Masking by weak localization of metallic behavior in a two-dimensional electron system in strong parallel magnetic fields

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The rise and saturation, in a parallel magnetic field $B_{\parallel}$, of the resistivity of two-dimensional electron systems on the metallic side of the apparent metal-insulator transition (MIT) can be interpreted as the destruction of a metallic state. It is also compatible, however, with the reduction of screening of charged impurities, resulting from the breaking of the spin degeneracy, in a traditional Fermi liquid. We demonstrate, using a Si-MOSFET, that electrons in a strong magnetic field $B_{\parallel}$ parallel to the Si-SiO_2 interface may exhibit a metalliclike behavior, analogous to the effect observed at $B_{\parallel}=0$, provided that weak-localization corrections are suppressed. Conventionally, this suppression is achieved by applying a perpendicular magnetic field, but it appears that this also happens when a strong electric field is applied between the source and drain. Both methods are used in this paper. At $B_{\parallel}=0$, weak-localization corrections are also visible but the metalliclike contribution to the resistivity is greater. These results suggest that spin polarization may simply reduce the strength of mechanisms, such as screening, which contribute a metalliclike temperature dependence to the resistivity, *relative* to weak-localization corrections. The polarized system undergoes a transition from weak to strong localization at a lower critical resistivity and at a larger density than the unpolarized system.

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A significant feature of two-dimensional (2D) systems exhibiting a "metallic" behavior (defined, at finite temperature T, as $\mathrm{d}\rho/\mathrm{d}T>0$, at carrier densities n greater than some value n_{c} and resistivities ρ less than some value ρ_{c}) is the rise and saturation of ρ when a magnetic field $B_{\parallel}$ is applied parallel to the plane of the electrons. This saturation is related to the spin polarization of the electrons.[1,2] Some authors (see Refs. 3 and 4 and references therein) argue that a "metallic" behavior at $B_{\parallel}=0$ indicates a genuine metallic state, in the sense that the charge carriers are delocalized at $T=0$ (a condition that is never achievable experimentally), so that ρ remains finite. Since, in the absence of electronic interactions, quantum interference (weak localization) makes any two-dimensional system insulating, this metal can only exist by virtue of interactions that cancel weak localization. They also argue that this metal is destroyed when the electrons are fully spin polarized. Recently, a theory by Zala *et al.*[5,6] suggested that, when the Zeeman spin splitting is much less than the Fermi energy, the sign and the magnitude of the gradient $\mathrm{d}\rho/\mathrm{d}T$, in the ballistic regime, depend on the Fermi-liquid interaction parameter F_0^{σ}, which measures the strength of exchange interactions. According to this theory, F_0^{σ} accounts for the metallic behavior $\mathrm{d}\rho/\mathrm{d}T$ in zero magnetic field, but a fully polarized system must always show $\mathrm{d}\rho/\mathrm{d}T<0$. This is in agreement with the reports that n_{c} diverges as $B_{\parallel}$ increases,[2] which would signal the destruction of the "metallic" state. Experimentally, however, positive[7] and negative[8,9] values of $\mathrm{d}\rho/\mathrm{d}T$ have been observed in the spin-polarized state by different authors. The present work seeks to clarify under which conditions and by what mechanism $\mathrm{d}\rho/\mathrm{d}T$ can change sign. This is necessary to determine whether or not the observed "metallic" behavior really shows a true metal.

In noninteracting theories, an insulating behavior arising from weak localization (WL),[10] is generally expected to dominate as T decreases. This is because the phase-coherence time τ_{ϕ} increases more rapidly with decreasing T than the momentum relaxation time τ. The greater the ratio τ_{ϕ}/τ, the greater the number of self-intersecting paths that contribute to weak localization. (When $\tau_{\phi}<\tau$, there are no such paths.) Some authors[11–14] have shown that, in samples with low disorder and non-negligible interactions, in zero magnetic fields, WL corrections in $\rho(T)$, while present, may escape notice at moderately low temperatures because τ_{ϕ} is not yet much larger than τ. The negative contribution to $\mathrm{d}\rho/\mathrm{d}T$, due to WL, is then overwhelmed by other T-dependent effects, such as screening,[15,16] a semi classical effect that can produce a large, positive $\mathrm{d}\rho/\mathrm{d}T$. In this case, WL corrections in $\rho(T)$ would become dominant only at low temperatures that may not be accessible experimentally. Nevertheless, WL corrections are still manifest in $\rho(B_{\perp})$, where $B_{\perp}$ is a weak magnetic field perpendicular to the plane of the carriers, because τ and τ_{ϕ} are not dependent on small $B_{\perp}$. Thus, such a system *approximates* to a metal but is not a genuine metal. In a recent paper, Rahimi *et al.*[17] observe the disappearance of the WL corrections as n approaches n_{c}, and interpret this as an indication that interactions suppress WL. However, their analysis relies on a theory of WL that assumes $\rho\ll h/e^2$; it should be noted that, in the vicinity of n_{c}, $\rho\approx h/e^2$, so that the scattering length is comparable to the Fermi wavelength (i.e., a disorder-related effect).

