PHYSICAL REVIEW B VOLUME 52, NUMBER 3 15 JULY 1995-I

Effect of a parallel magnetic field on the resonant-tunneling current through a quantum wire

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(Received 13 January 1995)

The effect of a magnetic field applied parallel to the current through quantum confined GaAs/AlAs resonant-tunneling diodes with submicrometer lateral dimensions is studied theoretically. A tunneling current peak associated with an intersubband transition is predicted for intermediate magnetic field. This results from a difference in the degree of confinement in the emitter and the well. The results are compared with experiments.

Resonant tunneling through quasibound states gives rise to strong peaks in the current voltage characteristics $I(V)$ of double barrier GaAs/AlAs structures.[1] Additional peaks have been observed in the $I(V)$ of resonant-tunneling diodes (RTD's) with submicrometer lateral dimensions.[2] The additional peaks may arise from either a quantum size effect or incorporation of donors in the quantum well.[3,4]

In a previous study on a RTD in the form of a quantum wire, we have shown that the peaks in $I(V)$ due to one-dimensional (1D) lateral quantization and donors may be distinguished by applying a magnetic field B perpendicular to the current direction.[2,5] In this paper we study the effect of a magnetic field applied parallel to the current direction in the 1D RTD. Electrons are confined electrostatically for $B = 0$ in such a device, and magnetic confinement dominates for sufficiently high parallel magnetic field. We therefore expect a transition from electrostatic to magnetic confinement as B is increased. For $B = 0$, only transitions confined between states with the same symmetry occur, while in the high magnetic field limit the only transitions allowed are between Landau levels with the same index in the absence of scattering. At around the transition when the magnetic and electrostatic confinement potentials are comparable, the parity conservation is broken, but inter-Landau-level transitions are permitted. In this paper we try to clarify the B dependence of $I(V)$ especially for B around this transition.

The schematic diagram of the active region of the device is shown in Fig. 1(a) together with the coordinate axes used for the following discussion. The x axis is defined to be the growth direction, and y and z are, respectively, perpendicular and parallel to the long axis of the quantum wire cavity. The diode is fabricated from a GaAs/AlAs heterostructure grown by molecular beam epitaxy using optical lithography and undercut etching in order to achieve submicrometer dimensions.[6] The GaAs quantum well width is 90 Å and the AlAs barrier width is 47 Å. The active area of the device is the region of overlap of two GaAs bars, one (thickness ℓ_t) etched in the top contact layer and the other (thickness ℓ_b) in the bottom contact layer. The physical device dimensions $\ell_t \times \ell_b \approx 0.5 \times 1.0\,\mu\text{m}^2$ but the conduction area is much smaller because of sidewall depletion and is estimated to be $\approx 0.06 \times 0.6\,\mu\text{m}^2$ [the depletion edge is schematically shown in Fig. 1(a) as a dashed line]. We have shown previously for this device that the electron motion in the y direction is quantized and that the confining potential is parabolic up to energies of $\approx$ 100 meV.[2] Thus a set of one-dimensional subbands is formed in the emitter and the well, with free electron motion assumed for the z direction. For $B = 0$ the levels in the emitter and the well (denoted by the indices i and j, respectively) have an energy separation ($\hbar = 1$ throughout) of ω_e and ω_w, which are the confining energy of the parabolic potential in the emitter and the well, respectively [see Fig. 1(b)]. Because of the asymmetric device structure, the degree

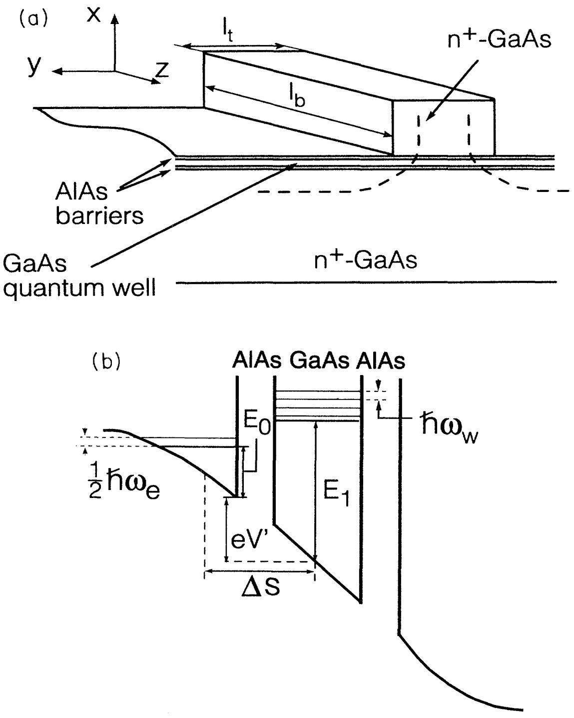

FIG. 1. (a) Schematic diagrams of active region of the diode and (b) the conduction-band profile for $B = 0$. The active area of the device is the region of overlap of two GaAs bars, one (thickness ℓ_t) etched in the top contact layer and the other (thickness ℓ_b) in the bottom contact layer.

0163-1829/95/52(3)/1504(4)/$06.00

of confinement is expected to be different in the emitter and the well. In the following we discuss only the polarity in which electrons flow from the top contact to the bottom contact. For this polarity $\omega_e > \omega_w$ and only the lowest ($i = 0$) state is occupied in the emitter for this diode.

