Resonant transmission through an open quantum dot

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We have measured the low-temperature transport properties of a quantum dot formed in a one-dimensional channel. In zero magnetic field this device shows quantized ballistic conductance plateaus with resonant tunneling peaks in each transition region between plateaus. Studies of this structure as a function of applied perpendicular magnetic field and source-drain bias indicate that resonant structure deriving from tightly bound states is split by Coulomb charging at zero magnetic field. [S0163-1829(97)06212-7]

Advancing technology has made it possible to define artificial semiconductor microstructures that confine electrons in all three spatial dimensions[1] with discrete zero-dimensional states. Such structures, often called quantum dots, provide uniquely simple systems for the study of few-electron physics. In particular, the Coulomb blockade (CB) of single electron tunneling through quantum dots[2] has been extensively investigated.[3] It has been demonstrated[4] that transport through small quantum dots is determined by charging effects[5,6] as well as quantum confinement effects.[7–9] Quantum dots can also be formed by impurities that are either directly in the electron gas, as for Si devices,[10] or are remote ionized donors in a spacer layer[11] as for the GaAs/Al$_x$Ga$_{1-x}$As heterojunction.[12] The CB effects in such unintentionally defined quantum dots have been studied extensively.[10,12,13]

Within a noninteracting picture Tekman and Ciraci[14] have predicted that resonant tunneling (RT) may occur through energy states bound to an attractive impurity potential in a split-gate device even when some one-dimensional (1D) channels are perfectly transmitted. Therefore in addition to 1D quantized conductance steps,[15,16] replicated resonant features between plateaus should be observed when a quantum dot formed by an impurity potential is present in a split-gate device. In this paper, we report the observation of such resonant structure from a quantum dot formed by an impurity potential in a split-gate device. We show how these RT features develop in a perpendicular magnetic field B and we investigate the energy spacings between different resonant states using source-drain bias measurements.

The Schottky gate pattern shown in the inset to Fig. 1 was defined by electron beam lithography on the surface of a GaAs/Al$_{0.3}$Ga$_{0.7}$As heterostructure, 90 nm above a two-dimensional electron gas (2DEG). The carrier concentration of the 2DEG was 3.3×10^{15} m^{-2} with a mobility of 90 m^2/V s. Experiments were performed in a dilution refrigerator at 100 mK and the two-terminal differential conductance $G=dI/dV$ was measured using an ac excitation voltage of 10 μV.

Figure 1 shows the differential conductance as a function of the voltage V_{g1} on gate 1, for various voltages V_{g2} on gate 2. For $V_{g2}=-1.7$ V (trace 3) we observe replicated resonant peaks in $G(V_{g1})$, reminiscent of those predicted.[14] As the temperature was increased, these structures became broader but were still discernible up to 650 mK. When the conduction channel through the split-gate structure was moved sideways by varying the voltage on gate 2 (Ref. 17) away from $V_{g2}=-1.7$ V, the sharp RT features gradually diminished until at $V_{g2}=-2.6$ V (trace 7) only quantized 1D ballistic conductance steps were seen. In a subsequent cooldown in a ^{3}He cryostat, we did not observe identical RT structure. Although the surface Schottky gate pattern was intended to define a quantum dot in the 2DEG electrostatically, both observations suggest that ionized impurities in the spacer layer[11] played an important role in determining the transport properties through the channel defined by the surface gates. Since we observe conductance peaks (resonant tunneling) rather than resistance peaks (resonant reflection), we believe that in our system an attractive impurity potential helped create a quantum dot. Previously McEuen *et al.*[18] claimed that two resonant transmission peaks they observed for $G<2e^2/h$ in a disordered split-gate device derived from the formation of a quantum dot by a *single* hydrogenic impurity. In this experiment only two peaks were observed because the electrons that filled the impurity bound states acted to screen

