Spin-orbit scattering time and dephasing time of carriers in two-dimensional HgTe-CdTe superlattices and heterojunctions

John K. Moyle*
Department of Physics, University of Southern California, Los Angeles, California 90089-0484

J. T. Cheung
Rockwell International Science Center, P.O. Box 1085, Thousand Oaks, California 91360

N. P. Ong
Department of Physics, Princeton University, P.O. Box 708, Princeton, New Jersey 08544
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Weak-localization studies of the two-dimensional electron gas in heterostructures made of HgTe and CdTe are reported. Because of the small effective mass, the diffusion constant D is unusually large (824 cm^2/s). By fitting the weak-field transverse magnetoresistance data from high-mobility samples (50 000 $cm^2/V\,s$, $R_{\square}=720\ \Omega/\square$) to the calculation of Hikami, Larkin, and Nagaoka, we have determined the temperature (T) dependence of the spin-orbit scattering time and the dephasing time of the carriers. The system is shown to have dominant spin-orbit scattering. (The spin-orbit rate is 0.6 times the elastic scattering rate.) The dephasing time which is linear in T agrees in magnitude with a calculation by Al'tshuler, Aronov, and Khmel'nitskii. Evidence for significant Coulomb interaction effects is obtained from the zero-field conductivity. However, the magnitude of the interaction parameter $\tilde{F}_\sigma$ derived from the plot of the resistance versus $\ln T$ strongly disagrees with existing theories. Prominent anomalous magnetoresistance is also seen in the longitudinal geometry. The data are compared with the interaction theory (Zeeman splitting) and the weak-localization theory. Results from high-resistance (30 $k\Omega/\square$) samples which indicate the breakdown of perturbation theory are also reported.

I. INTRODUCTION

In two-dimensional (2D) conductors the observed increase in the resistance as the temperature decreases is due to the combined effects of weak localization[1] and Coulomb interaction.[2,3] The weak localization effects are also observable as weak logarithmic anomalies in the magnetoresistance in the weak-magnetic-field regime.[4–6] Recent progress in the understanding of these universal phenomena has made possible the determination of various carrier scattering times from magnetoresistance data alone.[7] A growing list of 2D systems has been studied using this approach. The role of spin-orbit scattering has also been clarified by the controlled introduction of strong spin-orbit scattering centers;[8] weak localization with or without strong spin-orbit scattering is now reasonably well understood. Here we apply this technique to a novel 2D system which has dominant spin-orbit coupling and a very large diffusion constant.

Weak localization arises when two wave packets traversing the same closed loop in opposite directions interfere constructively to enhance the probability amplitude for returning to the origin (compared with the Drude value). The destruction of this coherent backscattering by a weak magnetic field **B** (normal to the 2D sample) is observed as a negative magnetoresistance which is apparent in very weak fields.[3] In a material with very strong spin-orbit scattering, however, the interference between the two wave functions at the origin is destructive rather than constructive when B is zero.[7–9] (The change in sign arises from the transformation properties of the $s=\frac{1}{2}$ spinor under rotations.) Hence the probability for returning to the origin is reduced from the Drude value (anti localization.) The destruction of this interference by a weak B field is observed as a positive magnetoresistance. These effects have been studied in a great many metallic films [Au-Pd (Ref. 10), Mg (Ref. 11), Cu (Ref. 12)] and systems made from semiconducting hosts (silicon field effect transistors[13] and $Al_xGa_{1-x}As$-GaAs heterostructures[14]). In contrast to the metallic films the 2D systems made from semiconducting hosts are characterized by very small carrier effective masses m^*, low carrier density, and large diffusion constants D. Transport in such systems extends the study of these effects to new regions in the parameter space. For instance, both weak-localization effects and the quantum Hall effect can be observed in the same semiconducting 2D sample as opposed to the situation in the metallic films where only localization effects can be studied. For these reasons the investigation of 2D semiconducting systems, especially those with parameters quite different from $Al_xGa_{1-x}As$-GaAs heterostructures and Si MOSFET's (metal-oxide-semiconductor field-effect transistors) promises to be quite fruitful. To date the observation of positive magnetoresistance due to the suppression of antilocalization has been confined to metallic films of the heavy elements (Au,Bi,Pt,Pd).[15,16] In the Si MOSFET and $Al_xGa_{1-x}As$-GaAs heterostructures the spin-orbit scattering is usually negligible so that the weak-field magnetoresistance is negative. In this paper we report a detailed study of the positive magnetoresistance due to the suppression of antilocalization in superlattices made of HgTe and CdTe. Both compounds have strong spin-orbit coupling parameters.

