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LETTER TO THE EDITOR

Low-temperature transport in Si:Sb ultra-thin doping layers

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Abstract. We present the results of low-temperature transport measurements on 'metallic' Si:Sb doping layers with nominal widths of 10 nm, 20 nm and 80 nm. There is clear evidence of weak localization and interaction corrections to the transport coefficients and for a 3D to 2D transition as the layer width decreases—the first observation of these effects in this system. Of particular interest is the fact that the experimental results are in good quantitative agreement with most aspects of the theory, despite the highly disordered and multi-sub-band nature of the samples.

Much effort is currently being invested in trying to understand the nature of transport in ultra-thin doping layers in semiconductors [1]. In such systems, which include the so-called delta-doped structures, the dopants (either donors or acceptors) are placed in a very narrow layer such that electrostatic confinement can result in quasi-two-dimensional (2D) behaviour being observed. The principal differences between this type of 2D system and the more conventional $GaAs/Al_xGa_{1-x}As$ heterostructures or Si MOSFETs are (typically) (i) enhanced impurity scattering (transport occurs in the plane of the dopants), (ii) greater carrier concentrations and (iii) many occupied sub-bands.

In this letter we present low-temperature transport measurements taken on three n-type Si:Sb layers with nominal doping widths of 10 nm, 20 nm and 80 nm. The Sb doping profiles were realized by coevaporation doping in a VG80 molecular beam epitaxy system at a growth temperature of 800 °C. During growth, the effective 3D doping density, N_B for each sample was kept at about 1×10^{19} cm^{-3} and the background doping was made p-type ($\sim 10^{16}$ cm^{-3}) by coevaporation of boron. In bulk Si the conduction band has a six-fold valley degeneracy ($g_v = 6$). In the ultra-thin doping layer, two of these valleys have the heavy, longitudinal mass (0.916 m_0) perpendicular to the growth surface, and they form the lowest-lying sub-bands (whose energy origins are labelled E_i). The other four have the lighter, transverse mass (0.190 m_0) perpendicular to the growth surface and these form another set of sub-bands (labelled E_i'). Assuming a simple square-well confining potential, the measured sheet carrier concentration, n_s, of the 10 nm sample (9.3×10^{12} cm^{-2}) implies that three sub-bands are occupied (E_0 ($g_v = 2$), E_1 ($g_v = 2$) and E_0' ($g_v = 4$)) and that the Fermi energy $\epsilon_F \simeq 22$ meV. By the same reasoning the 20 nm sample has six occupied sub-bands and the 80 nm sample has 23 occupied sub-bands (all three samples have $\epsilon_F \simeq 20$ meV). The precise number of occupied sub-bands is obviously sensitive to

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the form of the confining potential; nevertheless, as we shall argue below, the exact nature of the sub-band structure is relatively unimportant in our experiments.

Below 4 K, the samples have typical mobilities μ of only ~100–200 cm^2 V^{-1} s^{-1} and are heavily disordered in the sense of $\epsilon_F\tau/\hbar \simeq 1$ (where τ is the transport time). However, the samples are clearly metallic since the behaviour observed at low magnetic fields is quite distinct from the variable-range hopping seen in more lightly doped delta layers [2]. The dominant scattering mechanism is due to ionized donors, and if we estimate τ using the theoretical expression for a square well of width W [3, 4],

$$1/\tau = (\pi\epsilon_F N_B W/g_v n_s) F_B(2k_F) / \left\{ [1 - G(2k_F)] F_C(2k_F) + 2k_F/q_s \right\}^2 \tag{1}$$

(where q_s is the Fermi screening wavenumber and F_B, F_C and G are the impurity form-factor, the electron form-factor and the Hubbard local field correction respectively), then we can readily explain the value of μ; i.e. ~10^2 cm^2 V^{-1} s^{-1}. Equation (1) has not been derived for our particular experimental system, but it probably provides a reasonable order of magnitude estimate (note that the form factors and the local field correction are not especially sensitive to well shape [3, 4]).

