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Spin-dependent Anomalies in the Conduction Edge of Semiconductor Quantum Wires

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Anomalies near the conductance threshold of clean semiconductor quantum wires are explained in terms of singlet and triplet resonances of conduction electrons with a single weakly-bound electron in the wire. This is shown to be a universal effect for a wide range of situations in which the effective single-electron confinement is weak. The dependence on gate voltage, source-drain voltage and magnetic field is also explained by this model. Speculation is made as to the possibility that these anomalies are a signature for a new kind of spin-polarised ground-state when the electron density in the wire is low.

1. Introduction

Semiconductor quantum wires can be fabricated with effective wire widths down to a few nanometers; for example, by heteroepitaxial growth on 'v'-groove surfaces [1], epitaxial growth on ridges [2], cleaved edge over-growth [3], etched wires with gating [4], and gated two-dimensional electron gas (2DEG) structures [5, 6]. More recently, there has been considerable interest in carbon nanotubes for which the quantum wire cross-section can approach atomic dimensions. Such structures have potential for opto-electronic applications, such as light-emitting diodes, low-threshold lasers and single-electron devices.

Many groups have now observed conductance steps in all of these various types of quantum wire, following the pioneering work in Refs. [5, 6]. Whilst these experiments are broadly consistent with a simple non-interacting picture [7], there are certain anomalies, some of which are believed to be related to electron-electron interactions and appear to be spin-dependent. In particular, a

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structure is seen in the rising edge of the conductance curve, starting at around $0.7(2e^2/h)$ and merging with the first conductance plateau with increasing energy [8]. This structure, already visible in the early experiments [5], can survive to temperatures of a few degrees and also persists under increasing source-drain bias, even when the conductance plateau has disappeared. Under increasing in-plane magnetic field, the structure moves down, eventually merging with the e^2/h conductance plateau at very high fields. Theoretical work has attempted to explain these observation in various ways, including conductance suppression in a Luttinger liquid with repulsive interaction and disorder [9], local spin-polarized density-functional theory [10] and spin-polarized sub-bands [11]. In some of the more recent experiments, an anomaly is seen at lower energy with conductance around $0.2(2e^2/h)$ [12, 2]. This can also survive to a few degrees, though is less robust than the 0.7 anomaly and is more readily suppressed by a magnetic field [2].

In this paper, we suggest that these anomalies are related to weakly bound states and resonant bound states within the wire. These would arise, for example, if there is a small fluctuation in thickness of the wire in some region giving rise to a weak bulge. If this bulge is very weak then only a single electron will be bound. We may thus regard this system as an 'open' quantum dot in which the bound electron inhibits the transport of conduction electrons via the Coulomb interaction. Near the conduction threshold, there will be a Coulomb blockade and we show below that this also gives rise to a resonance, analogous to that which occurs in the single-electron transistor [13]. This is a generic effect arising from an electron bound in some region of the wire and such binding may arise from a number of sources, which we do not consider explicitly. For example, in addition to a weak thickness fluctuation, a smooth variation in confining potential due to remote gates, contacts and depletion regions could contribute to electron confinement along the wire or gated 2DEG. In this paper we consider only very weak confinement near the conductance threshold for which a single electron is bound. The confinement could even be due to electronic polarisation of the lattice caused by the electron itself in an otherwise perfect wire. In the next Section we introduce the basic model and show that it applies to a number of different situations and is in this sense universal, as are the results which are a consequence of it. This is followed by a detailed analysis of the two-electron approximation in which one electron is weakly bound in the wire and gives rise to spin-dependent scattering of the other, this scattering problem being solved exactly. In Section 4 we then show how the solutions of the scattering problem may be used to determine conductance by an extension of the Landauer-Büttiker formula. This gives excellent agreement with a number of experiments on different kinds of quantum wire. Section 5 deals with magnetic field dependence of the anomalies and shows how they are related the spin-split steps in a perfect quantum wires. Section 6 examines the dependence of the anomalies on asymmetry introduced by finite source-drain voltage an finally in Section 7 we summarise and speculate on the possibility of a new kind of spin-polarised ground state.