Bulk HgTe is a negative-band-gap material[17] in which the conduction band consists of electronic states with very light effective mass ($0.02m_0$ to $0.03m_0$) while the valence band consists of "heavy holes" ($0.5m_0$). The two bands[18] overlap slightly near the zone center which has Γ_8 symmetry. (The electronic states have the symmetry with total angular momentum $j=\frac{3}{2}$, $m=\frac{1}{2}$ while the hole states transform as $j=\frac{3}{2}$, $m=\frac{3}{2}$.) On the other hand, bulk CdTe is a good insulator with an energy gap of 1.6 eV. Because of the small lattice mismatch between the two compounds (0.3%) HgTe-CdTe superlattices with interesting properties have been fabricated by several groups in recent years.[19–22] The carriers are confined to the HgTe layers which are of the order of 100 Å wide in our samples. Estimates[23,24] of the valence-band offset between the HgTe and CdTe vary from 40 to 300 meV. The electronic band structure in the HgTe-CdTe heterojunction is not well known and various models exist in the literature. Experimentally, superlattice samples with fairly similar parameters have been reported by different groups to be semiconducting as well as metallic. In our superlattice samples the carriers are predominantly electrons with densities which are somewhat sample dependent. The conductance remains quasimetallic down to 100 mK showing an increase of the resistance (as $\ln T$) as T decreases below 10 K. In high-mobility samples this increase is quite weak for T above 0.5 K. Quantized Hall steps are observed[20,25,26] in samples 1 and 2 which have mobilities exceeding 50 000 $cm^2/V\,s$. These samples (consisting of 12 periods of HgTe-CdTe of thicknesses 90 and 40 Å, respectively) also show positive logarithmic magnetoresistance at fields of the order of 1 mT.

II. FIELD SUPPRESSION OF COHERENT BACKSCATTERING

The identification of the magnetoresistance (MR) anomalies with the suppression of coherent effects in a strong spin-orbit scattering material leads to a close fit of the data to the theoretical expressions and reasonable values of various scattering times extracted from the fits. Perhaps the strongest evidence in favor of this interpretation is the temperature dependence of the two scattering times $\tau_{s.o.}$ (spin-orbit scattering time) and τ_ϕ (the dephasing time.) In the two samples the variation of τ_ϕ and $\tau_{s.o.}$ with T are consistent with that found in other systems, and as expected from the theory.

The superlattices were grown using a laser-assisted molecular-beam epitaxial technique on substates of CdTe oriented in the [110] direction at a substrate temperature of 160 °C. Current and voltage leads were attached in the usual four-probe configuration with In solder. Conventional ac lock-in techniques were used to measure the low-field MR. For temperatures T above 40 K the R-versus-B curves are parabolic in the $B \rightarrow 0$ limit. At lower temperatures an anomaly centered at $B=0$ with wings which decay as $\ln B$ becomes increasingly prominent as T decreases. We digitized the recorder traces and performed a best fit to the expression derived by Hikami, Larkin and Nagaoka,[9] for the transverse MR in a 2D system

$$\Delta\sigma(B) = -\alpha_B\sigma_N[\psi(\tfrac{1}{2}+1/a\tau)-\psi(\tfrac{1}{2}+1/a\tau_1) + \tfrac{1}{2}\psi(\tfrac{1}{2}+1/a\tau_2)-\tfrac{1}{2}\psi(\tfrac{1}{2}+1/a\tau_3)] , \quad (1)$$

where α_B is a constant of order 1, $\sigma_N = e^2/2\pi^2\hbar$, ψ is the digamma function, $a=4DeB/\hbar$, τ is the elastic scattering time, and the $1/\tau_n$'s are combinations of the various scattering rates

$$1/\tau_1 = (4/\tau_{s.o.}+2/\tau_s+1/\tau_\phi) ,$$

$$1/\tau_2 = (6/\tau_s+1/\tau_\phi) ,$$

$$1/\tau_3 = (4/\tau_{s.o.}+2/\tau_s+1/\tau_\phi) .$$

We will set the spin-flip scattering rate $1/\tau_s$ to zero. We have assumed isotropic 3D scattering for the spin-orbit scattering time ($\tau^x_{s.o.}=\tau^z_{s.o.}$) and have incorporated the correction noted by Maekawa and Fukuyama[27] in $1/\tau_1$. They have also generalized Eq. (1) to include Zeeman splitting effects. These effects are negligible for the fields used in the transverse MR (< 10 mT). In Sec. VI we discuss the longitudinal MR.