The aim of this work is to investigate the "destruction" of the "metallic" behavior by a parallel magnetic field $B_{\parallel}$. Our approach consists of studying the dependence of ρ on T and $B_{\parallel}$ after having suppressed the WL corrections to ρ. Hence, it is possible to determine whether the *dominant* behavior is due to strong interactions (F_0^{σ}) or to a more familiar mechanism (WL). The results show that opposing tendencies coexist when the electrons are spin polarized: a *positive* contribution to $d\rho/dT$ exists, but is weak compared to the WL corrections. At $B_\parallel=0$, when the electrons are unpolarized, WL is also observed but it is now weaker than the positive contribution, in agreement with Refs. 11–14. This suggests that the polarized and unpolarized systems may be governed by similar mechanisms, and that the *relative* significance of these mechanisms may depend on $B_\parallel$. Once WL effects are subtracted, the two systems may be qualitatively similar. Indeed, our experimental data suggest that the spin-polarized electrons undergo a transition to strong localization, in a manner similar to the unpolarized electrons, albeit with different values of n_c and ρ_c. There need not be a fundamental transition brought about by spin polarization.

Measurements were performed on an *n*-type Si-MOSFET inversion layer with a peak mobility of 19000 cm^2/Vs, corresponding to a density $n\approx 4\times 10^{11}$ cm^{-2}. The device has a Hall-bar geometry ($1000\times 100\ \mu$m^2), permitting four-point measurements of the resistivity, using constant a.c. current at 17 Hz. Measurements were performed on a pumped helium-3 cryostat, with bath temperature ranging from 0.3 to 3.0 K, fitted with a 12-T magnet and a rotation mechanism allowing an arbitrary orientation of the 2D system relative to the magnetic field.

A magnetic field $B_\perp$ suppresses WL by introducing a phase difference between the partial waves that interfere to produce WL. This gives rise to a negative magnetoresistance[18] and to a method for measuring τ_ϕ. The conductivity $\sigma=\rho^{-1}$ is fitted to the function[19]

$$\sigma(B_\perp)=\sigma_\infty-\frac{\alpha_v g_s g_v e^2}{4\pi^2\hbar}\left[\Psi\left(\frac{1}{2}+\frac{\hbar}{4eB_\perp D\tau}\right)-\Psi\left(\frac{1}{2}+\frac{\hbar}{4eB_\perp D\tau_\phi}\right)\right]. \qquad (1)$$

Here, σ_∞ is the component of σ that does not depend on WL, g_s and g_v are the spin and valley degeneracies, $\Psi(x)$ is the digamma function, $D=2\pi\hbar^2/g_s g_v e^2 m\rho$ is the diffusion coefficient, τ is the scattering time and m is the effective mass. α_v is related to the ratio of inter-valley to intra-valley scattering rates and may take values between -0.5 and 1.

From Eq. (1), the magnitude of the WL correction, at $B=0$, is given by

$$\Delta\sigma_{WL}\equiv\sigma_\infty-\sigma(0)=\frac{\alpha_v g_s g_v e^2}{4\pi^2\hbar}\ln\left(\frac{\tau_\phi}{\tau}\right). \qquad (2)$$

$\Delta\sigma_{WL}$ is usually small compared to $\sigma(0)$ for typical electron densities and temperatures. Unfortunately, it is often difficult to determine τ_ϕ unambiguously since $\Delta\sigma_{WL}$ is a function of both α_v and τ_ϕ. The following results assume $\alpha_v=1$, corresponding to simple bands without spin-orbit scattering.

WL was measured for nonzero values of $B_\parallel$, i.e., when the electrons are spin polarized, by tilting the sample relative to the solenoid. Figures 1(a)–1(c) show the perpendicular magnetoresistivity $\rho(B_\perp)$ at a fixed density $n=2.05\times 10^{11}$ cm^{-2} for fixed total fields $B=6$, 9, and 12 T. Because $B_\perp\ll B$, we have $B\approx B_\parallel$. A weak source-drain current of rms amplitude 3 nA was used, as strong currents may