2. Basic model

We consider a straight quantum wire with a small fluctuation in thickness giving rise to a weak 'bulge' as shown in Figure 1. The precise details of the bulge are largely unimportant for what follows, the main requirement being that the change in the width of the wire is sufficiently gradual that inter-channel mixing of the transverse modes is negligible and that only one electron may be bound in the bulge region. The latter is always the case for a weak symmetric bulge, which has at least one bound state which can only sustain one electron due to Coulomb repulsion. The problem reduces to electrons moving in an effective weak potential well if we confine ourselves (by choice of gate voltage) to the Fermi energies for which no more than one transverse mode is occupied, i.e. the conductance threshold and the first conductance step. This effective potential well is shown in Figure 1. In fact the potential well may be due to an actual potential fluctuation due to a remote gate or charged impurities, or even some self-consistent effect due to the electrons themselves. We shall not consider the possible cause of this weak potential further but emphasise that because it may arise in many ways, the weak potential well model is very general with widespread applicability

Consider now the motion of electrons in the wire near the conductance threshold [14]. A single electron will be bound in the potential well region and the remaining electrons will undergo scattering from the localised electron via the Coulomb interaction as they propagate from source to drain. At sufficiently low Fermi energy, the electrons in the source contact will be totally reflected by the bound electron due to Coulomb repulsion and there will be no current from source to drain at $T = 0$. As the Fermi energy is raised, the energy of the electrons in the source contact will be sufficiently high for them to overcome the Coulomb repulsion of the bound electron and a current will flow. In calculating this current we will make the approximation that the electrons flowing from source to drain only interact with the bound electron via a screened Coulomb interaction. This is a reasonable approximation provided that the electron density is not too low in the region of interest, i.e. the rising edge to the first

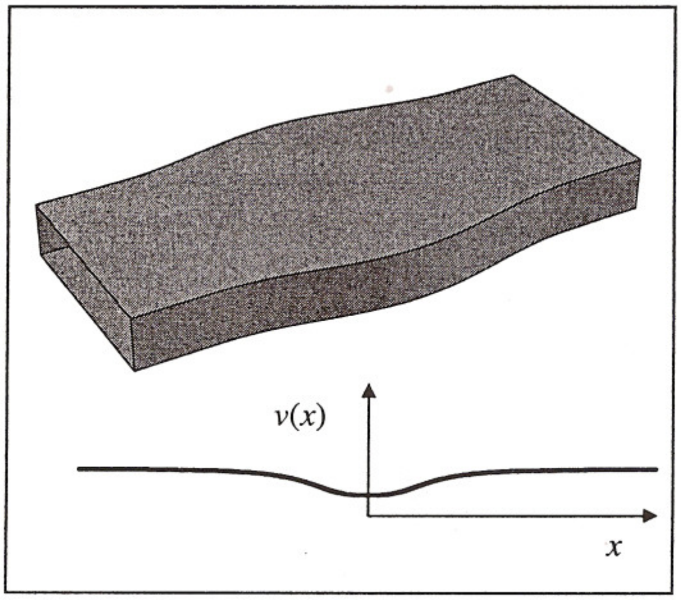

Figure 1. Schematic view of rectangular cross-section wire with bulge and corresponding effective one-dimensional potential.

conductance plateau. More precisely, the mean density of electrons in the wire (number per unit length) should be at least of order the inverse effective Bohr radius of the material. We return to this point again in the final Section. Within this approximation, the many-electron problem is reduced to an effective two-electron problem in which one electron is bound and the other is a representative electron at the Fermi energy in the leads. We show below that by solving this two-electron problem exactly and summing over all electrons near the Fermi energy we may compute the conductance.

3. Scattering problem in zero magnetic field

The two-electron problem in zero magnetic field is described by the effective Schrödinger equation,

$$\left[-\frac{\hbar^2}{2m^*}\left(\frac{d^2}{dx_1^2}+\frac{d^2}{dx_2^2}\right)+v\left(x_1\right)+v\left(x_2\right)+U(x_1,x_2)\right]\psi(x_1,x_2)=E\psi(x_1,x_2), \tag{1}$$

where $U(x_1,x_2)=e^2/(4\pi\varepsilon\varepsilon_0 d(x_1,x_2)\exp(-|x_1-x_2|/\rho)$, with $d(x_1,x_2)$ given by integrating over the Coulomb interaction over the lowest (nodeless) transverse wavefunction in the yz plane. $d(x_1,x_2)\to|x_1-x_2|$ when d is much larger than the width w of the wire, and $d\sim w$ for $x_1=x_2$. The dielectric constant is taken as $\varepsilon = 12.5$, appropriate for GaAs. Equation (1) is solved as a scattering problem with boundary conditions that one electron is bound in the wire and the other electron is in an unbound plane wave state asymptotically, i.e. there is an incident and reflected plane wave near the source contact and a transmitted plane wave moving towards the drain contact. In solving this scattering problem it is important to take into account the indistinguishability of the two electrons by antisymmetrising the states. Since in the case of two electrons, the spin-orbitals factor into spin part and orbital part, it is most efficient to deal with the orbital parts only since the effective Hamiltonian does not depend explicitly on spin. Thus the scattering problem was solved for both symmetric orbital wavefunctions corresponding to singlet states, and antisymmetric orbital functions corresponding to triplet states. The finite difference method was used in solving the equations, starting with the asymptotic outgoing plane wave for the unbound electron, with the bound electron in the lowest allowed bound state consistent with total spin. In these calculations the step-length was made progressively smaller to ensure convergence. We note that the choice of boundary condition is only consistent if the total energy of the two-electron system is sufficiently small, i.e. $\varepsilon_F < |\varepsilon_0|$, since this ensures that one electron remains bound asymptotically. This condition is fulfilled for most cases of interest near the conductance threshold. At higher energies, we would have to take into account quasi-inelastic processes in which either one or both electrons are unbound asymptotically. This case is considered elsewhere [15].