The fits (solid lines) and the data are shown in Fig. 1 for selected temperatures. From the fits at each T we extracted the scattering times τ_ϕ and $\tau_{s.o.}$ as well as the amplitude α_B. (The elastic scattering time τ was determined separately by the mobility and not varied in the fitting process.) From the plot in Fig. 2 we find that the spin-orbit scattering time $\tau_{s.o.}$ is insensitive to T. This is consistent with the nature of spin-orbit scattering which originates from scattering off impurities. On the other hand, the dephasing rate $1/\tau_\phi$ is observed to be linear in T

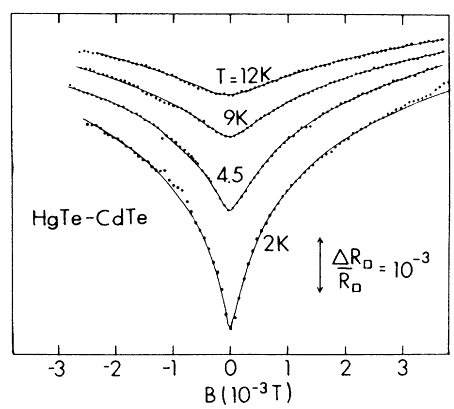

FIG. 1. The weak-field transverse magnetoresistance of HgTe-CdTe superlattice (sample 1) at various temperatures. The logarithmic anomaly at $B=0$ arises from coherent backscattering effects with dominant spin-orbit scattering. Solid lines are fits to Eq. (1). (See text.) The sample consists of 12 periods with thickness 90 Å (40 Å) for the HgTe (CdTe) layers. The electron mobility is 53 000 $cm^2/V\,s$.

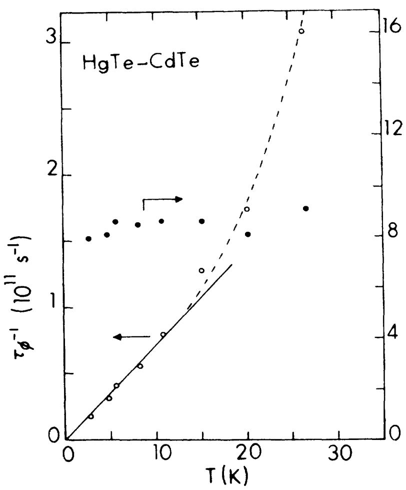

FIG. 2. The temperature dependence of the dephasing rate $1/\tau_\phi$ and the spin-orbit scattering rate $1/\tau_{s.o.}$ extracted from the fits to the data in Fig. 1. The solid line indicates linear T dependence. The dashed curve is a guide to the eye. Note that $1/\tau_{s.o.}$ is insensitive to T.

below 15 K. Above 15 K the dephasing rate increases faster than linear dependence. Recent theories interpret the linear dependence as arising from electron-electron scattering.[28–30] Although there exists some differences in the literature as to the theoretical magnitude of τ_ϕ (and its T dependence) we find that the value in Fig. 2 is in fair agreement with the expression of Al'tshuler, Aronov, and Khmel'nitskii.[30] The more rapid increase in the dephasing rate above 15 K is due to the contribution of inelastic processes due to to phonon scattering.

III. SCATTERING TIMES

From the fit of the weak localization theory to the MR data we can extract the characteristic fields which are related to τ_ϕ and $\tau_{s.o.}$ as follows:

$$B_\phi = \hbar/(4eD\tau_\phi), \quad B_{s.o.} = \hbar/(4eD\tau_{s.o.}) \ . \tag{2}$$

[The spin-orbit diffusion length and dephasing length can also be defined as $l_{s.o.}^2 = D\tau_{s.o.}$ and $l_\phi^2 = D\tau_\phi$, respectively. Using Eqs. (2) and the values of B_ϕ and $B_{s.o.}$ derived from the fits we find that $l_\phi = 3.47$ μm at 1 K and $l_{s.o.} = 0.310$ μm.] To convert these fields to relaxation times we determined the diffusion constant D by using the measured conductance and Hall data at low fields. An uncertain factor in our analysis is the number of layers N actively participating in current transport. However, because the CdTe layers are only 40 Å thick and In (the solder material used to attach current probes) has a very high diffusivity in both HgTe and CdTe we have assumed that N equals 12 in sample 1. The sample resistance at 4 K $R_\square$ equals 720 Ω/□ per layer. From the relation

$$m^* D/\hbar = \sigma/(2e^2/h) \ , \tag{3}$$

we calculate the metallic parameter $m^* D/\hbar$ ($=\varepsilon_F \tau/\hbar$) to be 17.9 (where ε_F is the Fermi energy.) From cyclotron resonance measurements[31] on samples cut from the same wafer the effective mass m^* is determined to be $0.025m_0$. Thus D calculates out to be 0.0824 m^2/s. This is comparable to the value found in high-mobility $Al_xGa_{1-x}As$-GaAs heterojunctions.[14] The Hall constant R_H in our sample which equals $322m^2/C$ in the limit of $B \to 0$ implies that the carrier density per layer n_s equals 1.62×10^{11} cm^{-2} and the electron mobility μ is 53 700 cm^2/V s. Using these numbers the potential (elastic) scattering time τ and mean free path l_0 are 7.69×10^{-13} s and 0.356 μm, respectively. Finally, using the inferred value of D we find from the data in Fig. 2 that the temperature-independent spin-orbit scattering rate $1/\tau_{s.o.} = (8.0\text{–}8.5) \times 10^{11}$ s^{-1} and the dephasing rate $1/\tau_\phi$ equals 6.93×10^9 s^{-1} at 1 K. The value of α_B obtained from our fits is weakly temperature dependent (Fig. 3), varying from 0.42 at 2.8 K to 0.3 at 20 K.