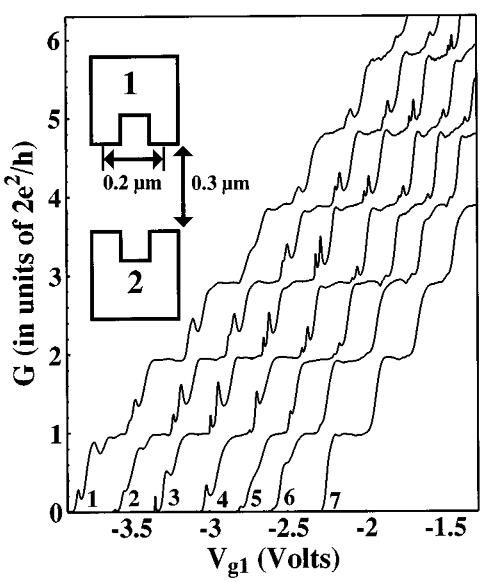

FIG. 1. $G(V_{g1})$ when the conduction path is electrostatically shifted by applying various gate voltages to gate 2. Traces 1 to 7: $V_{g2}=-1.3$, -1.5, -1.7, -1.9, -2.1, -2.3, and -2.6 V, respectively. The inset shows the Schottky gate geometry.

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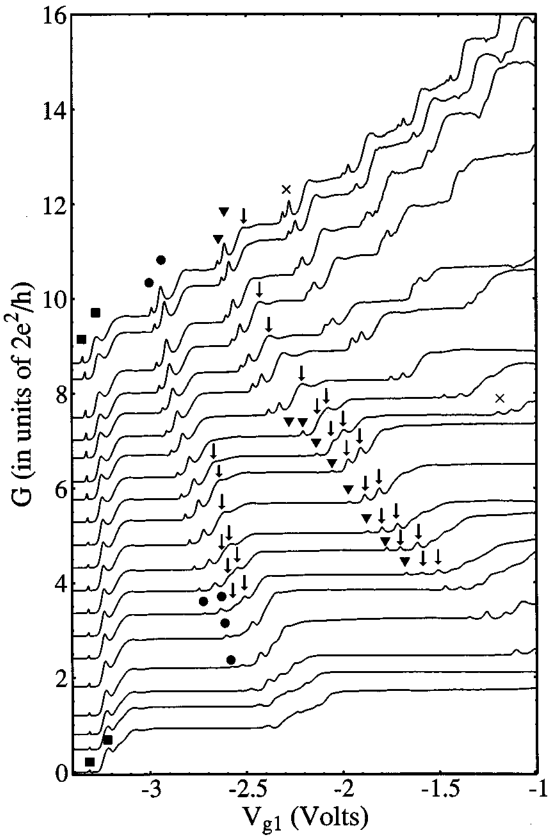

FIG. 2. $G(V_{g1})$ for $V_{g2}=-1.7$ V at various magnetic fields. The corresponding magnetic fields are, from bottom to top: 2.5, 2.4, 2.3, 2.2, 2.1, 2, 1.9, 1.8, 1.7, 1.6, 1.5, 1.4, 1.2, 1, 0.8, 0.6, 0.4, 0.2, and 0 T. Traces are vertically offset for clarity. The symbols indicate the evolution of the RT features for $G<2e^2/h$ (square), $2e^2/h<G<4e^2/h$ (circle), $4e^2/h<G<6e^2/h$ (triangle), and $6e^2/h<G<8e^2/h$ (cross) as the applied magnetic field is increased from 0 T. Arrows serve as a guide to the eye indicating a single resonant peak splits into two for $2e^2/h<G<4e^2/h$ and $4e^2/h<G<6e^2/h$, respectively.

it so that at higher energies only quantized conductance with no resonant structure was seen. In our experiment we observe at least fourteen resonant peaks (see trace 3 in Fig. 1), implying that the impurity potential does not become screened even after accommodating 14 electrons. We do not believe that such a potential could be generated by a single ionized impurity, only a cluster would be capable of this.

Figure 2 shows $G(V_{g1})$ for $V_{g2}=-1.7$ V at different B. For $G<2e^2/h$, the two RT peaks have a weak B dependence and persist to $B=4$ T. As the magnetic field is increased, the conductance plateaus and the RT peak positions for $2e^2/h<G<8e^2/h$ move to more positive V_{g1} as a result of the formation of hybrid magnetoelectric subbands.[8] At $B\approx 2$ T these resonant features are no longer seen. A broad RT peak adjacent to the sharp RT peaks for $4e^2/h<G<6e^2/h$ develops at $B=0.6$ T and splits into two at higher magnetic fields, as indicated by arrows. Similar but less pronounced results can be also seen for $2e^2/h<G<4e^2/h$.

