Observation of orientation- and $k$-dependent Zeeman spin-splitting in hole quantum wires on (100)-oriented AlGaAs/GaAs heterostructures

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Abstract. In this paper, we study the Zeeman spin-splitting in hole quantum wires oriented along the [011] and [01\bar{1}] crystallographic axes of a high mobility undoped (100)-oriented AlGaAs/GaAs heterostructure. Our data show that the spin-splitting can be switched ‘on’ (finite $g^*$) or ‘off’ (zero $g^*$) by rotating the field from a parallel to a perpendicular orientation with respect to the wire, and the properties of the wire are identical for the two orientations with respect to the crystallographic axes. We also find that the $g$-factor in the parallel orientation decreases as the wire is narrowed. This is in contrast to electron quantum wires, where the $g$-factor is enhanced by exchange effects as the wire is narrowed. This is evidence for a $k$-dependent Zeeman splitting that arises from the spin-$\frac{3}{2}$ nature of holes.

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1. Introduction

The use of spin instead of charge to carry information is a central goal in the fields of spintronics and quantum information, generating significant interest in routes to efficient spin manipulation in semiconductor devices [1, 2]. Low-dimensional hole systems in p-type AlGaAs/GaAs heterostructures hold considerable potential because the much stronger spin–orbit coupling in holes [3] may lead to devices where spin can be manipulated electrostatically [4, 5]. The strong spin–orbit coupling also presents some important fundamental physics questions, including how the peculiar spin-$\frac{3}{2}$ nature of holes [6] is manifested in the experimentally observable properties of low-dimensional GaAs hole devices [7]–[17].

Experiments to date have focused almost solely on devices fabricated in (311)-oriented AlGaAs/GaAs heterostructures. The Zeeman spin-splitting in two-dimensional (2D) hole systems formed in these heterostructures is highly anisotropic [7, 8], owing to spin–orbit coupling and the low symmetry of the (311) surface. Recent studies have also revealed a significant anisotropy in the Zeeman spin-splitting in one-dimensional (1D) hole systems fabricated on the (311) heterostructures [15]–[17], but it is not trivial to separate the competing influences of 1D confinement and 2D crystallographic anisotropy on the spin-splitting [17]. Hole systems fabricated on higher-symmetry planes such as (100) are not subject to such complex crystallographic effects, and are therefore a much better candidate for studying the spin physics of 1D hole systems. To achieve high-quality 1D hole systems, we use semiconductor–insulator–semiconductor field-effect transistor (SISFET) devices, where a 2D hole system is ‘induced’ using a voltage applied to a degenerately doped semiconductor gate rather than through modulation doping [18, 19]. Klochan et al have used this approach to fabricate 1D hole systems with highly stable gate characteristics and clear conductance quantization [20], and recently extended it to study the Zeeman spin-splitting anisotropy in 1D hole systems in (311)-oriented heterostructures [17].

In this paper, we extend this SISFET-based approach to study the Zeeman spin-splitting in hole quantum wires oriented along the [011]- and [01T]-directions of a (100)-oriented heterostructure. The crystallographic anisotropy that complicates transport studies of quantum wires on (311)-oriented heterostructures [16, 17] does not occur in these devices. Instead, we
find that the Zeeman spin-splitting is finite when the applied magnetic field $B$ is oriented parallel to the wire, and nearly zero when $B$ is oriented perpendicular to the wire. This behaviour is almost identical for both orientations of the wire relative to the dominant in-plane crystallographic directions. The ability to switch the spin-splitting ‘on’ or ‘off’ simply by rotating the applied magnetic field through 90° may have useful spintronic applications. Finally, for $B$ parallel to the wire, we observe $k$-dependent spin-splitting, where $g^*$ decreases as the wire is made narrower, in marked contrast to 1D electron systems, where $g^*$ instead increases as the wire becomes more 1D [21, 22]. This finding is reminiscent of the absence of exchange enhancement effects for 2D hole systems in (100)-oriented heterostructures [10].