There exist several checks for the consistency and reasonableness of these numbers. First, in order for the theory to be applicable the weak-localization anomaly should be observed at fields corresponding to a magnetic length l_B which is larger than the mean free path l_0. The physical picture is that the destruction of the interference effects occurs when the flux through the closed loop equals one flux quantum, i.e., when the area of the loop is of the order of l_B^2. For the diffusion picture to be valid the carrier must scatter off impurities repeatedly within this area, which implies that $l_0 < l_B$. From this condition we deduce that the field at which this picture breaks down equals 4 mT in sample 1. This is quite consistent with the data in Fig. 1 which show that the anomaly occurs at fields smaller than this value. Secondly, for the theory to be consistent the spin-orbit scattering rate cannot exceed the potential scattering rate. The derived numbers imply that $1/\tau_{s.o.}$ is of the order of 0.65 times $1/\tau$. In a semiconductor where spin-orbit scattering is dominant Al'tshuler *et al.*[32] have argued that the two rates $1/\tau_{s.o.}$ and $1/\tau$ are comparable as we find here.

Finally, the magnitude of the dephasing rate has been calculated by several groups assuming that the main inelastic process arises from electron-electron scattering. Using the expression given by Al'tshuler, Aronov, and Khmel'nitskii,[30]

$$1/\tau_\phi = k_B T/(2m^* D)\ln(m^* D/\hbar) \ , \tag{4}$$

where k_B is Boltzmann's constant, we find that the theoretical value of $1/\tau_\phi$ at 1 K equals 1.05×10^{10} s^{-1}

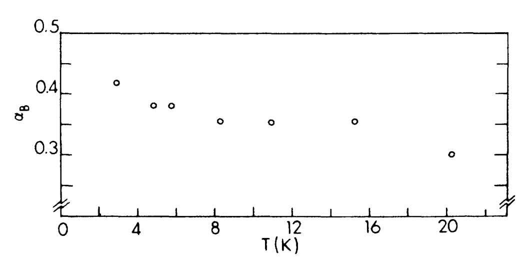

FIG. 3. The temperature dependence of α_B extracted from the fits to the data in Fig. 1.

[using the metallic parameter value found in Eq. (3)]. This is in order-of-magnitude agreement with the value found above.

IV. COMPARISON WITH ZERO-FIELD R VERSUS T

The zero-field resistance of all HgTe-CdTE superlattice and single-layer samples investigated by us shows an increase of R as T decreases qualitatively as predicted by both weak-localization and interaction theory, viz.

$$\Delta\sigma(T,B=0)=Ce^2/(2\pi^2\hbar)\ln T \ , \quad (5)$$

$$C=\alpha p+(2-\tfrac{3}{2}\widetilde{F}_\sigma)/2 \ . \quad (6)$$

C is comprised of the weak-localization term αp and the Coulomb interaction terms. α equals 1 for the case of zero spin-orbit scattering and $-\frac{1}{2}$ for dominant spin-orbit scattering. p is defined by $1/\tau_\phi \sim T^p$, and equals 1 from the data in Fig. 1. The interaction contribution[6,33] is the sum of the exchange term, proportional to 2 in Eq. (6), and the Hartree term $\frac{3}{2}\widetilde{F}_\sigma$ where $\widetilde{F}_\sigma=8(1+F/2)\ln(1+F/2)/F-4$, and F is the angular average of the statically screened Coulomb interaction. [The term within parentheses in Eq. (6) is written as $2-2F$ in earlier papers. The old Hartree term $2F$ is decomposed into a singlet term $-\widetilde{F}_\sigma/2$, which is absorbed into the exchange term, and a triplet term $-\frac{3}{2}\widetilde{F}_\sigma$.] Figure 4 shows that the $R_\square$ in sample 1 increases as ln T for decreasing T. The slope of the straight line corresponds to $\Delta\sigma/\ln T=4.04\times10^{-5}\ \Omega^{-1}$, where we used $\Delta\sigma=-\Delta R_\square/R_\square^2$. Comparing with Eq. (5), we find that C equals 3.27 for sample 1. Using this value in Eq. (6) with $\alpha=-\frac{1}{2}$ we find that $(2-\frac{3}{2}\widetilde{F}_\sigma)/2$ equals 3.77 or $\widetilde{F}_\sigma=-1.18$. The value for $\widetilde{F}_\sigma$ comes out negative because coherent backscattering with dominant spin-orbit scattering alone leads to a positive slope for the R-versus-ln T curve. The fact that a negative slope is observed means that Coulomb effects are significant. However, the magnitude of the slope is larger than can be accommodated by existing theory of the Coulomb effects unless the Hartree term is negative. This is in clear contradiction with existing theory. Similarly large values for C are obtained in two III-V heterojunctions. In $Al_xGa_{1-x}As$-GaAs Lin *et al.*[14] find that the factor $1-F$, which is equivalent to $(2-\frac{3}{2}\widetilde{F}_\sigma)/2$ in Eq. (6), varies from 0.5 to 3.2. In a GaSb-InAs-GaSb quantum-well structure with two hole layers and one electron layer Washburn *et al.*[34] find that C equals 4.8.