The sharp resonances correspond to tightly bound states and the broad resonances to weakly bound states within the picture of Tekman and Ciraci.[14] The application of a perpendicular magnetic field strengthens the confinement of states in a quantum dot by localizing electron wave functions to the sample boundaries.[19] This is consistent with the disappearance of the tightly bound states (they become immeasurably small) and the strengthening of the resonant structure from the weakly bound states at high field seen experimentally — Fig. 2. When the sharp RT structures for $2e^2/h<G<6e^2/h$ have disappeared, oscillations in $G(V_{g1})$ are still observed. Their structure is more complicated and possibly derives from a combination of resonant transmission and resonant reflection[20] from bound states.

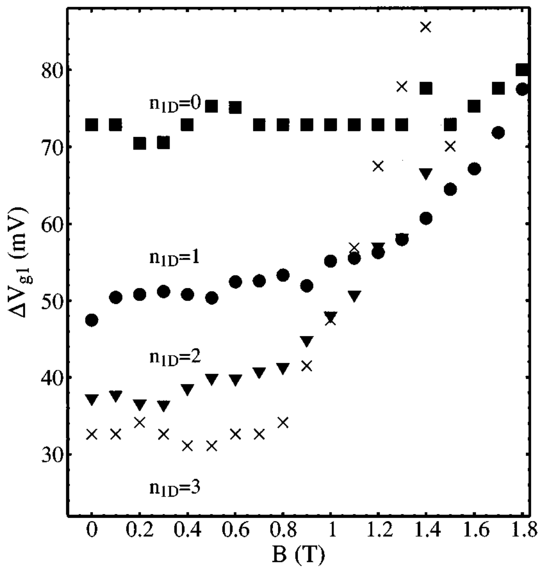

FIG. 3. $\Delta V_{g1}(B)$ for various n_{1D}.

We now discuss the separation in gate voltage ΔV_{g1} between each pair of tightly bound RT peaks at various magnetic fields (marked as square, circle, triangle, and cross in Fig. 2). Figure 3 shows $\Delta V_{g1}(B)$ for RT features that occur with different numbers of transmitted 1D channels n_{1D}. For $n_{1D}=0$, ΔV_{g1} shows only a weak magnetic-field dependence. For $n_{1D}=1$, 2, and 3, ΔV_{g1} shows saturation at low B and a linear B dependence at high B.

At $B=0$, ΔV_{g1} decreases as n_{1D} increases, as shown in the inset to Fig. 4. To obtain the energy spacing $\Delta E(n_{1D})$ between pairs of tightly bound RT peaks, we have used a standard source-drain bias technique.[6,9,12,21] ΔE decreases dramatically from $n_{1D}=0$ to $n_{1D}=1$ (see Fig. 4). Note that we were not able to measure ΔE between pairs of RT peaks for $G>8e^2/h$, perhaps because the application of a dc bias caused the quantum dot to break down.

Within the noninteracting picture[14] at $B=0$ the energy states through which RT occurs are spin degenerate. As B is increased, if there is no spin splitting, states with different angular momentum in the same Landau level become closer in energy.[22] If the Zeeman energy is included, electrons in the same Landau level with the same angular momentum but different spin move apart in energy causing individual resonant transmission peaks to split into two peaks. Using the minimum possible value 0.44 for the Landé g factor in our system, we estimate the Zeeman energy to be ≈ 0.1 meV at $B=4$ T, a factor of twelve larger than thermal smearing at

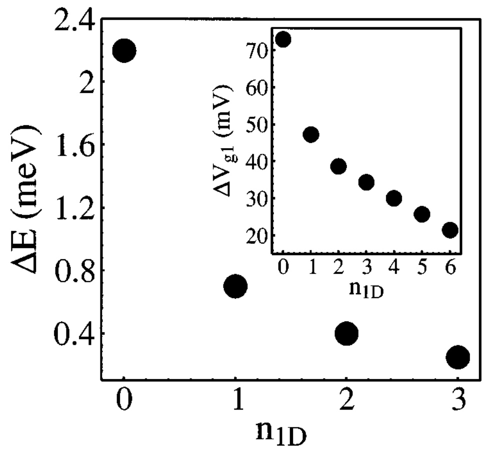

FIG. 4. The energy spacing ΔE between pairs of RT peaks as a function of n_{1D} deduced from the dc bias measurements. The inset shows $\Delta V_{g1}(n_{1D})$.

