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Correction to the two-dimensional density of states

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We report tunneling measurements on quench-condensed films of tin in order to determine the correction to the two-dimensional electronic density of states due to the effects of electron-electron interactions. Comparison with recent theory yields good agreement in the regime that the theories are valid. These measurements illustrate the nature of the corrections over a wider range of parameters.

Over the past several years, dramatic progress has been made in understanding conduction in disordered two-dimensional (2D) systems. In particular, it is generally accepted that the low-temperature resistance rise observed in such diverse 2D systems as thin metallic films and the inversion layer in silicon metal-oxide-semiconductor field-effect transitors[1–5] results from the combined effects of electron localization[6] and electron-electron interactions.[7] Localization effects have been isolated successfully by application of a small perpendicular magnetic field.[8–10] In this work, we use a tunneling measurement to isolate interaction effects.

In 1979, Altshuler and Aronov[11] predicted that interaction effects alone would lead to a change in the single-particle density of states near the Fermi level. This was extended to 2D by Altshuler, Aronov, and Lee,[7] who calculated that the first order corrections to the 2D density of states would be logarithmic. Within smearing factors of a few kT, a measure of the conductance ($G = dI/dV$) of a metal-insulator-metal tunnel junction is a determination of the relevant density of states. Imry and Ovadyahu[12] made some limited tunneling measurements on indium-oxide films. We report the results of a more consistent, comprehensive study using tin films which varied from 89 to 8900 Ω/□. Our results show conclusively that there are logarithmic corrections to the 2D density of states at low voltages, and we study these corrections up to the limit where exponential localization dominates.

In order to obtain clean tin films over a wide range of $R_{\square}$, 99.99% tin was evaporated onto fire-polished glass substrates at or below 4.2 K. First, however, four gold contacts were deposited at room temperature. Then the aluminum counterelectrode was added, also at room temperature, and oxidized in laboratory air for 10 min. This reproducibly yielded resistances between 1000 and 3000 Ω for a 1-mm^2 junction area. The sample was then cooled to helium temperatures in a low-temperature evaporator[13] and the tin film was deposited. These quench-condensed films are conducting at thicknesses much less than the equivalent film deposited at room temperature.[14] The sample geometry (inset to Fig. 1) was chosen to allow a four probe measurement of both the junction conductance and the tin film resistance.

The advantage to this technique is that we were able to use the *same* tunnel junction for an entire series of tin films. This was possible because the evaporator was designed to allow *in situ* conductance measurements. In addition, cyropumping on the walls of the chamber reduces the residual gas pressure to $< 10^{-10}$ torr. As a result, after evaporating and measuring a tin film, additional tin could be