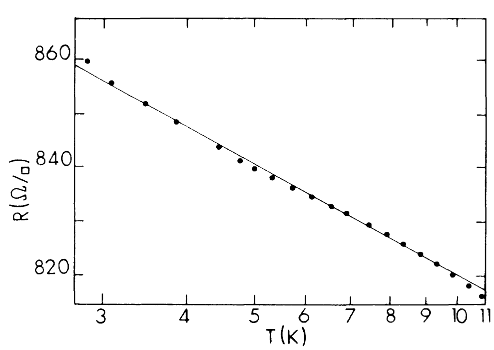

FIG. 4. The areal resistance of sample 1 (per HgTe layer) vs temperature. The slope of the line gives $C=3.27$. [See Eq. (6).]

V. STRONG SPIN-ORBIT SCATTERING IN ZINC-BLENDE CRYSTALS

A general formalism for treating coherent effects (without interactions) in semiconductors such as GaAs, Ge, InSb, and HgTe has been provided by Al'tshuler, Aronov, Larkin, and Khmel'nitskii.[32,4] They considered the Hamiltonian valid for semiconductors with the zinc-blende structure

$$H=\hbar^2k^2/2m^*+\mathbf{s}\cdot\mathbf{\Omega},\ \ \Omega_x=k_x(k_y^2-k_z^2)\delta \ , \quad (7)$$

where Ω is proportional to the spin-orbit scattering rate, s is the spin, and δ is expressed in terms of the electronic mass m^* and energy gap E_g as

$$\delta=\alpha'/[2(2m^{*3}E_g)^{1/2}] \ . \quad (8)$$

To treat spin-scattering mechanisms more generally, Al'tshuler *et al.*[32] resolve the particle-particle propagator ("Cooperon") into the singlet and triplet components ($j=0$ and $j=1$)

$$C(\mathbf{Q})=-\tfrac{1}{2}C^0(\mathbf{Q})+\tfrac{3}{2}C^1(\mathbf{Q}) \ . \quad (9)$$

The quantum correction to the conductivity is proportional to $C(\mathbf{Q})$. In the absence of spin scattering $C^0(\mathbf{Q})$ and $C^1(\mathbf{Q})$ are equal. However, when spin-orbit scattering is domant ($\tau_{\rm s.o.}\sim\tau$) the $j=1$ component is completely suppressed while the $j=0$ component is unaffected by the spin-orbit scattering. This leads to a $\Delta\sigma$ half as large and of the opposite sign as in the weak spin-orbit scattering case, in agreement with Hikami *et al.* [Eq. (1)] in the $\tau_{\rm s.o.}\sim\tau$ limit. The experimental fits and numbers obtained here are in quite good agreement with their calculations. From the magnitude of δ Al'tshuler *et al.* estimate that $\tau_{\rm s.o.}$ in $Al_xGa_{1-x}As$-GaAs heterostructures is between 2×10^{-9} to 2×10^{-10} s. Since this is much longer than τ_ϕ near 1 K no spin-orbit effects are observable at these temperatures in agreement with Ref. 14. (See also Ref. 35.) Note that in HgTe the much smaller magnitudes of m^* and E_g in Eq. (5) greatly enhance the magnitude of Ω and hence $1/\tau_{\rm s.o.}=\tau\Omega^2$.