2. Experimental details

Samples were fabricated from a (100)-oriented heterostructure that consisted of a heavily doped 20 nm C: GaAs cap, a 10 nm undoped GaAs, an 160 nm undoped AlGaAs barrier and an undoped GaAs buffer. The C-doped cap acts as a metallic gate [19], with a 2D hole system induced at the AlGaAs/GaAs interface for top-gate voltages $V_{TG} < -0.1$ V. Measurements of a separate unpatterned Hall bar of the same heterostructure gave a peak mobility $\mu = 4.8 \times 10^5$ cm$^2$/Vs at a density $p = 1.3 \times 10^{11}$ cm$^{-2}$ and a temperature $T = 100$ mK. The device studied here consists of two orthogonal 400 nm long quantum wires, as shown in figure 1 (inset), defined by EBL and shallow wet etching of the cap layer. Each wire has three gates, a top gate used to control the density, and two side gates used to narrow the wire. The two wires can be measured independently and are oriented along the [011] and [01$\bar{1}$] crystallographic directions of a Hall bar running along the [01$\bar{1}$] direction. The two wires are denoted as QW011 and QW01$\bar{1}$, respectively. The quantum wires were measured in a dilution refrigerator with
a base temperature of 20 mK using standard ac lock-in techniques with an excitation voltage of 50–100 µV at a frequency of 17 Hz. Measurements were obtained at a top-gate voltage $V_{TG}$ = −0.80 V, which corresponds to a 2D hole density of $2.56 \times 10^{11}$ cm$^{-2}$.

The width of the wire and its conductance $G$ can be gradually reduced by applying a positive voltage $V_{SG}$ to the two side gates. For both wires, we observe the well-known ‘staircase’ of quantized conductance plateaus [23] as the wire is narrowed by increasing $V_{SG}$, with the wire ‘pinching off’ at $V_{SG}$ $\sim$ 1.5 V. The similar pinch-off voltages indicate that the two wires are almost identical, with similar dimensions and confining potentials. The accurate quantization of the plateaus at $G = n \times 2e^2/h$, where $n$ is the number of occupied 1D subbands, confirms that transport through the wires is ballistic [20, 23]. Moving from left to right in figure 1 corresponds to strengthening the 1D confinement, taking the wire from being only quasi-1D (large $n$ and $G$) towards the 1D limit (small $n$ and $G$).

We study the spin properties of the hole quantum wires by measuring the Zeeman spin-splitting for different orientations of the wire and magnetic field with respect to the crystallographic axes. To obtain the $g$-factor for the various 1D subbands $n$, we use a technique that compares the 1D subband splitting due to an in-plane magnetic field [24] (see figure 2) and an applied dc source–drain bias [25] (see figure 3). These two sets of measurements are repeated in two cool-downs to allow for rotation of the sample with respect to the magnetic field, thus providing data for the four different combinations of wire and magnetic field orientation with respect to the crystallographic axes.

3. Results

3.1. 1D subband spacings and source–drain bias measurements

The 1D subband spacing of the wires is obtained by adding a dc bias $V_{SD}$ to the ac bias used to measure the conductance. In figure 2(a), we plot the transconductance $d g / d V_{SG}$, where $g = dI / dV$ is the differential conductance, as a colour-map against $V_{SG}$ and $V_{SD}$ using data obtained from QW01T. Figure 2(b) shows the conductance $G$ versus $V_{SG}$ measured at $V_{SD} = 0$ V and corresponds to taking a vertical slice through the centre of the colour map in figure 2(a). The dark regions in figure 2(a) correspond to high transconductance (risers between plateaus) and the bright regions correspond to low transconductance (the plateaus themselves). Thus, the dark regions indicate when a particular 1D subband crosses the Fermi energy. As $V_{SD}$ is increased, the plateaus at multiples of $2e^2/h$ evolve into plateaus at odd multiples of $e^2/h$. The subband spacing $\Delta E_{n,n+1} = eV_{SD}$ is obtained from the source–drain bias where adjacent transconductance peaks cross (i.e. from the dark regions at nonzero $V_{SD}$). The subband spacings for the two wires are plotted in figure 2(c), and increase monotonically from $\sim$100 to $\sim$300 µeV as the wire is made narrower and more 1D. The subband spacings for the two wires agree within 10 µeV, again highlighting the similarity of the two wires fabricated along different crystallographic axes.