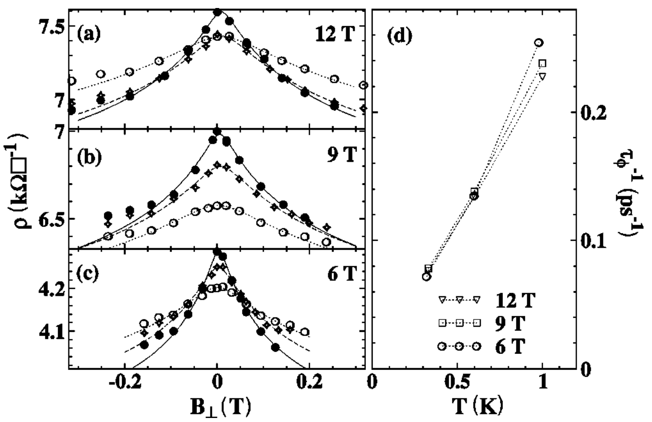

FIG. 1. (a)–(c) Magnetoresistivity in weak perpendicular fields for total fields of 12, 9, and 6 T, at temperatures (in order of decreasing line solidity) of 0.3, 0.6, and 1.0 K. The electron density is fixed at $n=2.05\times 10^{11}$ cm^{-2}. ρ was measured with a small current $I_{sd}=3$ nA. (d) The corresponding phase-coherence times ${\tau_\phi}^{-1}$ plotted as a function of T.

result in heating and may affect τ_ϕ. The peaks, characteristic of weak localization, were measured for three temperatures. The corresponding decoherence rates ${\tau_\phi}^{-1}$ were obtained from Eq. (1), with the assumptions that $\alpha_v=1$ and that the system is always near full polarization (i.e., $g_s=1$ and $D=\pi\hbar^2/e^2 m\rho$). The use of Eq. (1) is justified, in spite of the large degree of spin polarization, because the resistivity remains low at this density and an ohmic behavior is retained. In this paper, all fits to Eq. (1) treat σ_∞ as a free parameter, to accommodate any dependences on temperature and electric field. Also, since $B_\perp\ll B_\parallel$, the perpendicular and parallel magnetoresistivities can be considered independent. The results are plotted in Fig. 1(d) and are consistent with the relation ${\tau_\phi}^{-1}\propto T$ predicted by theory.[20] It is therefore unlikely that τ_ϕ is significantly affected by the small amplitude of this current.

Figures 1(a)–1(c) suggest that the sign of $d\rho/dT$ depends on the *relative* strengths of competing tendencies, negative and positive. The data were fitted to Eq. (1) to extract the resistivity in the wings of the curves, $\rho_\infty=\sigma_\infty^{-1}$, and the height of the peaks above this baseline, $\Delta\rho_{WL}=\rho(B_\perp=0)-\rho_\infty\approx-\rho_\infty^2\Delta\sigma_{WL}$. The effects of WL are contained solely in $\Delta\rho_{WL}$. As T increases from 0.3 to 1.0 K, ρ_∞ increases by (a) 4%, (b) 0.7%, and (c) 3%, while $\rho(B_\perp=0)$ decreases by (a) 2%, (b) 6%, and (c) 2%. [The values of ρ_∞ at $T=0.3$ K are (a) 6.1, (b) 5.7 and (c) 3.8 k$\Omega/\square$.] We note that, in the three cases, the net change in the height of the peak between these temperatures is approximately 6%, independent of $B_\parallel$. Thus, in Figs. 1(a) and 1(c), a negative $d\rho/dT$ is observed at $B_\perp=0$ because $d\rho_\infty/dT<-d(\Delta\rho_{WL})/dT$, whereas a positive $d\rho/dT$ is apparent for large $|B_\perp|$, where $d\rho/dT\approx d\rho_\infty/dT$. This effect is not apparent, however, in Fig. 1(b), where $d\rho_\infty/dT$ is small. Similar data obtained by other authors[11–14] for unpolarized electrons ($B_\parallel=0$) show $d\rho/dT>0$ for all values of $B_\perp$, even $B_\perp=0$. This behavior was reproduced in the present samples at $B_\parallel=0$ and is consistent with $d\rho_\infty/dT>-d(\Delta\rho_{WL})/dT$ in this regime, within the range of accessible temperatures. [At lower temperatures,

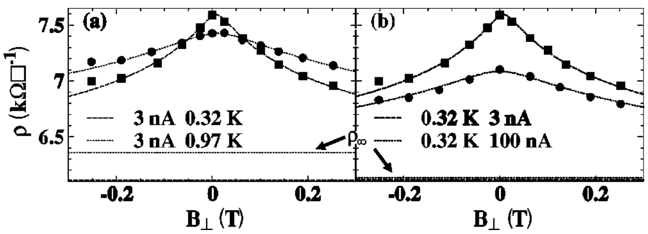

FIG. 2. Suppression of the weak localization correction $\Delta\rho_{\mathrm{WL}} = -\rho_\infty^2 \Delta\sigma_{\mathrm{WL}}$ by (a) temperature or (b) a large source-drain current I_{sd}. Electron density $n=2.05\times10^{11}$ cm^{-2}, parallel field $B_\parallel = 12$ T. For the dashed curves in (a) and (b), $I_{\mathrm{sd}}=3$ nA, $T=0.32$ K ($\tau_\phi=13.1$ ps). Dotted curves: (a) $I_{\mathrm{sd}}=3$ nA, $T=0.97$ K ($\tau_\phi=4.4$ ps), and (b) $I_{\mathrm{sd}}=100$ nA, $T=0.32$ K ($\tau_\phi=4.2$ ps). The horizontal lines below the curves show the corresponding baseline resistivities $\rho_\infty=\sigma_\infty^{-1}$ that remain when the WL corrections are subtracted.