The solutions of the scattering problem gives the singlet and triplet transmission coefficients for an 'incident' electron at the Fermi energy. Typical transmission probabilities are shown in Figure 2 near the conductance threshold in which the Fermi energy is gradually increased by application of a gate voltage. We see that in all cases, the transmission undergoes a resonance (becomes unity) at a certain Fermi energy/gate voltage and that this resonance peak is sharper for the triplet than for the singlet. Furthermore, the resonances shift to lower energy with increased screening. This behaviour may be understood as follows. As the unbound electron moves towards the bound electron it sees a progressively increasing screened Coulomb potential until it enters the weak potential well region when there is a small decrease in total energy, since the decrease in the one-electron potential energies of the two electrons outweighs their increase in energy due to Coulomb repulsion. For a symmetric potential well this clearly leads to a double-barrier structure, as shown in Figure 3. For an asymmetric potential, we also get a double barrier provided that the asymmetry is not too large (see also Section 6). Associated with this double barrier structure will be at least one resonance, as indicated in Figure 3. Furthermore, for a weak

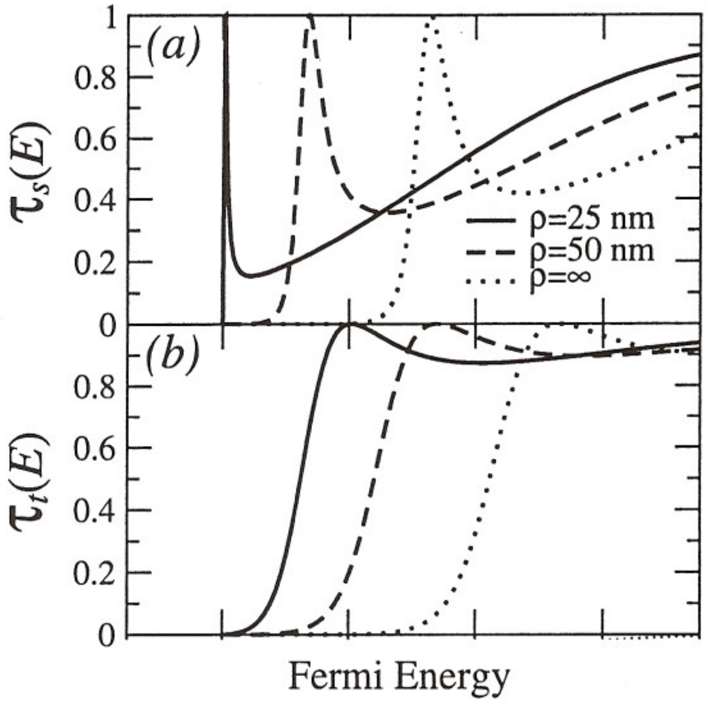

Figure 2. Typical singlet (a) and triplet (b) transmission probabilities.

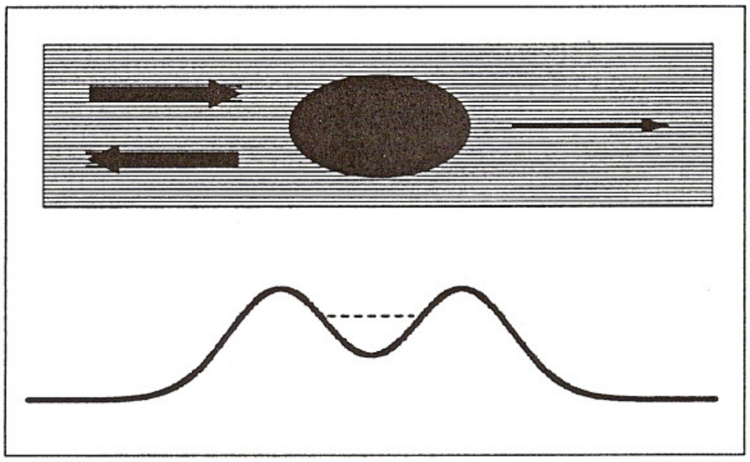

Figure 3. Scattering of conduction electron from localised electron showing effective Coulomb double barrier and resonant bound state.