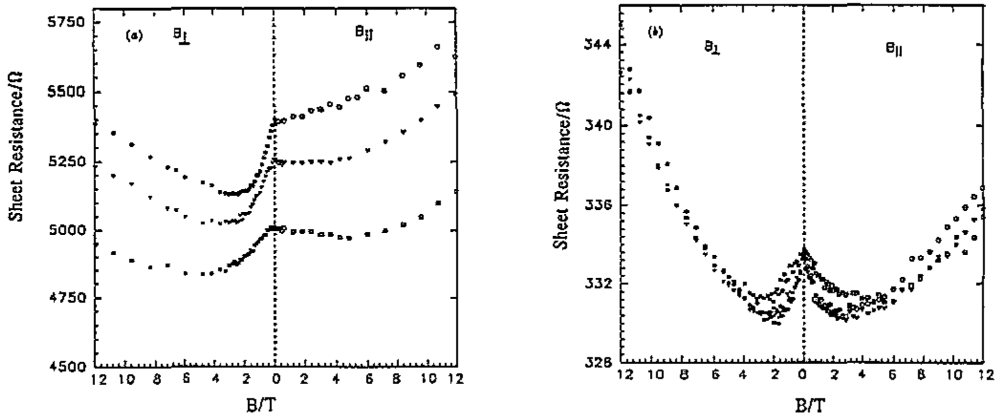

Figure 1. (a) Sheet resistance as a function of magnetic field applied both perpendicular ($B_\perp$) and parallel ($B_\parallel$) to the plane of the 10 nm doping layer (○, ●, 1.4 K; ▽, ▼, 2.0 K; □, ■, 4.0 K); (b) the same as (a) for the 80 nm sample.

In figure 1 we show the magnetoresistance (MR) of the 10 nm sample and the 80 nm sample at three different temperatures and in fields of up to 12 T, applied both perpendicular and parallel to the doping plane. The MR of the 10 nm sample is highly anisotropic, with almost no low-field ($\leqslant$ 2 T) negative MR observed in the parallel field direction. This is consistent with quasi-2D transport, the negative MR resulting from the suppression of quantum interference by the field. If we take the usual expression for the 2D weak localization (WL) correction to the conductivity [5],

$$\delta\sigma(B,T) = \left(g_v\alpha e^2/2\pi^2\hbar\right)\left[\Psi\left\{\tfrac{1}{2} + \hbar/4eBL_\varphi^2\right\} - \Psi\left(\tfrac{1}{2} + \hbar/4eBl^2\right)\right] \tag{2}$$

where $\Psi(x)$ is the digamma function, L_φ is the phase relaxation length, l is the elastic scattering length and, for the moment, we assume that all effects relating to valley

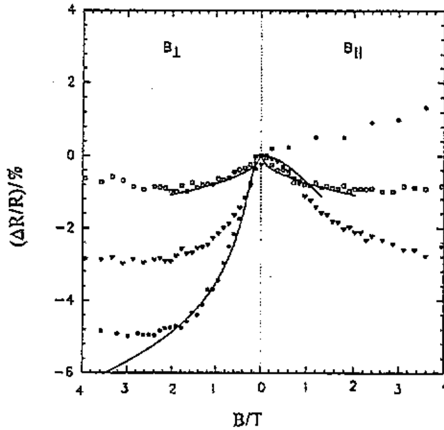

Figure 2. Fractional change in resistance as a function of magnetic field for all three samples (●, 10 nm; ▽, 20 nm; □, 80 nm). The full curves show the theoretical predictions.

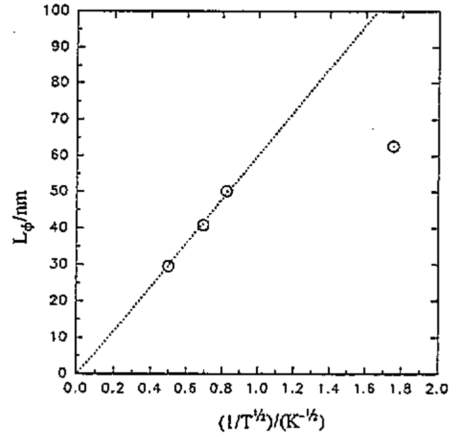

Figure 3. The phase relaxation length L_φ plotted against $T^{-1/2}$.