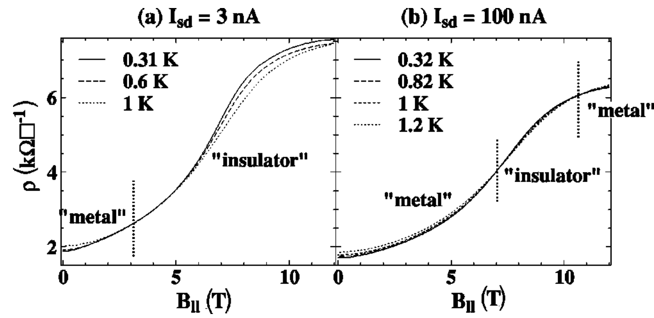

FIG. 3. Temperature dependence of $\rho(B_\parallel)$ for (a) $I_{\mathrm{sd}}=3$ nA and (b) $I_{\mathrm{sd}}=100$ nA. $n=2.05\times10^{11}$ cm^{-2}. The vertical dotted lines indicate transitions from ''metallic'' ($\mathrm{d}\rho/\mathrm{d}T>0$) to ''insulating'' ($\mathrm{d}\rho/\mathrm{d}T<0$) behavior. There is no ''metallic'' regime at large $B_\parallel$ in the low-current case (a) within the available magnetic field. For the temperatures shown, the crossing points of the curves in (b) occur at the dotted lines.

though, we expect $\mathrm{d}\rho_\infty/\mathrm{d}T < -\mathrm{d}(\Delta\rho_{\mathrm{WL}})/\mathrm{d}T$.] Figure 1 thus illustrates the principal qualitative conclusion of this paper: under some conditions, the behavior of ρ may depend strongly on whether WL is present ($B_\perp=0$) or absent [i.e., far in the wings of the $\rho(B_\perp)$ curves].

An electric field between the source and the drain can reduce τ_ϕ.[21] This may be interpreted as an electron-heating effect analogous to raising the substrate temperature.[20] However, we find that these two sources of ''heating'' affect ρ differently. Figure 2 shows the effect on $\rho(B_\perp)$ of (a) raising the substrate temperature, or (b) applying a large low-frequency source-drain current I_{sd} (or, equivalently, an electric field). The constant electron density $n=2.05\times10^{11}$ cm^{-2} and the parallel magnetic field $B_\parallel \approx B = 12$ T. At this density, $\Delta\rho_{\mathrm{WL}}$ and the heating effects are small compared to the parallel magnetoresistivity and can be treated as perturbations. [$\rho(B_\parallel=12\text{ T})-\rho(B_\parallel=0)\approx 5$ k$\Omega/\square$. This value is not strongly dependent on I_{sd} in this regime.] The dashed curves in Figs. 2(a) and 2(b) correspond to the same data, with $I_{\mathrm{sd}}=3$ nA and $T=0.32$ K. $I_{\mathrm{sd}}=3$ nA is a *weak* current that does not limit τ_ϕ. We note that the data obey $|B_\perp| \leqslant \hbar/4eD\tau \approx 0.26$ T, the range of validity of Eq. (1). We obtain $\tau_\phi=13.1$ ps. For the dotted curves, τ_ϕ is reduced to remarkably similar values [$\tau_\phi=4.3$ ps in (a) and 4.2 ps in (b)], through the two different ''heating'' mechanisms: (a) $I_{\mathrm{sd}}=3$ nA, $T=0.97$ K and (b) $I_{\mathrm{sd}}=100$ nA (corresponding to an electric field ≈ 7 mV/m between the source and the drain), $T=0.32$ K. Interestingly, these curves differ in their baseline resistivities ρ_∞, as indicated by the horizontal lines at the bottom of the figures. The temperature change (0.33–0.97 K) in Fig. 2(a) raises ρ_∞ by 4%, whereas the current increase (3–100 nA) in Fig. 2(b) raises ρ_∞ by only 0.3%. This result suggests that we cannot associate unambiguously an effective temperature with the ''heating'' produced by the large current. Despite the uncertainty in ρ_∞, due to the limited number of data points used for fitting Eq. (1), Fig. 2(a) highlights the central qualitative result of this paper: in a strong $B_\parallel$, the T dependence of ρ may be dominated by WL corrections; having subtracted these corrections, other mechanisms with a positive contribution to $\mathrm{d}\rho/\mathrm{d}T$ become apparent.