VI. THE ZEEMAN TERM AND LONGITUDINAL MAGNETORESISTANCE

In addition to the logarithmic anomalies in the MR caused by coherent backscattering there are other causes arising from the Zeeman energy of the spin which can lead to logarithmic anomalies in both the transverse and longitudinal MR. Within interaction theory an important contribution to the magnetoresistance arises from the suppression of Hartree terms to the self-energy when the Zeeman energy (of the electron spins) exceeds the thermal energy or $g\mu_BB>k_BT$, where g is the g factor and μ_B the Bohr magneton. Lee and Ramakrishnan,[36] and Kawabata[36] considered this effect on the Hartree and exchange diagrams in the particle-hole channel. [The applied field $\mathbf{B}$ is assumed to be large enough to suppress the particle-particle (or "Cooperon") channel divergence.] Whereas the exchange terms are unaffected by the Zeeman splitting the Hartree terms proportional to $\tilde{F}_\sigma$ in Eq. (6) with antiparallel spins are suppressed when B greatly exceeds $k_B T/(g\mu_B)$. The suppression leads to a positive MR which is observable in the longitudinal geometry (**B** parallel to current) since it affects only the spins. The field dependence is given by Lee and Ramakrishnan[36] as

$$\delta\sigma(B)-\delta\sigma(0)=-(e^2/4\pi^2\hbar)\tilde{F}_\sigma g_2(h) , \quad (10)$$

$$g_2(h)=\int_0^\infty d\Omega \left[\frac{d^2}{d\Omega^2}[\Omega N(\Omega)]\right]\ln|1-h^2/\Omega^2| , \quad (11)$$

$$h=g\mu_B B/k_B T , \quad (12)$$

where $N(\Omega)=1/[\exp(-\Omega)-1]$. The function g_2 is quadratic in h for $h \ll 1$, and increases as $\ln h$ for $h \gg 1$.

It turns out that in our superlattice samples the longitudinal MR shows a pronounced logarithmic anomaly near zero B which becomes more prominent as T decreases. Although mindful of the inapplicability of Lee and Ramakrishnan's calculations to a material with such dominant spin-orbit scattering we have compared the longitudinal data in Fig. 5 with Eqs. (10)–(12). Sample 3 on which these measurements were performed is from the same batch as the sample used in Ref. 20, and has a lower mobility than sample 1 in Figs. 1 and 2 (8000 compared with 50 000 cm^2/V s). It consists of 155 periods of HgTe and CdTe of layer thicknesses 110 and 330 Å, respectively. As is clear from the data the logarithmic anomaly is superposed on a gentle quadratic background which is insensitive to temperature. The origin of this quadratic background is unclear. We have found that by adding a B^2 term to Eq. (10) the longitudinal MR data are fit very well (solid lines in Fig. 5) at all temperatures from 0.5 to 4.2 K. From these fits we determined the value of the characteristic field B_1 at each temperature where $B_1 \equiv B/h$. The temperature dependence of B_1 is plotted in Fig. 6. In contrast with the expectation that B_1 should scale linearly with T we find that it is highly nonlinear and appears to saturate at high T. The Zeeman interaction theory is clearly inapplicable to the present system. Nonetheless, it might be hoped that the characteristic field B_1 extracted from the fit to the data provides a convenient field scale which summarizes the temperature variation of the longitudinal MR anomaly. Al'tshuler and Aronov[37] point out that in films with finite thickness weak-localization effects alone (without interactions) will give rise to a longitudinal MR of the form (with no spin-orbit or spin scattering)

$$\sigma(B)-\sigma(0)=(e^2/2\pi^2)\ln(1+\tau_\phi/\tau_B) , \quad (13)$$

$$1/\tau_B=4e^2DB^2a^2/12 , \quad (14)$$

where a is the film thickness. As in the transverse case the negative MR is due to suppression of the Cooperon by weak magnetic fields. However, note that at large fields $\delta\sigma$ increases as $\ln(B/B_2)$ where $B_2 \equiv (12B_\phi\hbar/a^2e)^{1/2}$. Thus the ratio of longitudinal and transversal fields needed to suppress the logarithmic anomaly B_2/B_ϕ equals $7l_\phi/a$ which is $\gg 1$. The inelastic length l_ϕ equals $(D\tau_\phi)^{1/2}$. If we make a heuristic generalization to the strong spin-orbit scattering case (not treated in Ref. 37) the ratio $B_2/B_{\text{s.o.}}$ equals $7l_{\text{s.o.}}/a$ where the diffusion length between spin-orbit scatterings $l_{\text{s.o.}}$ is 0.3 μm for sample 1. Using $a=110$ Å we find that $B_2/B_{\text{s.o.}}$ equals 220 which is roughly of the right order of magnitude. (From the data for sample 1 in Figs. 1 and 4, respectively, $B_{\text{s.o.}}=1.7$ mT and $B_2 \sim B_1=0.15$ T.) A reliable comparison is not possible because of the large mobility difference between the high mobility samples (1 and 2) and sample 3. A more systematic study is underway to compare the longitudinal MR in higher mobility samples to the weak-localization theory.