degeneracy and multiple sub-band occupancy are contained in the phenomenological factor $g_v\alpha \simeq 1$ (see below), then from fitting to the data we extract a value of $L_\varphi(1.4\ \mathrm{K}) \simeq 34$ nm. The criterion for the validity of equation (2) is that $L_\varphi > W$ [5], which is satisfied. Conversely, for the 80 nm sample $L_\varphi < W$ (theoretically L_φ should not be a strong function of W), and so a 3D WL correction is to be expected [6],

$$\delta\sigma(B) = 0.605 g_v \alpha e^2 W/2\pi^2\hbar L_c \tag{3}$$

where $L_c = (\hbar/eB)^{1/2}$ is the cyclotron length and equation (3) has been scaled by W to compare with the measured quantity. Both qualitative and quantitative agreement are found, with a $B^{1/2}$ scaling at low fields and an almost isotropic low-field MR. On the basis of this, one would expect the 20 nm sample to be close to the critical width governing the 2D to 3D transition in the WL correction. In figure 2 we compare the low-field MR for all three samples. A detailed modelling of the 20 nm sample is complicated by the fact that the three length scales, W, L_φ and L_c, are all roughly comparable at $B \simeq 1$ T. Perhaps the most appropriate formula is that derived by Altshuler and Aronov [7] for the MR of a thin film in a parallel field,

$$\delta\sigma(B,T) = \left(g_v\alpha e^2/2\pi^2\hbar\right)\ln\left(1 + W^2L_\varphi^2/3L_c^4\right). \tag{4}$$

Taking the previous estimate for $L_\varphi(1.3\ \mathrm{K}) \simeq 34$ nm we find that $W \simeq 25$ nm, which suggests that assuming a square-well confining potential is not that unreasonable. Moreover, a fit to values of L_φ at different temperatures (taking $W \simeq 20$ nm) suggests that $L_\varphi = (D\tau_\varphi)^{1/2} \propto T^{-1/2}$, where D is the diffusion constant and τ_φ is the phase relaxation time, or $\tau_\varphi \propto T^{-p}$ where the exponent $p \simeq 1$ (figure 3). This is in accordance with theory [5] which gives values of $p \simeq 1$ in 2D and $p \simeq 3/2$ in 3D. It is interesting to reflect that in p-type Si:B delta layers the weak-field negative MR observed here is absent, due to the strong spin–orbit scattering found in p-type Si:B samples [8, 9].

For fields greater than about 4 T the MR becomes positive in all three samples for both perpendicular *and* parallel field orientations (see figure 1). The most immediate explanation for this is Zeeman splitting of the spin-up and spin-down bands affecting the electron–electron interaction contribution to the conductivity, an effect known to lead to positive MR [5]. Evidence for the presence of interaction effects is clearly found by studying the temperature dependence of the resistivity R and Hall coefficient R_H (see figures 4 and 5). For the 10 nm sample the pronounced logarithmic dependence observed is entirely in accordance with 2D theory, which predicts that [5]

$$\delta R/R \simeq -R(e^2/2\pi^2\hbar)\left(1-\tfrac{3}{4}F^* + g_v\alpha p\right)\ln T \tag{5a}$$

$$\delta R_H/R_H \simeq -R(e^2/2\pi^2\hbar)\left(1-\tfrac{3}{4}F^*\right)\ln T \tag{5b}$$

where F^* is an interaction screening parameter whose theoretical value should be of order unity [10] (note that the Hall coefficient is sensitive only to interaction effects). From the data we extract values of $F^* \simeq 1$ and $g_v\alpha p \simeq 0.86$ (so taking our previously determined value of $p \simeq 1$, we have $g_v\alpha \simeq 0.86$). The dimensionality of the interaction contribution is 2D as long as the thermal diffusion length $L_T = (\hbar D/kT)^{1/2}$ is greater than the width W [5]. At 1.4 K, we estimate that $L_T \simeq 40$ nm. The relative weakness of the correction to R_H in the 80 nm sample is therefore to be expected (figure 5). According to theory [5], in 3D the interaction correction to the conductivity varies as $mT^{1/2}$ and the WL correction as $BT^{p/2}$, where $p \simeq 3/2$. Fitting to the data we extract values of $m \simeq -8\ \Omega^{-1}\ \text{cm}^{-1}\ \text{K}^{-1/2}$ and $B \simeq +7\ \Omega^{-1}\ \text{cm}^{-1}\ \text{K}^{-3/4}$, in sensible agreement with other workers [11]. The finite values of F^* and m confirm the importance of interaction effects in these samples.