potential well, which is the case of interest with our 'clean' quantum wires, there will be only one such resonance. This partly explains the transmission results since with such a double barrier structure there will always be an energy for which there is perfect transmission. However, we emphasise that this is not a double barrier in the usual sense of a variation in one-electron potential, but depends intrinsically on the repulsion between the two electrons—there is no such double barrier and resonance for non-interacting electrons. The spin-dependence of these resonances is a little more subtle but may be understood within the framework of an effective Anderson impurity model. Consider the effect of gradually switching on the Coulomb interaction. For a range of parameter sets which correspond to a very weak bulge/confining potential, there are two bound states for one electron. With no interaction both electrons may thus occupy one of 4 states (3 singlets and a triplet). If we now switch on a small Coulomb interaction then the lowest two-electron state will be a singlet, derived from both electrons in the lowest one-electron state. We may regard one electron as occupying the lowest bound-state level and the other electron of opposite spin also in this same orbital state but energy U higher, where U is the intra-'atomic' Coulomb matrix element, as in the Anderson impurity model [16]. As the Coulomb interaction is increased, U eventually exceeds the binding energy and this higher level becomes a virtual bound state giving rise to a resonance in transmission. An estimate of the energy of the virtual bound state is given by the Anderson-Coulomb matrix element with both electrons in the one-electron orbital ψ_0, i.e., $U = \int dx dx' |\psi_0(x)|^2 |\psi_0(x')|^2 U(x, x')$. We have computed this and obtained reasonable agreement with the exact result. We can in addition approximate the full scattering problem by solving the Hartree-Fock equations without iteration in which one of the electrons again occupies ψ_0. The agreement is also very good and reproduces all the resonance features. When both electrons have the same spin, they must occupy different orbitals in the dot region when the Coulomb interaction is switched off. With small Coulomb interaction the resulting triplet is the lowest two-electron excited state and this develops into a resonant bound state with the full Coulomb interaction, with energy at approximately $E_1 + U_1 - J_1$, where E_1 is the energy of the second one-electron state with U_1 and J_1 the respective Coulomb and exchange integrals. We can now see why the singlet resonance is somewhat sharper than the triplet in Figure 2; this is simply because it is lower in energy and closer to a 'real' bound state. This is also the reason why screening pushes the resonances to lower energy, since the Coulomb integrals become smaller.

4. Conductance in zero magnetic field

Conductance arising from ballistic transport of non-interacting electrons in a one-dimensional quantum wire may be calculated using the Landauer-Büttiker formula [17],

$$G = \frac{2e^2}{h} \int T(\varepsilon) \frac{\partial f(\varepsilon_F - \varepsilon, T)}{\partial \varepsilon} d\varepsilon \tag{2}$$