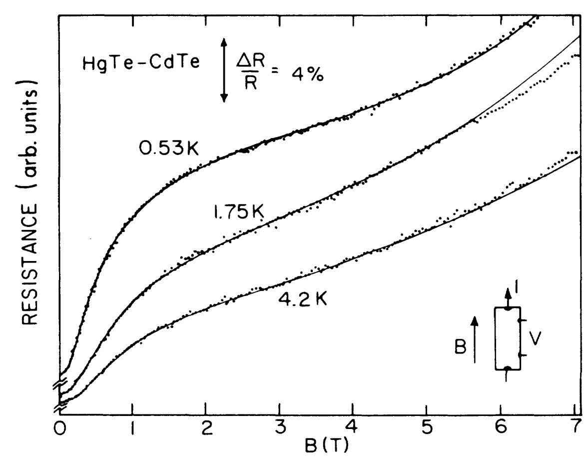

FIG. 5. The longitudinal magnetoresistance of HgTe-CdTe superlattice (sample 3) at various temperatures. The inset shows the field orientation. The sample has 155 periods with a carrier mobility of 8000 cm^2/V s. The solid lines are fits to Eq. (10) (plus a quadratic background).

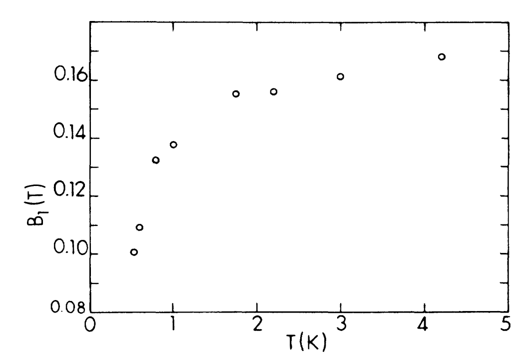

FIG. 6. The temperature dependence of the characteristic field $B_1=B/h$ [where h is defined in Eq. (12)] extracted from data partially displayed in Fig. 5. The interaction theory calculations (Ref. 36) (Zeeman effect) predict that B_1 is linear in T.

VII. HIGH-RESISTANCE SAMPLES

The results above are for samples in which the metallic parameter $\varepsilon_F\tau/\hbar$ is very large so that perturbation-theory calculations can be compared with experiment with confidence. We have also studied a number of samples with much larger areal resistance. These samples (usually single layers of HgTe grown on a CdTe substrate) have lower mobility and show rather complicated MR behavior. As an example we show in Fig. 7 data from sample 4 which consists of a single layer of HgTe of thickness 70 Å. The areal conductivity $1/R_\square$ decreases as $\ln T$ for T below 20 K with a coefficient C equal to 1.11 [using Eq. (6)]. However, since the areal resistance is 30 k$\Omega/\square$ at 2 K (9 k$\Omega/\square$ at 50 K) we are clearly out of the realm in which perturbation theory is valid. It is interesting that logarithmic variation with C close to 1 persists in this sample where localization cannot be considered incipient. The transverse MR has been measured at various temperatures at fields up to 15 T (Fig. 8). Although superficially similar to Fig. 1 we are in a regime where perturbation theory[9] of the magnetoresistance is invalid. Some of the difficulties we faced in interpreting the data in Fig. 8 may be brought out by the following discussion. At 0.4 K the MR is positive below 2 T. Above this field it becomes strongly negative before turning positive again. Within perturbation theory the usual interpretation[8,9] is that below 2 T the diffusion length $l_{\text{s.o.}}$ between scattering events which rotate spin (by spin-orbit coupling) is much shorter than the magnetic length l_B which is in turn shorter than the dephasing length l_ϕ. Thus the carriers are antilocalized and the increasing field suppresses the interference, resulting in positive MR. Above 2 T, l_B becomes shorter than $l_{\text{s.o.}}$ so that within the coherent region smaller in area than πl_B^2 the scattering events do not rotate spin. Thus the carriers are localized and field suppression of the localization leads to negative MR. (This also implies that $1/\tau \gg 1/\tau_{\text{s.o.}}$.) Finally, the upturn at 10 T is identified with the classical effect $\sim(\mu B)^2$. However, on closer examination of the trends versus temperature we find that this interpretation is invalid. As T increases l_ϕ becomes shorter than $l_{\text{s.o.}}$. Within a coherent region spin-orbit scattering is unimportant. Thus the MR is negative at weak fields. In other words as T increases the central anomaly (positive MR) should disappear much faster than the outer wings associated with negative MR. (A clear exposition of these trends is provided by Bergmann.[8]) The data in Fig. 8 indicate just the opposite behavior. The central anomaly persists up to 4.2 K, while the negative MR at higher fields vanishes as T is raised from 0.4 to 4.2 K. Theoretical understanding of the regime beyond weak localization has not been as successful as in the weak regime. It is hoped that further experiments on the higher-resistance samples may lead to further progress.