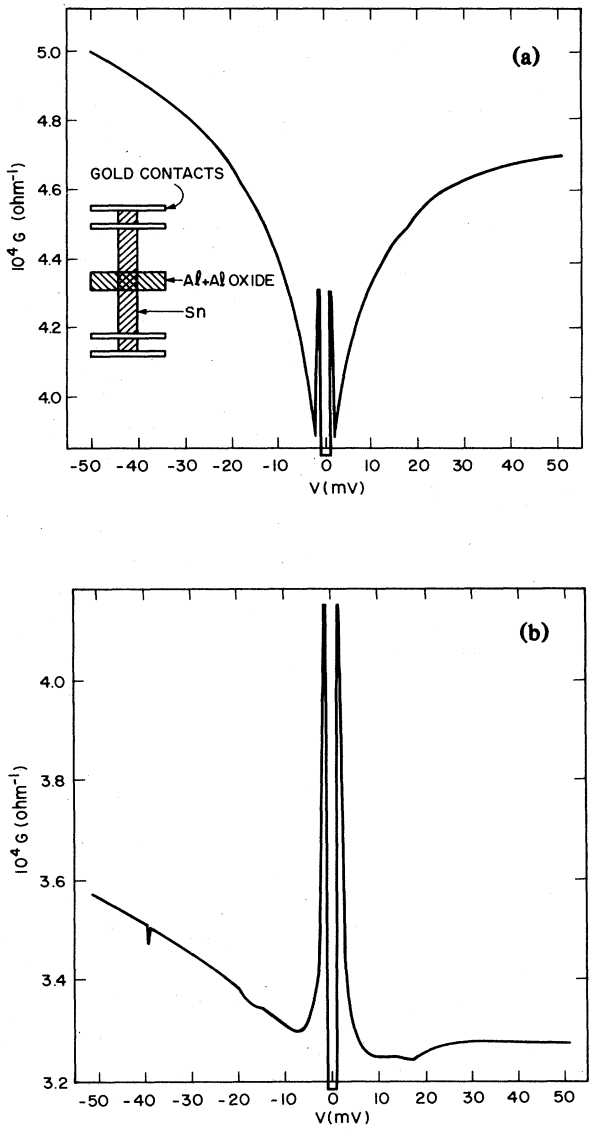

FIG. 1. (a) Trace of conductance (G) vs applied voltage (V) for a 1500 Ω junction at ~1.4 K. The tin film has $R_{\square} = 3057$ Ω. Note the superconducting energy-gap structure at 0 V and the phonon structure at ±15 mV. Inset shows sample geometry. (b) Baseline curve for the junction in (a).

evaporated without formation of an intermediary oxide layer.

The conductance measurements were made with the tin film in the superconducting state to assure that the measured voltage was being dropped entirely across the junction. Furthermore, observation of superconducting energy gap structure is strong evidence for the quality of the junction. Except for structure at the energy gap, the superconducting corrections to G are small compared to the conductance variations in the measurement so the data are interpreted as the density of states, $N(E)$. The junctions were of high quality with conductances below the gap limited by thermal excitations.

A conductance trace for a typical junction is shown in Fig. 1(a). The decrease in the conductance about the Fermi energy is large and extends well beyond 50 meV. At zero voltage, sharp features due to the superconducting energy gap appear. Additional structure observed at ± 15 meV is induced by phonons in the tin and provides further evidence of junction quality.

The component of the conductance which is an odd function of the applied voltage (V) is attributed to tunnel barrier asymmetry[15] and is also seen in the conductance of thicker, clearly 3D tin films. In order to separate out these effects, we took a baseline curve for each junction with the final, thicker film at the end of the set of measurements. Such a curve is illustrated in Fig. 1(b) for the junction of Fig. 1(a).

We measured four sets of tin films with $R_\square$'s ranging from 8900 $\Omega/\square$ ($k_F\lambda = 0.12$) down to 89 $\Omega/\square$ ($k_F\lambda \gtrsim 1$). In order to analyze the data for each set, we assumed that the measured current could be expressed as a convolution of the density of states of the superconducting tin (N_{1s}), the normal state aluminum (N_{2n}), and the asymmetric barrier potential $P(E)$:

$$I_{ns} = \int_{-\infty}^{\infty} N_{1s}(E) N_{2n}(E)[f(E) - f(E + \mathrm{eV})] P(E)\, dE \quad .$$

Taking $N_{2n}(E)$ to be a constant on this scale, the differential conductance becomes

$$G_{ns} = \frac{dI_{ns}}{dV} = N_{2n}(0) \int_{-\infty}^{\infty} N_{1s}(E) \left[-\frac{\partial f(E + \mathrm{eV})}{\partial(\mathrm{eV})} \right] P(E)\, dE \quad .$$

The Fermi function weighting factor peaks at $-\mathrm{eV}$ and has a width of $\sim \mathrm{kT}$, so at low temperatures this simplifies to

$$G_{ns} = N_{2n}(0) N_{1s}(\mathrm{eV}) P(\mathrm{eV}) \quad .$$

A reasonable deconvolution can then be accomplished by dividing the data (normalized to the conductance at 2 meV) by the similarly normalized baseline at each value of V.