Figure 3 compares the T dependences of the *parallel* magnetoresistivity $\rho(B_\parallel)$, for the same density $n=2.05\times10^{11}$ cm^{-2} and fixed $B_\perp=0$, for small (3 nA) and large (100 nA) I_{sd}. For small I_{sd} [Fig. 3(a)], there is no positive $\mathrm{d}\rho/\mathrm{d}T$ at $B_\parallel=12$ T. As confirmed in Fig. 1(a), this is the regime where WL dominates. For large I_{sd} [Fig. 3(b)], a *positive* $\mathrm{d}\rho/\mathrm{d}T$ is observed at $B_\parallel=12$ T. This is the same qualitative behavior as that of ρ_∞ in Fig. 1(a). (The range 7 T$<B_\parallel<$11 T, shows the opposite tendency, but the behavior in this regime is likely to be complicated by the fact that the electrons are *partially* polarized.)

There is therefore no simple correspondence between the dephasing brought about by the use of a strong current and by the substrate temperature. Our results suggest that the large current reduced τ_ϕ but did not affect τ significantly ($\rho_\infty = m/ne^2\tau$). This effect will be used to probe the behavior of $\rho(T)$ once the WL corrections have been subtracted. It is emphasized that the justification for using this technique is empirical, and awaits a rigorous treatment. In the meantime, the *qualitative* conclusions of these results are consistent.

Figure 4 shows the magnetoresistivity $\rho(B_\parallel)$, measured with a current amplitude of 100 nA, for a number of electron densities n and a range of T. At $B_\parallel=0$, this system has a critical resistivity $\rho_c=42$k$\Omega/\square$, and the densities shown correspond to the ''metallic'' regime (with $\mathrm{d}\rho/\mathrm{d}T>0$). For all the densities shown, ρ increases and then saturates as $B_\parallel$ is increased. For each n, at the point $B_\parallel=B_\parallel^{c1}$, $\rho=\rho_{c1}$ (upward arrows), ρ is T independent. The value of ρ_{c1} decreases as $B_\parallel$ increases, in agreement with other experiments (for example, Ref. 22). The plateau resistivities at large $B_\parallel$ for $n \leqslant 1.25\times10^{11}$ cm^{-2} increase strongly with decreasing temperature. It is likely that this tendency will continue at still lower temperatures. This is the ''destruction'' of the ''metallic'' state that has been discussed by other authors.[4] However, this does not happen for $n \geqslant 1.56\times10^{11}$ cm^{-2}, where a *second* crossing point is visible at $B_\parallel=B_\parallel^{c2}$ (downward arrows). For these densities, a positive $\mathrm{d}\rho/\mathrm{d}T$ is recovered at large $B_\parallel$. The discussion of Fig. 2, above, suggests that this results from the *suppression* of WL by the strong current. Although magnetic fields beyond 12 T were not available,

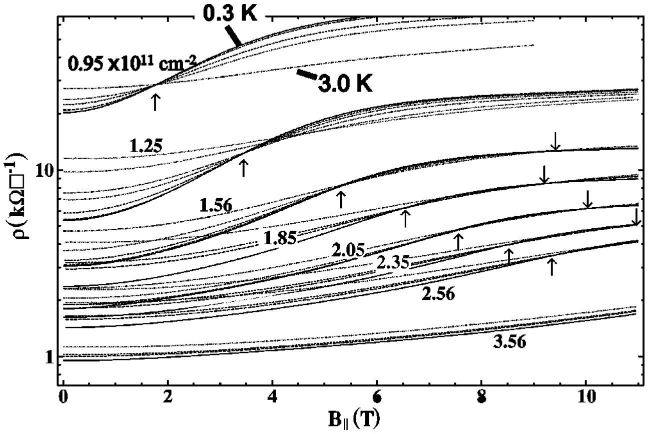

FIG. 4. Magnetoresistivity in a parallel magnetic field for a range of n and T, measured with a (strong) current $I_{sd}=100$ nA. Number labels give the density($/10^{11}\ \mathrm{cm}^{-2}$). T ranges from 0.3 (solid line) to 3 K. Arrows ($\uparrow,\downarrow$) at $B_\parallel = B_\parallel^{c1}, B_\parallel^{c2}$ indicate the crossing points of the low-T curves.

one may reasonably expect that a second crossing point is not likely to exist for $n \leqslant 1.25\times 10^{11}\ \mathrm{cm}^{-2}$.