3.2. Zeeman spin-splitting measurements

The effect of an in-plane magnetic field $B$ on the 1D subbands is shown in figures 3(a)–(d) for different orientations of the quantum wire and magnetic field. In each case, we plot a colour map of the transconductance $d g / d V_{SG}$ versus $B$ and $V_{SG}$, with the dark regions marking the
Figure 2. (a) The conductance $G$ versus $V_{SG}$ measured at $V_{SD} = 0$ V, which corresponds to a vertical slice through the centre of the colour map in (b). (b) A colour map of the transconductance $d_g/dV_{SG}$ versus $V_{SD}$ on the $x$-axis and $V_{SG}$ on the $y$-axis for QW01T. The bright regions in (b) correspond to conductance plateaus (low transconductance) and dark regions correspond to the risers between plateaus (high transconductance). The superimposed numbers in (b) indicate the conductance $G$ of the corresponding plateau in (a) in units of $2e^2/h$. (c) The measured 1D subband energy spacings $\Delta E_{n,n+1}$ obtained from the subband crossings plotted as a function of 1D subband index $n$.

1D subband edges (high transconductance corresponding to the risers between conductance plateaus). The superimposed white dashed lines in figure 3 are guides to the eye, tracking the evolution of the 1D subbands with $B$.

Figure 3 shows that there is only a Zeeman splitting of the 1D subbands if $B$ is aligned along the wire, independent of the crystallographic orientation of the wire: if the field is aligned perpendicular to the wire, as in figures 3(a) and (d), then the Zeeman spin-splitting is extremely weak. In figure 3(d), no splitting is evident up to the highest fields available in the experiment $B = 10$ T, while in figure 3(a), some splitting is just apparent near $B \sim 10$ T. In stark contrast, if $B$ is aligned parallel to the wire, as in figures 3(b) and (c), then the Zeeman spin-splitting is quite strong with clear splitting evident at quite modest fields $B \sim 1$ T, crossings between adjacent subbands at moderate fields $B \sim 5$ T, and ultimately, crossings between subbands differing in $n$ by two at high fields $B \sim 10$ T. The directional dependence of the Zeeman spin-splitting in these (100)-oriented quantum wires is much simpler than in wires fabricated on (311)-oriented heterostructures, where a complex interplay between 1D confinement and 2D crystallographic anisotropy is observed [12, 16, 17].

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3.3. Obtaining the $g$-factors for the four magnetic field and wire orientations

We now extract the effective Landé $g$-factors [12, 17]. When $g^*$ is large enough that the spin-down level of the $n$th subband crosses the spin-up level of the $(n+1)$th subband within the maximum available magnetic field $B = 10$ T (see figure 3), then $g^*$ can be obtained from a measurement of the field $B^C(n)$ at which this crossing occurs. Combining this crossing field with the corresponding dc bias $V^C_{SD}$, where the $n$ and $(n+1)$th subbands cross in figure 2, gives

$$
\langle g^*_n, g^*_{n+1} \rangle = \frac{eV^C_{SD}}{\mu_B B^C}.
$$

Data obtained in this way are plotted as solid symbols at $(n + 1)/2$ in figure 4, since the quantity $\langle g^*_n, g^*_{n+1} \rangle$ represents the average $g$-factor for the two subbands.

When the spin-splitting is small, as in figures 3(b) and (d), we can only measure an upper bound on $g^*$, i.e. $g^*$ must sit between zero and this upper bound; otherwise the spin-splitting would be resolvable. We determine this upper bound from the width $\Delta V^C_{SG}$ of the transconductance peak in the colour map at $B_{\text{max}} = 10$ T, which would be the maximum possible