The results of such an analysis are shown in Fig. 2 for one series of tin films where we have plotted the normalized junction conductance G_N vs $\ln V$. We find that all curves are approximately logarithmic over an admittedly small range at low voltages, but deviate from this dependence at higher voltages. Moreover, the deviation changes sign as values of $k_F\lambda$ for the films become greater than one.

If these low-voltage slopes, determined by a least-squares fit, are then plotted versus the $R_\square$ of the tin film, universal behavior is found (Fig. 3). The data from four sets of tin films are included and the reproducibility is impressive. Interestingly, the shape of the curve is almost logarithmic at the low $R_\square$ end and then diverges as expected when $R_\square$ approaches 10 000 $\Omega/\square$ and the screening length diverges.

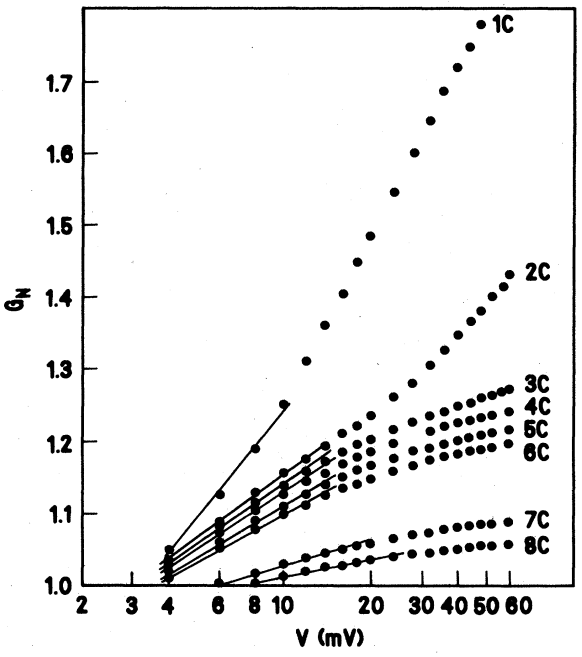

FIG. 2. Normalized junction conductance (G_N) vs applied voltage on a logarithmic scale for one series of tin films with resistances: 1C—8058 $\Omega/\square$; 2C—7092 $\Omega/\square$; 3C—6116 $\Omega/\square$; 4C—4879 $\Omega/\square$; 5C—3849 $\Omega/\square$; 6C—3067 $\Omega/\square$; 7C—1044 $\Omega/\square$; 8C—699.5 $\Omega/\square$; 9C—89.91 $\Omega/\square$. Solid lines are least-squares fits to the low-voltage data.

The result of the Altshuler, Aronov, and Lee calculation[7] of the correction to the 2D density of states due to interaction effects is

$$\frac{\delta N}{N} = \frac{-\hbar}{2\pi E_F \tau} \left(\frac{\pi}{k_F d} \right) \ln \frac{E\tau}{\hbar} \ln \frac{E}{\hbar D \kappa^2} \quad , \tag{1}$$

where τ is the elastic scattering time of the electron, d is the film thickness, $\kappa = 2me^2/\hbar$ is the inverse 2D screening length, and D is the 3D diffusivity (since the film thickness is $> \lambda$). In this expression, we have included an extra factor of $\pi/k_F d$ to partially account for the quasi-2D nature of the films ($\lambda < d < \hbar D/\mathrm{kT}$). The first logarithm results from including impurities in the calculation of the exchange term in the self-energy and the second logarithm is the consequence of an energy-dependent screening factor.