An alternative method used for detecting $B_\parallel^{c1}$ and $B_\parallel^{c2}$ was to vary n and T, keeping $B_\parallel$ constant. The crossing points were again clearly identifiable. The results are summarized in Fig. 5, where the circles represent the crossing points obtained by sweeping $B_\parallel$, and the squares by sweeping n. The consistency between the two sets of data confirms that any hysteresis in either the field or gate-voltage sweeps, if present, is weak. Figure 5 is not intended to be a precise phase diagram, but rather to discriminate roughly regions of phase space according to the *sign* of $\mathrm{d}\rho/\mathrm{d}T$, under the condition $I_{sd}=100$ nA. Following the discussion of Fig. 2 previously, we argue that this behavior (at $I_{sd}=100$ nA) reflects the behavior of ρ_∞, i.e., the component of the total resistivity

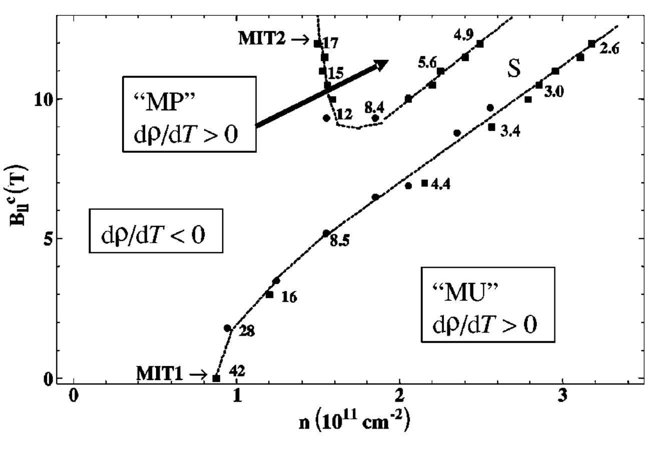

FIG. 5. Map of T-independent points, $B_\parallel^{c1}(n)$ (lower curve) and $B_\parallel^{c2}(n)$ (upper curve), obtained by sweeping $B_\parallel$ (circles) or n (squares). The current used for the measurements was $I_{sd}=100$ nA, which seems to suppress WL, as discussed in the main text. Regions on the map are labeled for convenience as MU, MP, and S. Numbers indicate the resistivity (in $\mathrm{k}\Omega/\Box$) at the crossing points. MIT1,2 (at $n=n_{c1}, n_c 2$) show the location of the MIT for unpolarized and polarized electrons.

that is independent of WL. For convenience, let us name the "metallic" region $B_\parallel < B_\parallel^{c1}$ "MU," the "metallic" region $B_\parallel > B_\parallel^{c2}$ "MP," and the intermediate region "S." The data suggest that region S marks the boundary between fully polarized (region MP) and unpolarized or partially polarized (region MU) electrons. Indeed, Figs. 4 and 5 show that S corresponds to the "knee" in $\rho(B_\parallel)$, i.e., before the electrons are fully polarized. Although $\mathrm{d}\rho/\mathrm{d}T<0$ in region S, it must not be considered as a real insulator, in the sense that ρ increases exponentially as T decreases.

The lowest point on Fig. 5 (labeled MIT1) represents the usual "$B=0$ metal-insulator transition." Here, $n_c=0.88\times 10^{11}\ \mathrm{cm}^{-2}$ and $\rho_c=42\ \mathrm{k}\Omega/\Box$. As $B_\parallel$ is increased, the transition shifts to larger n and ρ_c drops rapidly, as indicated by the numbers written along the curves. Other authors[3,22] have observed a similar effect which they interpreted as the destruction of the "metallic" state by the magnetic field. Our data indicate that, when $B_\parallel$ is sufficiently large so that the electrons become fully polarized, metallic behavior becomes apparent provided that WL corrections are subtracted. At large enough $B_\parallel$, region MP extends to arbitrarily large values of n. This is in contrast to the prediction of Zala *et al.*[6] which predicts $\mathrm{d}\rho/\mathrm{d}T<0$ even in the *absence* of WL.

There is therefore no reason, *a priori*, to discard traditional theories that predict a positive $\mathrm{d}\rho/\mathrm{d}T$ (neglecting WL). T-dependent screening[15,16] has the advantage of providing a reasonable agreement with experiment on the basis of semiclassical mechanisms. It assumes that interactions affect, through screening, only the effective disorder seen by the carriers. The screened disorder potential, treated in the Boltzmann formalism, has the form[23]

$$V(q)=V_0(q)\left(1+\frac{q_s(q)}{q}\right)^{-1}, \tag{3}$$

where $V_0(q)=e^2/4\pi\epsilon\epsilon_0 q$ is the unscreened potential, q the scattered wave vector, and q_s the screening wave vector. At $T=0$, within the Thomas-Fermi approximation ($q\rightarrow 0$), $q_s(q)$ is a constant:

$$q_s^{TF}=g_s g_v m e^2/4\pi\epsilon\epsilon_0\hbar^2=1.2\times 10^9\ \mathrm{m}^{-1}, \tag{4}$$