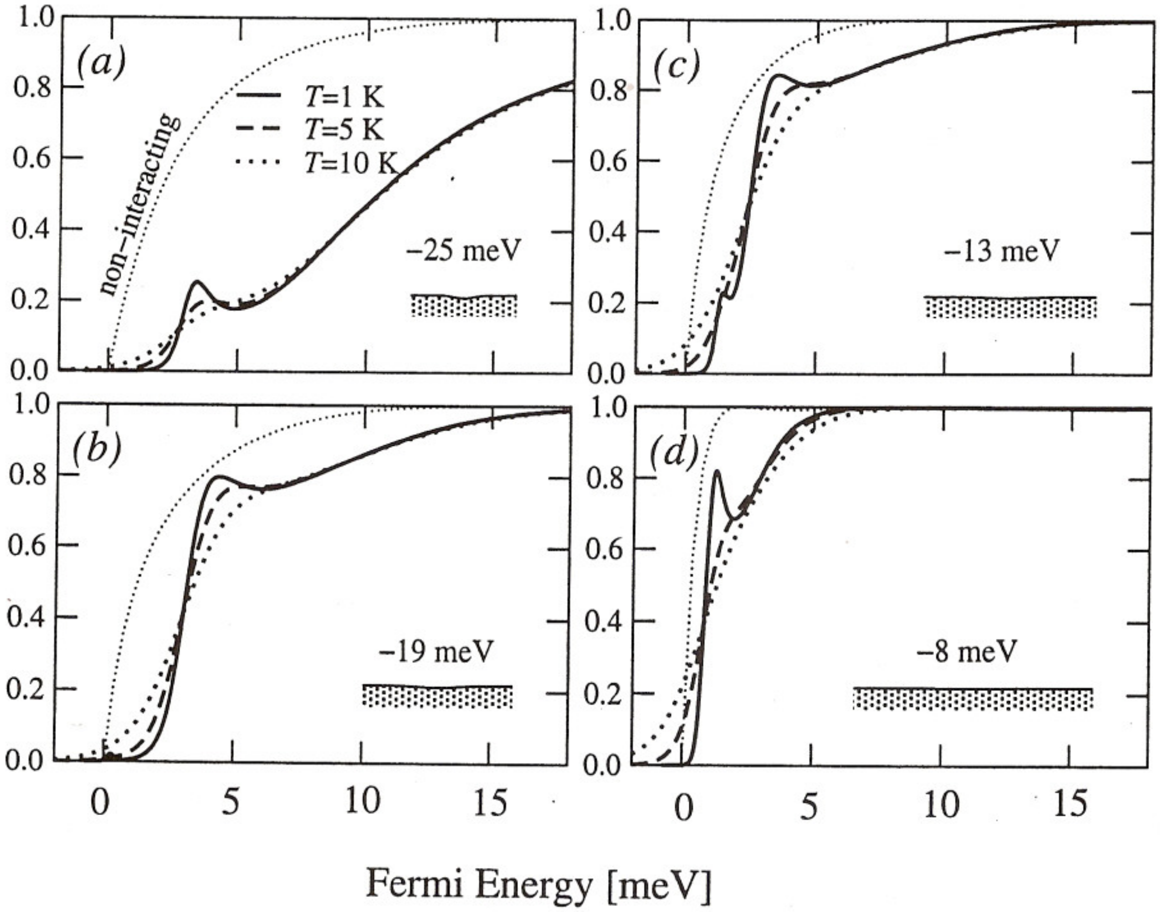

Figure 4. Conductance in units of $2e^2/h$ vs Fermi energy in meV for various single-electron confinement potentials. Insets schematically represent confinement potentials together with depth $V_{\rm eff}$ in meV.

where $\mathcal{T}(\varepsilon)$ is the transmission probability for an unbound electron of energy ε near the Fermi energy ε_F in the source lead and $f(\varepsilon, T) = [1 + \exp(\varepsilon/k_{\rm B}T)]^{-1}$ is the usual Fermi function which describes the thermal distribution of electrons in the leads. At zero temperature the Fermi function is a step function and hence $G \to \frac{2e^2}{h}\mathcal{T}(\varepsilon_F)$, which for a perfect wire gives the well-known step function from zero to a conductance plateau at $G = \frac{2e^2}{h}$. We may extend this formula to the case described in the previous Section in which one electron is bound in the wire and the remaining electrons are transmitted with energy-dependent probability, as shown for example in Figure 2. Let P_σ be the probability that the bound electron has spin σ. It follows directly that the conductance due to all spin-up electrons in the leads is given by the extended Landauer Büttiker formula:

$$G_\uparrow = \frac{e^2}{h}\int \left[P_\uparrow \mathcal{T}_{\uparrow\uparrow}(\varepsilon) + P_\downarrow \mathcal{T}_{\uparrow\downarrow}(\varepsilon)\right] \frac{\partial f(\varepsilon_F - \varepsilon, T)}{\partial \varepsilon} d\varepsilon,$$

where $\mathcal{T}_{\uparrow\uparrow}$ is the transmission probability when the bound electron is spin up and $\mathcal{T}_{\uparrow\downarrow}$ is the transmission probability when the bound electron is spin down. We have a similar expression for spin-down electrons in the leads and hence the total conductance is

$$G = \frac{e^2}{h}\int \left[P_\uparrow \mathcal{T}_{\uparrow\uparrow}(\varepsilon) + P_\downarrow \mathcal{T}_{\uparrow\downarrow}(\varepsilon) + P_\uparrow \mathcal{T}_{\downarrow\uparrow}(\varepsilon) + P_\downarrow \mathcal{T}_{\downarrow\downarrow}(\varepsilon)\right] \frac{\partial f(\varepsilon_F - \varepsilon, T)}{\partial \varepsilon} d\varepsilon. \quad (3)$$

The transition probabilities $\mathcal{T}_{\uparrow\uparrow}$ and $\mathcal{T}_{\uparrow\downarrow}$ are different since in the former case the conduction and bound electrons both have the same spin (up) before and after scattering whereas in the latter case there are two possible final states, with or without spin flip, i.e.

$$\mathcal{T}_{\uparrow\downarrow} = |\langle\uparrow,\downarrow|\mathbf{t}|\uparrow,\downarrow\rangle|^2 + |\langle\downarrow,\uparrow|\mathbf{t}|\uparrow,\downarrow\rangle|^2,$$

where $\mathbf{t}$ is the transition operator connecting ingoing and outgoing states. For the zero magnetic field case, for which we clearly have $P_\uparrow = P_\downarrow = 1/2$, G may be written in a simpler form by transforming to singlet and triplet states with $S_z = 0$, i.e.