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**Figure 3.** Colour maps of the transconductance $dg/dV_{SG}$ versus in-plane magnetic field $B$ (x-axis) and $V_{SG}$ (y-axis) for QW011 with (a) $B \parallel [011]$ (field perpendicular) and (b) $B \parallel [01\bar{1}]$ (field parallel), and QW011 with (c) $B \parallel [011]$ (field parallel) and (d) $B \parallel [01\bar{1}]$ (field perpendicular). The superimposed numbers in (a) indicate the conductance $G$ of the corresponding plateau in units of $2e^2/h$. The white dashed lines are guides to the eye that track the evolution of the various 1D subbands with $B$.
Figure 4. The effective Landé $g$-factor $g_n^*$ plotted as a function of 1D subband index $n$ for four cases: QW011 with $B \parallel [01\bar{1}]$ (red circles), QW011 with $B \parallel [011]$ (blue squares), QW01$\bar{1}$ with $B \parallel [011]$ (red hatched region) and QW011 with $B \parallel [011]$ (blue hatched region). In the latter two cases, hatching is presented because at best we can determine the upper bound on $g^*$, as minimal spin-splitting is observed up to $B = 10$ T (see text).

splitting if it could be resolved. We convert this width into a splitting rate due to the field $\partial V_{SG}/\partial B = \Delta V_{SG}/B_{\text{max}}$, and combine it with dc biasing data $\partial V_{SG}/\partial V_{SD}$ to obtain the upper bound as

$$|g^*| \leq \frac{e}{\mu_B} \frac{\Delta V_{SG}}{\partial V_{SG}/\partial V_{SD} B_{\text{max}}}.$$  (2)

These upper bounds are indicated by the hatched regions in figure 4.

4. Discussion

In order to discuss the two key results in figure 4, namely the $g$-factor anisotropy and the decrease of $g^*$ as the wire is made narrower, it is first necessary to review some of the complexities of Zeeman splitting in the presence of spin–orbit coupling.

For electrons in free space, an applied magnetic field causes the spins to align along $B$, with a spin splitting $\Delta E = \frac{1}{2} g^* \mu_B B$, where $g^* = 2$. In a solid, deviations from $g^* = 2$ are driven by the spin–orbit interaction [26, 27]; for example, in bulk GaAs $g^* = -0.44$. In addition to changing the magnitude of $g^*$, the projections of spin and orbital angular momenta $L$ and $S$ are no longer good quantum numbers if the spin–orbit coupling becomes strong. Only the total angular momentum $J = L + S$ is conserved, and an applied magnetic field causes $J$ to align along $B$. However, if the electron is in a non-symmetric environment, such as a polar GaAs crystal, a quantum well, or a quantum wire, the quantization axis for $J$ does not automatically align with the applied $B$ and it is rarely possible to find eigenstates of both $B$ and $J$. This complicates the theoretical analysis of spin splitting considerably, since the microscopic details of the host crystal and the confinement due to the quantum well/wire must all be taken into account. Furthermore, the practical definition $g^* = \Delta E/\mu_B B$ generally yields a $B$-dependent
g-factor into which any existing contributions to the Zeeman splitting $\Delta E$ that are nonlinear in $B$ are absorbed. As our experimental technique obtains the $g$-factor at some low but still finite magnetic field, additional modifications due to nonlinear terms could arise. Finally, the effect of the Coulomb interaction on the Zeeman splitting of confined holes is currently an open question.

4.1. Zeeman splitting in 2D and quasi-1D holes

The uppermost valence band in bulk GaAs consists of ‘heavy-hole’ (HH) and ‘light-hole’ (LH) branches that are degenerate at the valence band edge ($k = 0$). Confinement to a quantum well breaks this HH–LH degeneracy, such that only the lowest HH subband ($m = \pm \frac{3}{2}$) is occupied in a 2D hole system. However, a residual HH–LH coupling at finite wavevector not only results in highly non-parabolic bands but also plays a significant role in determining the electronic properties in lower-dimensional hole structures.

In the simplest approximation, the 2D confinement forces the quantization axis for $\mathbf{J}$ to point out of the 2D plane. To lowest order, there is only a spin-splitting of the HH states if $B$ is applied perpendicular to the quantum well, since $\langle \mathbf{B} \cdot \mathbf{J} \rangle = 0$ for in-plane magnetic fields [7]. In practice, however, the cubic crystalline anisotropy terms [6], as well as higher order terms in the in-plane wavevector $k_\parallel$, can result in a finite in-plane Zeeman splitting. For quantum wells on the (311) GaAs surface, the cubic anisotropy terms result in a linear in $B_\parallel$ spin splitting at $k_\parallel = 0$. For (100) oriented quantum wells, the zeroth-order contributions due to cubic crystalline anisotropies are absent [7], but a substantial linear spin-splitting can still be achieved because of LH–HH mixing at $k_\parallel \neq 0$, as discussed in section 7.4 of [6]. Because the Zeeman splitting on (100) surfaces arises from $k_\parallel$ dependent LH–HH mixing, it is hard to define a $g$-factor for (100) 2D holes since $g^*$ must be averaged over all occupied states, and is strongly dependent on carrier density. One of the main advantages of quasi-1D systems compared to 2D systems is the ability to perform energy spectroscopy, and thereby measure the $g$-factor directly [21].