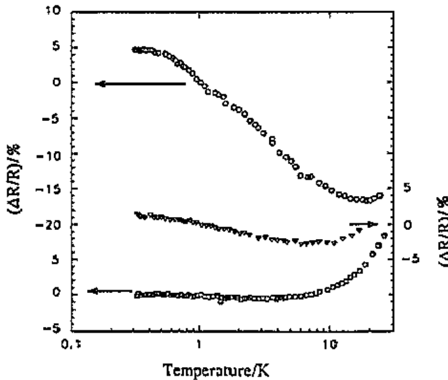

Figure 4. Fractional change in sheet resistance plotted against temperature for all three samples (○, 10 nm; ▽, 20 nm; □, 80 nm).

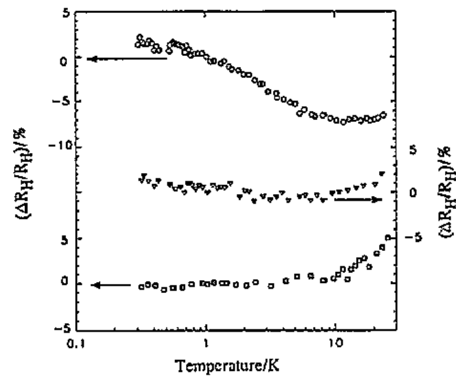

Figure 5. Fractional change in Hall coefficient plotted against temperature for all three samples (○, 10 nm; ▽, 20 nm; □, 80 nm).

The Zeeman splitting contribution to the MR mentioned above has been calculated for both 2D and 3D situations [5]. It is predicted to be independent of field orientation and to become significant once the Zeeman splitting energy $g\mu_B B$ exceeds kT (where g is the electronic g factor and μ_B is the Bohr magneton). Assuming a value of $g \simeq 2$,

then this means $B \geqslant 1.2$ T at 1.4 K. At high fields, the MR is expected to vary as $F^* \ln(B/T)$ in 2D and $F^* B^{1/2}$ in 3D [5]. At a qualitative level, the data support the theory (see figure 1): the high-field MR is roughly isotropic; the effect becomes weaker with increasing T in the 2D (10 nm) sample and is approximately independent of T in the 3D (80 nm) sample. However, serious deficiencies within the theoretical predictions are also apparent. In particular, the currently accepted theory is unable to account for the magnitude of positive MR (without invoking unacceptably high values of F^*) nor its scaling with B, in either 2D or 3D. These difficulties have been noted many times before in highly disordered samples which are close to the metal–insulator transition, and they suggest that the theory is not adequate to describe such situations. On the insulating side of the transition, theories predict that the resistance R varies as $R_0 \exp[(T_0/T)^\nu]$, where the 'Mott' parameter $T_0 \propto B$ and the exponent ν depends upon the assumptions being made (such as whether there is a Coulomb gap present or not) [12–14]. It is possible that in the present case (as we have suggested elsewhere for Si:B delta layers [9]) the samples are close to undergoing a transition to an insulating state once the magnetic field exceeds a certain value and so exhibit high-field behaviour more reminiscent of hopping conduction. A detailed theory of transport in this interesting transition regime is, however, unfortunately still unavailable.

On the face of it, it is puzzling why the simple theoretical expressions quoted above lead to any reasonable answers at all for the present samples. One important simplifying feature is the significant inter-valley ($1/\tau_v$) and inter-sub-band ($1/\tau_{ij}$) scattering rates due to the heavy disorder. Clearly, if no such scattering took place then each band and valley would contribute additively. However, it has been shown (by considering the effect of inter-valley scattering on WL in isolation) that when $1/\tau_v \gg 1/\tau_\varphi$, the parameter $g_v\alpha$ (which equals g_v in the limit of no inter-valley scattering) tends towards unity, i.e. the system behaves as if it consisted of essentially a single valley [15]. Similar results have been shown for the case $1/\tau_{ij} \gg 1/\tau_\varphi$, when the system behaves as if it consisted of a single 'effective' band [16, 17]. This may go some way towards explaining the success of our simplified analysis.

In conclusion, we have presented the first detailed measurements of localization and interaction corrections in Si:Sb ultra-thin doping layers. Despite the multiple sub-band occupancy and the valley degeneracy, excellent agreement is found between experiments and theories derived for much simpler systems, apart from the high-field magnetoresistance, which remains a topic of great theoretical interest.

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