$$|s\rangle = \frac{|\uparrow,\downarrow\rangle - |\downarrow,\uparrow\rangle}{\sqrt{2}}, \quad |t\rangle = \frac{|\uparrow,\downarrow\rangle + |\downarrow,\uparrow\rangle}{\sqrt{2}}$$

and hence

$$|\uparrow,\downarrow\rangle = \frac{|t\rangle + |s\rangle}{\sqrt{2}}, \quad |\downarrow,\uparrow\rangle = \frac{|t\rangle - |s\rangle}{\sqrt{2}}$$

giving

$$\mathcal{T}_{\uparrow\downarrow} = \frac{\mathcal{T}_s + \mathcal{T}_t}{2},$$

where $\mathcal{T}_s = |\langle s|\mathbf{t}|s\rangle|^2$, $\mathcal{T}_t = |\langle t|\mathbf{t}|t\rangle|^2 = \mathcal{T}_{\uparrow\uparrow} = \mathcal{T}_{\downarrow\downarrow}$. Hence, finally,

$$G = \frac{2e^2}{h}\int \left[\frac{3}{4}\mathcal{T}_t(\varepsilon) + \frac{1}{4}\mathcal{T}_s(\varepsilon)\right] \frac{\partial f(\varepsilon_F - \varepsilon, T)}{\partial \varepsilon} d\varepsilon \quad (4)$$

which reduces to the non-interacting case (2) in the appropriate limit for which $\mathcal{T}_t = \mathcal{T}_s = \mathcal{T}$. Equation (4) clearly shows the origin of the anomalies in conduction which occur near $\frac{1}{4}\frac{2e^2}{h}$ and $\frac{3}{4}\frac{2e^2}{h}$, when either $\mathcal{T}_s$ or $\mathcal{T}_t$ are close to unity. These anomalies are shown in Figure 4 for typical effective confining potentials and temperatures. For ease of comparison we have chosen wires with small expansions in thickness as shown in the insets together with confining potential

V_{eff} in meV. We see that the effect of temperature is to smooth out the resonance features into small shoulders, as observed in experiment. Furthermore, the singlet resonance is less pronounced or even unresolvable for the 'more perfect' wires [Figures 4(b)–(d)], whereas the case with the largest deviation in width (a) clearly shows the singlet resonance. This behaviour is simply due to the narrowness of the lower-energy singlet resonance which inhibits its resolution, except for case (a) in which the increase in energy of the triplet resonance has already pushed it into the plateau region where $G \sim \frac{2e^2}{h}$.

5. Finite magnetic field

The effect of a finite magnetic field along the length of the wire is to Zeeman-split the energy of spin-up and spin-down electrons, i.e. we add to the Hamiltonian of (1) the Zeeman term $\frac{1}{2}g^*\mu_B B(\sigma_{1z}+\sigma_{2z})$. Near the conductance threshold, the localised electron will be in its lowest spin-down Zeeman state before each scattering event and hence, from equation (3), the conductance becomes

$$G(\varepsilon_F, B, T) = \tag{5}$$

$$= \frac{e^2}{h}\int \left[\mathcal{T}_{\downarrow\downarrow}(\varepsilon+\varepsilon_B) + \mathcal{T}_{\uparrow\downarrow}(\varepsilon-\varepsilon_B)\right] \frac{\partial f(\varepsilon_F-\varepsilon, T)}{\partial \varepsilon} d\varepsilon =$$

$$= \frac{e^2}{h}\int \left[\mathcal{T}_t(\varepsilon+\varepsilon_B) + \frac{1}{2}\mathcal{T}_t(\varepsilon-\varepsilon_B) + \frac{1}{2}\mathcal{T}_s(\varepsilon-\varepsilon_B)\right] \frac{\partial f(\varepsilon_F-\varepsilon, T)}{\partial \varepsilon} d\varepsilon$$

where we have used $\mathcal{T}_{\downarrow\downarrow}(\varepsilon, B) = \mathcal{T}_{\downarrow\downarrow}(\varepsilon+\varepsilon_B, 0) \equiv \mathcal{T}_{\downarrow\downarrow}(\varepsilon+\varepsilon_B)$ and $\mathcal{T}_{\uparrow\downarrow}(\varepsilon, B) = \mathcal{T}_{\uparrow\downarrow}(\varepsilon-\varepsilon_B, 0) \equiv \mathcal{T}_{\uparrow\downarrow}(\varepsilon-\varepsilon_B)$, where $\varepsilon_B = \frac{1}{2}g^*\mu_B B$. This follows from the fact that a spin up (down) electron in the lead will have its energy raised (lowered) by ε_B. We have included the spin-flip term in this equation, which assumes that the scattered electron, which lies $2\varepsilon_B$ below the Fermi energy, is not reflected by the collector. This necessitates inelastic processes in the collector and the approximation may break down in some circumstances which we shall not consider further here. However, at low temperatures and in a high magnetic field, the number of higher-energy spin-up electrons in the lead becomes negligible near the threshold and we get