We can predict the expected spin splitting in our quantum wire using a quasi-1D model in which we take the 2D results in [6] and add on the effects of quantization of the transverse wavevector by the 1D confinement. We define the components of the wavevector $\mathbf{k}_l = (k_l, k_l)$ with respect to the axis of the quasi-1D wire, where $k_l$ and $k_\parallel$ are the in-plane wavevector components parallel (longitudinal) and perpendicular (transverse) to the quantum wire. In the experiments on 1D holes, the spin-splitting is measured at the 1D subband edges, where $k_\parallel = 0$. Since $k_\parallel$ is quantized by the lateral 1D confinement, we can express the $g$-factor of the $n$th 1D subband as

$$g^*_{[011], \parallel}(n) = g^*_{[01\overline{1}], \parallel}(n) = 3\gamma_3|\kappa Z_1 - 4\gamma_2 Z_2| k_n^2,$$

$$g^*_{[011], \perp}(n) = g^*_{[01\overline{1}], \perp}(n) = 3\gamma_3|\kappa Z_1 - 4\gamma_2 Z_2| k_n^2. \tag{4}$$

Here $k_n$ is the quantized transverse wavevector $k_\parallel$ of the $n$th 1D subband, and we have rotated the expressions in equation (7.22) of [6] to align along the [011]- and [01\overline{1}]-axes. The subscripts on $g^*$ indicate the direction of the wire relative to the crystal and the field relative to the wire, respectively. Here $\kappa$, $\gamma_2$ and $\gamma_3$ denote Luttinger band-structure parameters ($\kappa$ is the isotropic bulk-hole $g$-factor) [28] and $Z_{1,2}$ are LH–HH coupling terms (see p 147 of [6]).

The first inference that we can draw from equations (3) and (4) is that the $g$-factor for both [011] and [01\overline{1}] quantum wires should exhibit the same anisotropy with respect to the magnetic
field, i.e. \( g^{*}_n / g^{*}_1 \) is the same for both wires. This is evident in figures 3 and 4. For both QW011 and QW01T, \( g^* \) is the same for \( B \) parallel to the wire (see figures 3(b) and (c)) and very small for \( B \) perpendicular to the wire (see figures 3(a) and (d)). This behaviour is quite different from quantum wires on (311) surfaces, where the anisotropy depends both on the orientation of the quantum wire with respect to the magnetic field and on the orientation of the field with respect to the crystal axes [17].

However, the quasi-1D theory disagrees with experiment on whether \( g^{*}_n > g^{*}_1 \) or \( g^{*}_n < g^{*}_1 \). Using expressions for \( Z_{1,2} \) for square quantum wells [6] and GaAs bandstructure parameters, equations (3) and (4) predict \( g^{*}_\perp > g^{*}_\parallel \). The experimental data exhibit exactly the opposite trend, \( g^{*}_\perp < g^{*}_\parallel \), as shown in figure 4. This is a surprising result, and we have repeated our experiment to confirm that this is indeed the case, obtaining identical results (to within 10%). We can only surmise that this discrepancy lies in the dependence of \( Z_{1,2} \) on the quantum-well confinement, as our 2D holes are confined in a triangular potential well at a single heterojunction, not in a square quantum well. Unfortunately \( Z_{1,2} \) are not available for a self-consistent triangular quantum well.