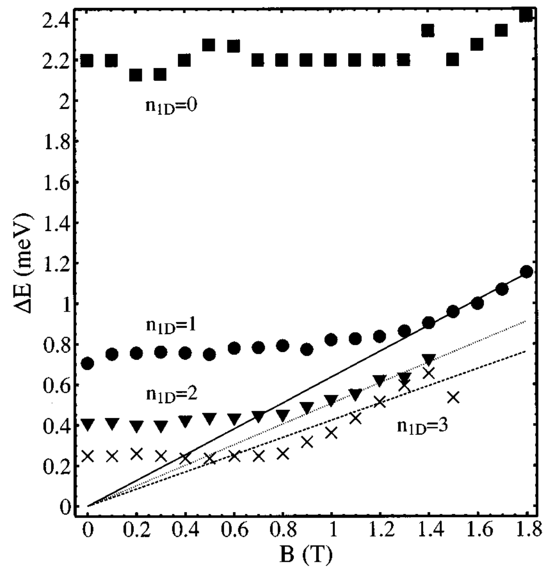

FIG. 5. The energy spacing ΔE between pairs of RT peaks as a function of B determined from data shown in Figs. 3 and 4. The straight line fits are discussed in the text.

100 mK, and equal the full width at half maximum of the tightly bound peak closest to pinch-off, suggesting that such splitting would be observable in our system. However, as shown in Fig. 2, the individual peaks in each pair of tightly bound RT peaks *do not* split at any magnetic field. In addition pairs of peaks *do not* come closer together and each pair of peaks remains in the same transition region. These factors imply that each pair of peaks derives from the same single-particle state. They split at zero magnetic field due to the energy difference between single and double occupation of a single state.[23] However, the case of charging-induced splitting in mesoscopic devices, where two adjacent single-electron tunneling peaks are related to states with different spin quantum numbers,[24] is only well understood in the Coulomb blockade regime[25] for $G<2e^2/h$. Assuming that the relations $\Delta E=30.1\Delta V_{g1}$ meV/V ($n_{1D}=0$), $\Delta E=14.8\Delta V_{g1}$ meV/V ($n_{1D}=1$), $\Delta E=10.3\Delta V_{g1}$ meV/V ($n_{1D}=2$), and $\Delta E=7.27\Delta V_{g1}$ meV/V ($n_{1D}=3$) (determined from the data shown in Fig. 4 and the inset), which hold at $B=0$ are still valid at high field, $\Delta E(B)$ for the tightly bound peaks with $n_{1D}=1$, 2, and 3 shown in Fig. 5 also implies charging-induced splitting at $B=0$. If the splitting arose solely from Zeeman splitting, then one would expect $\Delta E(B)\rightarrow 0$ as $B\rightarrow 0$. Instead $\Delta E(B)$ shows saturation at low B, suggesting that the splitting at low fields is due to some effect other than Zeeman splitting. The linear fits $\Delta E=0.636B$ (solid line), $\Delta E=0.507B$ (dotted line), and $\Delta E=0.424B$ (dashed line) shown in Fig. 5 yield Landé g factors of 10.9, 8.7, and 6.9 for $n_{1D}=1$, 2, and 3, respectively. Such large g factors have been measured in the quantum Hall regime where exchange energy is important.[26] For the case $n_{1D}=0$, $\Delta E(B)$ has a weak B dependence since near pinch-off the Coulomb charging effect is much stronger than the Zeeman term.

Having established the role of Coulomb charging effects in our system, we can now explain the splitting of the broad resonant tunneling peak, indicated by arrows in Fig. 2, which occurs as B is increased. For $B=0.6$ T, the broad resonant tunneling peak is spin degenerate, as the state through which RT occurs is weakly bound and the Coulomb charging arising from confinement is not pronounced. At higher B this state becomes more tightly bound, increasing the contribution of Coulomb charging effects. Therefore when the applied magnetic field is increased from $B=0.6$ T, both the Zeeman term and the Coulomb charging lift the electron spin degeneracy, causing the broad resonant tunneling peak to split into two.