where $g_s=g_v=2$ are the spin and valley degeneracies in silicon. This signifies that screening is directly related to the total density of states, since electrons of either spin participate equally in screening. The random-phase approximation is more general and allows finite values of q. For $q<2k_F$ (where k_F is the Fermi wave vector), $q_s(q)=q_s^{TF}$; for $q>2k_F$, q_s is a sharply decreasing function of q. A kink (the Kohn anomaly) separates these two regimes; the smearing of this anomaly by finite temperatures produces the linear $\rho(T)$ usually associated with metallic systems[16]:

$$\rho(T)=\rho(0)[1+CT/T_F]; \quad T\ll T_F, \tag{5}$$

with $C>0$. This theory, of course, neglects WL. (Zala *et al.*[5] predict a similar form, but with C being unpredictably positive or negative.) In principle, Eq. (5) also applies to polarized electrons, but (1) $\mathrm{d}\rho/\mathrm{d}T$ should decrease since T_F doubles upon polarization, and (2) $\rho(T=0)$ should increase by a factor of approximately 4, as a consequence of the breaking of the spin degeneracy ($g_s \rightarrow 1$) (Ref. 24):

$$\rho \propto |V(q=2k_F)|^2 = |e^2/4\pi\epsilon\epsilon_0(q_s+2k_F)|^2 \propto g_s^{-2}. \quad (6)$$

(The last relation uses the facts that, for typical experimental electron densities, $q_s \gg 2k_F$ and $q_s \propto g_s$.) Both of these predictions are confirmed in Fig. 4, for those densities $n > n_{c2}$ showing the crossing at $B_\parallel = B_\parallel^{c2}$.

Of course, screening theory is only applicable when $k_F l \gg 1$ (corresponding to $\rho \ll h/2e^2 = 13\ \mathrm{k}\Omega/\square$), where l is the mean free path. Mott-Anderson localization occurs when the level of disorder reaches a critical value, in the regime where $k_F l \sim 1$.[25,26] Since the strength of screening is g_s-dependent through Eq. (4), it is possible that a system just on the "metallic" side of MIT1 in Fig. 5 undergoes a Mott-Anderson localization when $B_\parallel$ is applied. The absence of $B_\parallel^{c2}$ for $n \leq 1.25 \times 10^{11}\ \mathrm{cm}^{-2}$ in Fig. 4 is consistent with this hypothesis. If, however, the initial value of $k_F l$ at $B_\parallel = 0$ is sufficiently large (when $n > n_{c2}$), its value in region MP may remain above the critical value.

MIT2 in Fig. 5 can then be regarded as the transition analogous to MIT1 for the case of polarized electrons. Strictly speaking, MIT2 is located on the upper curve where ρ takes its asymptotic value; the data suggest that this value is less than 20 $\mathrm{k}\Omega/\square$. We shall assume that the transitions MIT1 and MIT2 can be discussed in the framework of the Mott-Anderson model.[25,26] This may be justified by considering the consequences, on the critical concentration n_c and resistivity ρ_c of MIT1, of increasing the effective disorder by applying a large negative substrate bias relative to the Si-SiO_2 interface. This reduces the mobility for a given n as a result of enhanced scattering by interface roughness or by interface charges.[27–29] In going from $V_{sub} = 0$ to -25 V, MIT1 shifts from $n_{c1} = 0.88 \times 10^{11}\ \mathrm{cm}^{-2}$ to $n_{c1} \rightarrow 1.22 \times 10^{11}\ \mathrm{cm}^{-2}$, but $\rho_{c1} = 42\ \mathrm{k}\Omega/\square$ remains constant to within 5%. This is consistent with the Anderson-Mott model[26] for a disorder-induced transition, where it is the value of ρ_c (the "maximum metallic resistivity") and not n_c that defines the MIT.

The discrepancy between our two critical resistivities is intriguing: $\rho_{c1} = 42\ \mathrm{k}\Omega/\square$, whereas $\rho_{c2} \approx 17\ \mathrm{k}\Omega/\square$, although this value, obtained at the largest available $B_\parallel$, is certainly underestimated. Nevertheless, Fig. 4 suggests that ρ_{c2} should be less than $\sim 20\ \mathrm{k}\Omega/\square$. On the basis of the Anderson localization model[25] and neglecting the effects of multiple scattering, Mott[26] predicted that the value of ρ_c should in fact be universal:

$$\rho_{max} \sim \frac{2}{g_s g_v} \frac{h}{e^2} X, \quad (7)$$

where X is a constant expressing the degree of disorder; it is related to the suppression of the density of states resulting from disorder broadening.[26] Although this formula was originally derived for states in an impurity band, Mott argued that it should apply to all Anderson-like transitions, regardless of the origin of the disorder or of the form of the conduction band. Assuming that this model applies to MIT1 (where $g_s = 2$) and MIT2 (where $g_s = 1$), it follows from the experimental data that X is roughly four times as large for MIT1 as for MIT2. In other words, a greater degree of disorder is necessary to localize the *un*polarized electrons. This is at odds with Mott's prediction that X should be universal.