A second conclusion that we can draw from equations (3) and (4) is that in the quasi-1D limit, the \( g \)-factor of the wires should decrease with decreasing \( k_n^2 \). In the 1D constriction, \( k_n \) is given by the difference between the Fermi energy in the 2D reservoirs \( E_{F,2D} \) and the bottom of the 1D saddle-point potential [29]. At large subband index, \( k_n \) approaches \( k_n^{2D} \), and \( g^* \) should saturate as seen in figure 4. At small subband index, the wire becomes narrower, the saddle-point rises up in energy, \( k_n \) decreases and so does \( g^* \). Additionally, the increase in 1D confinement increases the LH–HH separation \( \Delta E_{LH,HH} \), which reduces the magnitude of the higher-order Zeeman terms, and thereby reduces \( g^* \). The decrease in \( g^* \) with decreasing subband index is consistent with the data shown in figure 4, but is different from almost all other studies of 1D systems, where a strong exchange enhancement of \( g^* \) is observed at low subband index [12, 21, 22, 30, 31]. It is also different from previous studies of 1D holes in (311) quantum wells, where the Zeeman splitting is believed to be due to a combination of crystal anisotropies at large \( n \) and re-orientation of the quantization axis for \( J \) at small \( n \).

4.2. Zeeman splitting in the 1D limit

In the quasi-1D description, the 1D confinement is a weak perturbation, so that \( \hat{J} \), the quantization axis for \( J \), remains perpendicular to the 2D system. The lowest-order terms for the spin-splitting are zero, since \( \langle B \cdot \hat{J} \rangle = 0 \), and \( g^* \) is only finite due to the higher order \( k_\parallel \) terms. It is thus interesting to consider what happens in the 1D limit where the wire width becomes equal to the width of the 2D confinement. In this case, \( \hat{J} \) is aligned with the wire axis and the lowest order spin-splitting is large and positive for \( B \) applied along the wire, but is zero for \( B \) perpendicular to the wire. This is consistent with the anisotropy measured in figure 4, where \( g^{*}_n > g^{*}_1 \). If the 1D confinement causes a re-orientation of \( \hat{J} \), then one might expect that \( g^* \) would increase as the system is made more 1D, as seen in previous experiments on (311)-based hole wires [12, 17]. Furthermore, it is predicted theoretically that the sign of \( g^* \) is opposite for wires in the (011) and (01T) orientations for a square 2D confinement, so one would expect the measured \( g \) factors for the two quantum wires to show different behaviours as we go from the quasi-1D to the 1D limit.

Thus, the quasi-1D model can explain the observed dependence of \( g^* \) on \( k_\parallel \), but not the anisotropy of \( g^* \), whereas the 1D-limit model can explain the observed anisotropy of \( g^* \), but not the dependence on \( k_\parallel \), since the latter depends on the quasi-1D model. To be able to resolve this
conundrum it will be essential to perform more detailed calculations in the quasi-1D limit for realistic 2D confining potentials.

5. Conclusion

In conclusion, we have studied the Zeeman spin-splitting in hole quantum wires fabricated in (100)-oriented AlGaAs/GaAs heterostructures, and find two new results: Firstly, if the applied in-plane magnetic field $B$ is aligned along the wire, we see strong spin-splitting, and if it is perpendicular to the wire, then we observe negligible spin-splitting up to $B = 10$ T. This behaviour is independent of the orientation of the wire on the heterostructure surface. Although this latter finding is consistent with theoretical predictions, our finding that the spin-splitting is maximized for $B$ aligned along the wire is at odds with a quasi-1D theory, which predicts maximum splitting instead for $B$ perpendicular to the wire. We propose that the possible solution to this disagreement may reside in the sensitivity of the theoretical calculations on the 2D confining potential—theoretical results have only been obtained for a square potential well so far, whereas the single heterojunction in our device leads to a more triangular confinement. Secondly, we report a decreasing $g^*$ as the 1D confinement is increased, which is at odds with previous experiments of both 1D electron systems in GaAs and InGaAs [21, 22, 31] and 1D hole systems in (311)-oriented GaAs heterostructures [12, 30]. This suggests that despite the strong hole–hole interactions there is no exchange enhancement in our 1D wires, consistent with recent measurements of (100)-oriented 2D hole systems [10]. These results highlight the complex and interesting spin-physics associated with $j = \frac{3}{2}$ hole systems, and suggest that much more theoretical work is needed before we understand the physics of holes, even on ‘simple’ (100) surfaces.

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