The decrease of $\Delta E(n_{1D})$ as n_{1D} is increased, at $B=0$, shown in Fig. 4, arises from two mechanisms: the Coulomb force between electrons bound in the quantum dot is increasingly screened as n_{1D} is increased; and the conduction channel defined by the surface Schottky gates becomes wider, increasing the spatial extent of the bound state wave functions, and hence reducing the Coulomb charging energy as n_{1D} is increased.

Although we can explain our results in terms of Coulomb charging effects qualitatively, it is important to note that ascribing the pairs of sharp RT features to zero-field splitting, for $G>2e^2/h$, requires an extension of the Coulomb charging picture to the metallic regime where some 1D channels are transmitted, and that the Coulomb interactions between pairs of electrons are partially screened by these 1D channels. In principle the results we present here are able to give information on the ability of 1D states to screen 0D states.

In conclusion, we have reported an observation of transmission resonances through an open quantum dot. The magnetic field dependence of pairs of tunneling peaks provides experimental evidence for Coulomb charging effects at zero-field magnetic field even when some one-dimensional channels are perfectly transmitted through the open quantum dot.

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[1] C.G. Smith *et al.*, J. Phys. C **21**, L893 (1988).
[2] U. Meirav, M.A. Kastner, and S.J. Wind, Phys. Rev. Lett. **65**, 771 (1990).
[3] For a review, see H. van Houten, C.W.J. Beenakker, and A.A.M. Staring, in *Single Charge Tunneling*, edited by H. Grabert and M.H. Devoret (Plenum, New York 1992).
[4] P.L. McEuen *et al.*, Phys. Rev. Lett. **66**, 1926 (1991).
[5] C.J.B. Ford *et al.*, Nanostruct. Mater. **3**, 283 (1993).
[6] J. Weis *et al.*, Phys. Rev. Lett. **71**, 4019 (1993).
[7] B. Su, V.J. Goldman, and J.E. Cunningham, Science **255**, 313 (1992).
[8] R.C. Ashoori *et al.*, Phys. Rev. Lett. **68**, 3088 (1992).
[9] A.T. Johnson *et al.*, Phys. Rev. Lett. **69**, 1592 (1992).
[10] J.H.F. Scott-Thomas *et al.*, Phys. Rev. Lett. **62**, 583 (1989).
[11] J.H. Davies and J.A. Nixon, Phys. Rev. B **39**, 3423 (1989).
[12] J.T. Nicholls *et al.*, Phys. Rev. B **48**, 8866 (1993).
[13] M. Pepper, M.J. Uren, and R.E. Oakley, J. Phys. C **12**, L897 (1979).
[14] E. Tekman and S. Ciraci, Phys. Rev. B **42**, 9098 (1990).
[15] B.J. van Wees *et al.*, Phys. Rev. Lett. **60**, 848 (1988).
[16] D.A. Wharam *et al.*, J. Phys. C **21**, L209 (1988).
[17] R.J. Stroh and M. Pepper, J. Phys. Condens. Matter **1**, 8481 (1989); J.G. Williamson *et al.*, Phys. Rev. B **42**, 7675 (1990); L.I. Glazman and I.A. Larkin, Semicond. Sci. Technol. **6**, 32 (1991).
[18] P.L. McEuen *et al.*, Surf. Sci. **229**, 312 (1990).
[19] D.R. Mace *et al.*, Phys. Rev. B **52**, R8672 (1995).
[20] J.K. Jain and S.A. Kivelson, Phys. Rev. Lett. **60**, 1542 (1988).
[21] N.K. Patel *et al.*, Phys. Rev. B **44**, 13 549 (1991).
[22] C.G. Darwin, Proc. Cambridge Philos. Soc. **27**, 86 (1930).
[23] N.W. Ashcroft and N.D. Mermin, *Solid State Physics*, edited by H. Ehrenreich and G. Turnbull (Academic, New York, 1976).
[24] J. Weis *et al.*, Surf. Sci. **305**, 664 (1994).
[25] Y. Meir, N.S. Wingreen, and P.A. Lee, Phys. Rev. Lett. **66**, 3048 (1991); **70**, 2601 (1993).
[26] R.J. Nicholas *et al.*, Phys. Rev. B **37**, 1294 (1988).