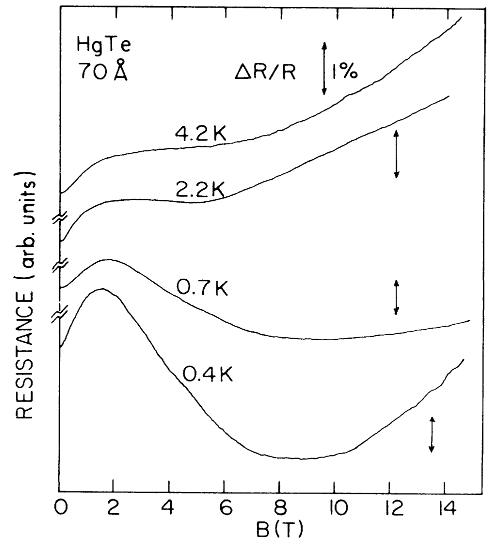

FIG. 8. The transverse magnetoresistance of sample 4 at various temperatures. The arrows indicate a 1% variation of each T. Despite the similarity to the data in Fig. 1 at low fields the data are inconsistent with the predictions of weak-localization theory. Note that as T increases the negative magnetoresistance contribution is strongly suppressed whereas the anomaly near zero field is relatively unaffected. (See text.)

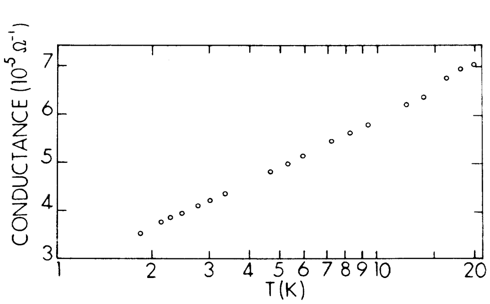

FIG. 7. The conductance of a high-resistance HgTe-CdTe heterojunction (sample 4) versus temperature. The areal resistance increases from 9 kΩ at 50 K to 30 kΩ at 2 K. This sample is not in the weak localization regime although the conductance varies as $\ln T$.

VIII. DISCUSSION

The successful fabrication of superlattices and 2D heterostructures of the HgTe-CdTe and $Hg_{1-x}Cd_xTe$-CdTe family has provided a rather rich testing ground for the many predictions of the localization and interaction calculations that have been reported. The high mobility samples (μ exceeding 50 000 $cm^2/V\,s$) discussed here also show[20,25] pronounced quantized Hall plateaus at high fields. In general, the interpretation of all transport measurements on these superlattices is hampered by uncertainties of the electronic band structure[38,39] in the superlattices. Nonetheless, we find that the low-T transverse MR anomalies in the high mobility samples are rather well described by calculations based on coherent backscattering. From the close fits to the data we have derived numbers for $\tau_{\text{s.o.}}$ and τ_ϕ which are quite reasonable and in good agreement with weak-localization theory. The parameter α_B is found to be 0.3–0.4 compared with 1 predicted by theory. Both the mean free path and the spin-orbit diffusion length are of the order of 0.3 μm, whereas the dephasing length is 3 μm at 1 K. These long-length scales suggest the feasibility of experiments which exploit the large coherent area in these samples at low T. The interaction effects are more difficult to determine quantitatively. Because C [defined in Eq. (6)] remains positive despite the dominant spin-orbit scattering, we infer that interaction effects must be quite large in this system. The large value of C (negative $\tilde{F}_\sigma$) is inconsistent with existing calculations, a situation also found in GaAs-$Al_xGa_{1-x}As$ and GaSb-InAs-GaSb heterostructures. This discrepancy is the most glaring failure of the existing interaction theory. Hopefully, with the coherent backscattering effects well accounted for in this system one may be able to isolate the remaining prominent interaction effects, and thereby confront perturbation calculations more critically. We have attempted to analyze the longitudinal MR using Lee-Ramakrishnan-Kawabata theory. However, the T dependence of the characteristic field shows that the application is not valid, a conclusion consistent with the very strong spin scattering deduced from the transverse MR analysis. We propose that the data are more consistent with the suppression of the particle-particle channel (weak localization effects) by the longitudinal field. Unfortunately, a calculation for the strong spin-orbit scattering case is not available. We have also displayed data form a high-resistance sample as a basis for discussing the difficulties encountered in extending the coherent backscattering arguments to the strongly localized regime.

Because of the dominant spin-orbit scattering and the very high diffusion constant of the carriers 2D systems based on HgTe and the alloy $Hg_{1-x}Cd_xTe$ present a very interesting alternative to the III-V compounds for testing the ideas of quantum transport. Aside from the magnetoresistance studies reported here the use of other probes such as tunneling may prove helpful in disentangling the interaction effects.

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*Present address: Rockwell International Science Center, P.O. Box 1085, Thousand Oaks, CA 91360.

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