More recently, Altshuler, Aronov, and Zuzin[16] and Lopes dos Santos and Abrahams[17] have done the calculation explicitly for thin metal films. The latter find

$$\frac{\delta N}{N} = \frac{-e^2}{8\pi^2\hbar} R_\square \ln \left(\frac{E}{\hbar D (2\pi/d)^2} \right) \ln \left(\frac{(2\pi)^2 E}{\hbar D \kappa^4 d^2} \right) \quad , \tag{2}$$

with $\kappa = 2me^2/\hbar(\pi/k_F d)$, which agrees within numerical factors with a guess at a dimensionally correct version of the former. As in (1) there are two logarithmic terms, but the first logarithm now includes an energy-dependent cutoff length which was first mentioned by Imry and Ovadyahu.[12] When this length, defined as $(\hbar D/E)^{1/2}$, becomes smaller than the film thickness, the system undergoes a dimensional crossover from two to three dimensions. In our films, this length corresponds to an energy of approximately 10–20

meV which is where the conductance data deviate from the logarithm. The extra factor of $\pi/k_F d$ in κ means that the second log term is considerably larger than the first.

To first order, both (1) and (2) go as $\ln V$ at low voltages, so we can compare the magnitude of the logarithmic derivative of (1) and (2) to the data. The theoretical predictions are included in Fig. 3 and both agree in order-of-magnitude with the data for $R_\square \lesssim 1000\ \Omega/\square$ with *no adjustable parameters.* The films at larger $R_\square$ have values of $k_F\lambda$ less than 1 and are, therefore, out of the range of applicability of the theory.

Although (1) and (2) are similar, the latter may allow more insight into the situation. The dependence of (1) on $R_\square$ is strictly linear, but, because of the thickness dependence of the second logarithm, (2) has some curvature with $R_\square$. Although this is not as pronounced as the curvature in the data, the trend is the same. In the regime between $k_F\lambda = 1 (R_\square \cong 1000\ \Omega/\square)$, and the onset of exponential localization ($R_\square \cong 10000\ \Omega/\square$), the energy-dependent screening term might become even more important than is indicated by the first-order theory, enhancing the downward curvature in (2) before diverging at 10000 $\Omega/\square$.

In conclusion, we have measured the correction to the 2D density of states due to electron-electron interaction effects in quench-condensed tin films and find that it is logarithmic at low voltages. The calculations agree in order of magnitude with the data for $k_F\lambda \gtrsim 1$; however, there is a substantial region between $k_F\lambda \approx 1$ and $R_\square = 10000\ \Omega/\square$ (because these films are quasi-2D) which has not been addressed theoretically. In this region, we find an almost logarithmic variation of $dG/d\ln V$ with $R_\square$. Although the more recent calculations,[16,17] do not strictly apply here, they appear to reproduce the trends in the data and may be useful in understanding the detailed dependence.

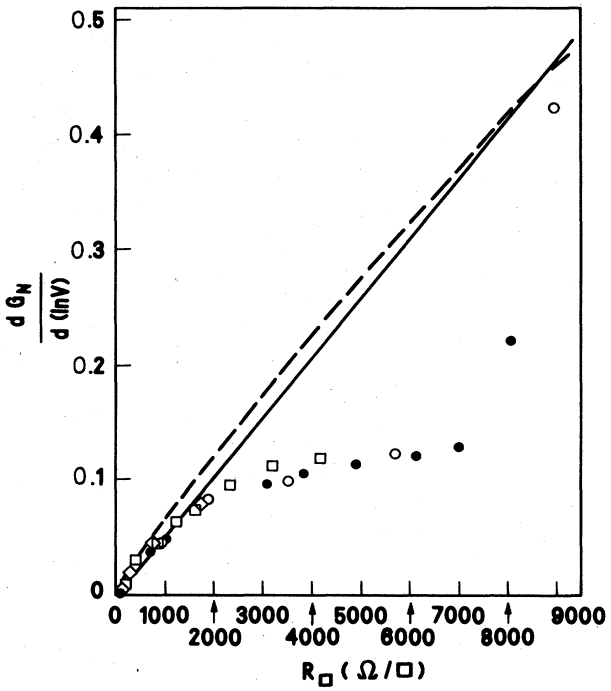

FIG. 3. Low-voltage slopes (see Fig. 2) vs $R_\square$ of the tin film for four sets of tin films. The solid and dashed lines are the theories of Refs. 16 and 17, respectively. $k_F\lambda \approx 1$ at 1000 $\Omega/\square$.

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