As discussed previously, some authors have argued that a real metallic state exists at $B=0$ (see Refs. 3 and 4) but that it is destroyed once the spins are polarized. This occurs possibly via the quenching of the exchange interactions. However, a simpler scenario is that of a transition from weak to strong localization whose condition is polarization dependent. Mott's original treatment of the Anderson model is of course strictly valid only for a spin-degenerate system where all the quantum states are doubly occupied. This is likely to be true for extended and localized wave functions whose decay length R is large, so that the Coulomb repulsion of same-site electrons ($\propto e^2/R$) is small compared to the Fermi energy, but it may be incorrect when R is small. Reference 30 predicts a ferromagnetic ground state for disorder-localized electrons; this magnetization is destroyed when the electrons are delocalized. In this case, it is possible that localizing an unpolarized electron gas would require larger levels of disorder, to create enough localized states to accommodate the carriers of both spin orientations. This hypothesis requires further investigation.

It would be highly desirable, as a support to the above interpretation, to show that MIT1 and MIT2 represent transitions to insulators with an activated temperature dependence, as predicted by the Mott-Anderson model. Unfortunately, as MIT1 and MIT2 are approached, the system becomes increasingly non-ohmic, making Eq. (1) unreliable. It is difficult, in this regime, to discriminate the effects of interference and those arising from other mechanisms. It must also be borne in mind that the Mott-Anderson transition is an approximation that neglects interference effects. The present data suggest that, at large $B_\parallel$, these effects may well be significant. Nevertheless, the Mott-Anderson model provides a convenient framework to picture the underlying process.

The question why, in region S, the sign of the temperature dependence should differ from that in MU and MP also remains to be answered. In this regime, spin-up and spin-down electrons coexist and differ greatly in their concentrations ($n_\uparrow$ and $n_\downarrow$) and Fermi wave vectors [$k_{F\uparrow,\downarrow} \propto (n_{\downarrow,\uparrow})^{1/2}$]. Moreover, since minority spin electrons, at the Fermi energy, are screened by both majority- and minority-spin electrons, and majority-spin electrons are not screened significantly by the minority-spin electrons, the two subsystems experience different effective disorders. Scattering times and mean free paths differ correspondingly. It is therefore possible for the minority spins to satisfy the Ioffe-Regel criterion ($k_{F\downarrow} l \sim 1$) and become localized, whilst the majority spins are metallic ($k_{F\uparrow} l \gg 1$). In this model, the minority spins ($\downarrow$) conduct by hopping while the majority spins ($\uparrow$) propagate as waves. The overall behavior near full polarization is a balance between opposing tendencies, whose T dependence may be very different. A quantitative treatment would involve calculating screening effects in the regime of low $k_F l$ for partial spin polarization; this is not achievable at present.

The existence of a spin-polarized metal has been discussed by other authors. Shashkin *et al.*[8] report a transition, at both $B=0$ and large $B_\parallel$, where the activation energy E_a for conduction vanishes. (The corresponding critical densities are very similar to those for MIT1 and MIT2 in Fig. 5.) On the "metallic" side, $E_a=0$ but $d\rho/dT<0$. On the other hand, Mertes *et al.*[7] observe positive and negative $d\rho/dT$ at large n in strong fields; their results show consistency with the existence of region S in Fig. 5 without applying a perpendicular magnetic field. Neither of these papers, however, addresses the issue of WL and the need to suppress it before considering the T dependence of other mechanisms. The present work demonstrates that this is crucial. It is likely that a comprehensive experimental study of the resistivity *without* WL will provide a useful insight into the similarities in the scattering mechanisms in the unpolarized and polarized "metals."

In conclusion, our results suggest that the reason for the disappearance of the positive $d\rho/dT$ in large $B_\parallel$ is due to a change in the significance of WL corrections *relative* to other mechanisms, such as screening. The existence of weak localization at all polarizations in this regime, including at $B_\parallel=0$, suggests that the system is an insulator at $T=0$, in agreement with Refs. 13, 14, and 31. At finite T, WL corrections may be small enough to make the positive $d\rho/dT$ apparent. It is possible that this behavior can be described on the basis of screening theory, given that, for $n>n_{c2}$, $\rho(B_\parallel\rightarrow\infty)\approx 4\rho(B_\parallel=0)$. However, the WL corrections may dominate the weak positive contribution of screening to $d\rho/dT$. This role of WL was observed by quenching it, either by applying a weak perpendicular magnetic field or a large source-drain current to dephase the electrons. (The latter method, although justified empirically in this paper, awaits a full characterization and a theoretical treatment.) Thus, the system *appears* metallic for $n>n_{c1}$ and $B=0$, and insulating at large B. An interesting, and so far unresolved, issue is the difference in the "maximum metallic resistivity" for polarized and unpolarized electrons.

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