$$G(\varepsilon_F, B, T) = \frac{e^2}{h}\int \mathcal{T}_t(\varepsilon+\varepsilon_B)\frac{\partial f(\varepsilon_F-\varepsilon, T)}{\partial \varepsilon} d\varepsilon. \tag{6}$$

This is plotted in Figure 5 for $T = 3$ K together with the corresponding results for non-interacting electrons and a straight wire. We see that these curves are very similar with a plateau at e^2/h but with the interacting case displaced to the right (due to the Coulomb repulsion) and showing a slight but distinctive dip, due to the broad triplet resonance. This curve is very similar to high-field experimental curves on gated 2DEG wires which show the "0.7" anomaly [8], further supporting the view that an electron is weakly bound in the wire.

6. Dependence on source-drain bias

The existence of the conductance anomalies described above is a direct consequence of an effective double-barrier potential seen by the conduction electrons propagating from source to drain contacts under the influence of a bound electron. For a symmetric one-electron confining potential, the existence of a bound state is guaranteed but this is not necessarily the case when the confinement is asymmetric. Such asymmetry in the confining potential may be easily achieved under a finite source drain bias and indeed, this was reported in some of the experiments on gated quantum wires [8, 18]. These experiments show that as the source-drain bias is increased from zero, an anomaly appears at $G \sim 0.25(2e^2/h)$, coexisting with the $0.7(2e^2/h)$ anomaly. This sharpens as the bias is increased and, in the example of [18], for $V_{sd} \sim 6$ mV the 0.25 anomaly is very pronounced whilst the 0.7 anomaly has disappeared. Eventually, at much larger bias, the remaining anomaly also disappears but only when the conductance steps themselves are on the point of disappearing, showing that the singlet anomaly is extremely robust. This behaviour is consistent with our model since under bias the triplet resonant bound-state will eventually disappear because the confining potential in the x-direction will only accommodate a single one-electron bound state, giving rise to a singlet resonance only. This is shown schematically in Figure 6 where we also indicate the surviving becoming broader with increasing bias resulting in a more pronounced step, as observed.

7. Summary and discussion

In summary, we have shown that quantum wires with weak longitudinal confinement, or open quantum dots, can give rise to spin-dependent, Coulomb blockade resonances when a single electron is bound in the confined region. This is a universal effect in one-dimensional systems with very weak longitudinal confinement. The emergence of a specific structure at $G(E) \sim \frac{1}{4}\frac{2e^2}{h}$ and $G \sim \frac{3}{4}\frac{2e^2}{h}$ is a consequence of the singlet and triplet nature of the resonances and the probability ratio 1:3 for singlet and triplet scattering and as such is a universal effect. A comprehensive numerical investigation of open quantum dots using a wide range of parameters shows that singlet resonances are always at lower energies

