Koike, M. Sasao, J. Pêter, and T. E. Ward [Phys. Rev. Lett. 43, 683 (1979)].

The incident beam energies cited were not corrected for energy loss in the beam transport system. Recent range measurements have shown that these corrections are substantial. Also, the average beam energies in the targets should have been corrected to take account of the sharply dropping production rate as a function of bean energy.

Table I gives the nominal energies quoted in the paper and the corresponding corrected incident and cross-section–averaged beam energies. The resulting shift puts the observed peak in the π^- spectrum closer to the projectile velocity and thereby strengthens the proposal that it is due to the existence of a projectilelike object after the collision. An important consequence of these new energies is to make the cross section exceed the predictions of the models described in the Letter by a factor of 5–10 at the lower beam energies.

TABLE I. The corrected beam energy (in units of MeV/u) and its average value in the targets as compared to the nominal values cited in the paper.

Nominal beam energy	Corrected beam energy	Average in the target NaF	Cu	U
125	101 ± 5	80 ± 10	⋯	⋯
150	130 ± 4	110 ± 7	⋯	118 ± 7
200	182 ± 3	164 ± 8	⋯	172 ± 8
250	235 ± 3	219 ± 5	214 ± 5	226 ± 5
400	388 ± 2	383 ± 3	377 ± 3	381 ± 3

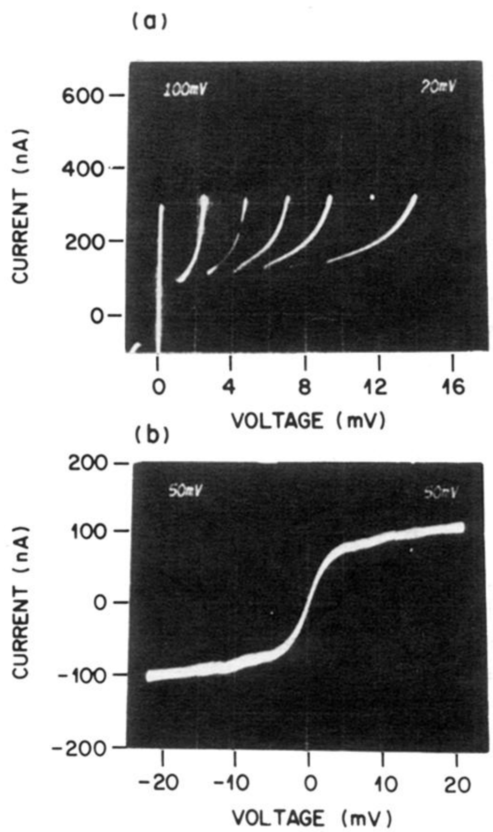

FIG. 2. Current-voltage characteristics at 1.4 K corresponding to (a) a film 300 Å thick by 20 μm wide by 2 mm long with $R_{\square}^{N} = 3.1$ kΩ/sq and (b) the 200-μm-wide film with resistive transition curve (*i*) of Fig. 1 where $R_{\square}^{N} = 28.3$ kΩ/sq.