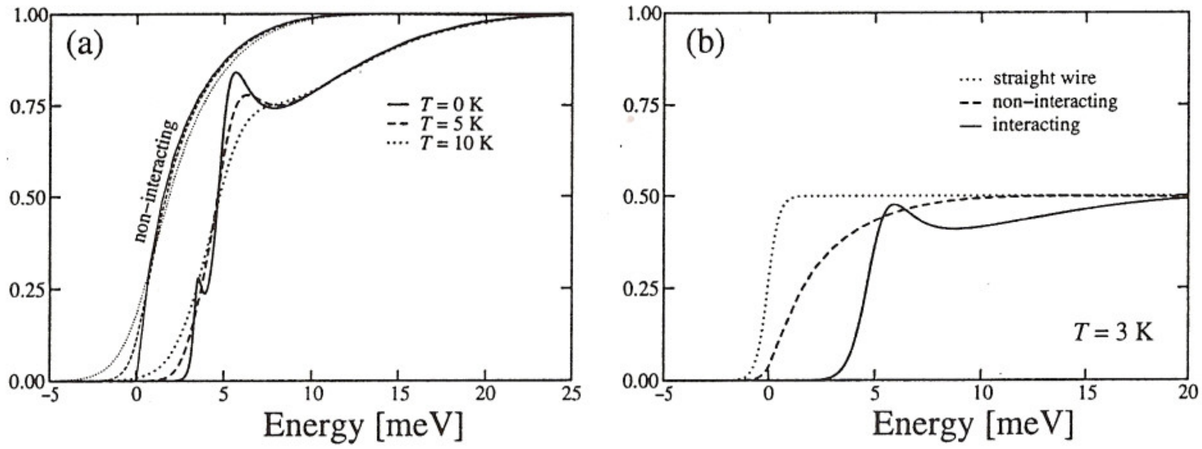

Figure 5. Change in conductance (in units of $2e^2/h$) when a high magnetic field along the wire is switched on. (a) $B = 0$ showing singlet and triplet anomalies as the Fermi energy is changed. (b) Large B showing Zeeman splitting and the experimentally observed characteristic resonance of spin-polarised electrons.

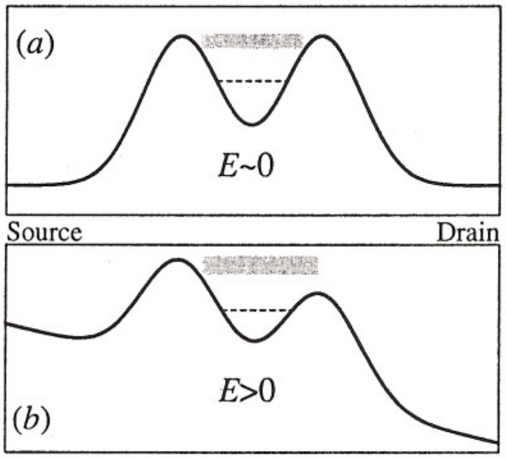

Figure 6. Effective double barrier showing the singlet and triplet resonances with very small source-drain bias (a) and large source-drain bias (b).

than the triplets, in accordance with the corresponding theorem for bound states [19]. With increasing in-plane magnetic field, the resonances shift their position and eventually merge in the conductance plateau at $G \sim e^2/h$. With increasing source-drain bias we have shown why the higher triplet resonance weakens at the expense of the singlet, with the latter surviving to the point where the conductance steps themselves disappear.

Finally, we speculate on the exciting possibility that these anomalies in conduction are themselves a signature for a new kind of conducting state in ultra-clean wires close to the conduction threshold. Indeed, there is some experimental evidence for this in that the anomalies discussed above merge into a conductance step at e^2/h under quite moderate magnetic fields and in the cleanest samples this behaviour is sometimes even seen in zero magnetic field. This suggests that there may be an underlying spin-polarised state associated with the propagating electrons in the quasi-1D region. Such a spin-polarised state would appear to violate the Lieb-Mattis theorem and would also need to be made consistent with our above explanation in terms of singlet and triplet resonances. In this respect we emphasise that the above theory must break down at very low electron density in the wire such that the mean separation between electrons in the wire is somewhat greater than the effective Bohr radius, the so-called strong correlation regime. In practical situations it is very difficult to avoid some kind of weak potential fluctuation which traps one electron. Indeed this may ultimately be impossible since even in a nominally perfect wire, the presence of a single electron will polarise its environment leading to a potential well which will bind the electron giving rise to a Coulomb blockade for the remaining electrons. The main question is whether or not this confinement is sufficiently large for the electron density to exceed the inverse Bohr radius when the wire begins to conduct. If the density remains low at this conductance threshold then we cannot ignore the mutual interaction between all electrons in the wire region, or even treat them self-consistently. In this situation, a more appropriate picture would be one in which the Coulomb repulsion dominates and maintains roughly equal separation between the electrons as in a Wigner chain. However, this would be a 'sliding' Wigner chain, carrying the current from source to drain contacts. How could this lead to a spin-polarised state? The precise details of this difficult problem are unknown at present but we do know of closed systems where spin-polarisation does indeed give rise to quasi-1D spin-polarised states. An example is the case of three electrons in a ring or a disc in the strong correlation regime for which the ground state is $S = 3/2$ and corresponds to all three electrons rotating together, as in a rigid rotor. A similar behaviour occurs for three electrons in a square confining potential which again has a spin-polarised ground state [20]. The important feature of these results is that the spin alignment is a kinetic-energy effect associated with the simultaneous propagation of the electrons and not a ferromagnetic exchange interaction between them. This can be seen directly by introducing a barrier which inhibits the movement of the electrons, resulting in a low-spin ground state. Thus there is an intrinsic relationship between motion of strongly correlated electrons and spin polarisation which may be the underlying cause of a new conducting state in quantum wires, but more research is needed to further develop and quantify these ideas. Further experiments in which the widths of quantum wires and/or the confinement potentials are engineered to control longitudinal confinement should throw also further light on the problem of spin-dependent ballistic transport.

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