Figure 2 shows the low temperature $I(V)$ in the presence of a magnetic field between $B = 0\,\mathrm{T}$ and $B = 6\,\mathrm{T}$. For $B = 0$, several peaks are observed in $I(V)$: (i) the main resonance ($V = 0.29\,\mathrm{V}$) due to tunneling between the $i = 0$ state in the emitter and the $j = 0$ state in the well; (ii) an additional shoulder ($V = 0.32\,\mathrm{V}$), which, as will be shown later, is due to tunneling between the $i = 0$ state and the $j = 2$ state; and (iii) peaks due to GaAs and AlAs longitudinal-optical phonon related tunneling,[7,8] which occur at $V = 0.49\,\mathrm{V}$ and $V = 0.6\,\mathrm{V}$, respectively. For a lower magnetic field the $j = 2$ shoulder[9] is only weakly affected by B. For $B > 3\,\mathrm{T}$, however, its voltage position has a strong, approximately linear, dependence on B [see Fig. 2(b)]. In this high field regime other peaks [$j = 1, 3$ in Fig. 2(b)] are observed. These peaks also have an almost linear dependence on B.

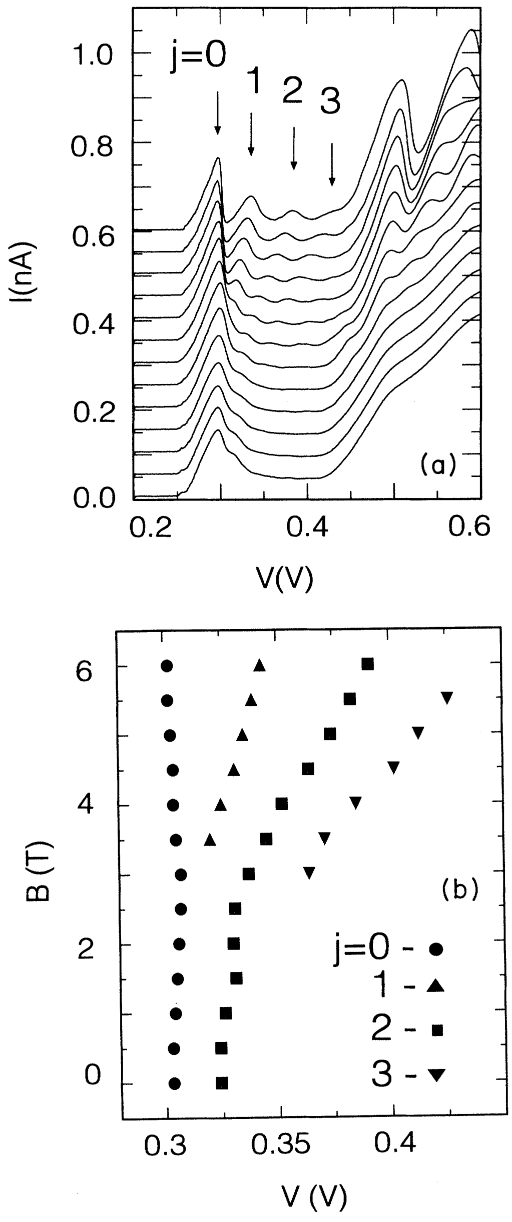

FIG. 2. (a) Experimental plots of $I(V)$ at $T = 0.3\,\mathrm{K}$ for a device with $\ell_t \times \ell_b \approx 0.5 \times 1.0\,\mu\mathrm{m}^2$ in the presence of a magnetic field between $B = 0\,\mathrm{T}$ (lowest curve) and $B = 6\,\mathrm{T}$ (top curve) in $0.5\,\mathrm{T}$ steps. The field is oriented parallel to the current direction. (b) The fan diagram shows the voltage position of peaks $j = 0, 1, 2, 3$.

We use the effective mass approximation and the transfer Hamiltonian formalism to calculate the tunneling current for the 1D RTD. Electronic states in the emitter and the well are considered separately. For a magnetic field oriented in the x direction, the single-electron Hamiltonian for the emitter states $\mathcal{H}_e$ and that for the well states $\mathcal{H}_w$ can be written, using the Landau gauge $\mathbf{A} = (0, 0, By)$, as

$$\begin{aligned}\mathcal{H}_\alpha = \frac{p_x^2}{2m} + \frac{p_y^2}{2m} + \frac{1}{2m}(p_z + eBy)^2 \\ + V_\alpha(x) + \tfrac{1}{2} m\omega_\alpha^2 y^2 \quad (\alpha = e \text{ or } w),\end{aligned} \tag{1}$$

where $V_\alpha(x)$ is the confining potential along the x direction due to the AlAs barriers and ω_α the parameter describing the approximately parabolic confinement along the y direction.[2] The eigenenergies and the corresponding eigenfunctions in the emitter and the well are given by

$$\varepsilon^e_{ik} = E^e_0 + (i + \tfrac{1}{2})\tilde{\omega}_e + \frac{k^2}{2m_e} \tag{2}$$

$$\Psi^e_{ik}(x, y, z) = \chi_e(x)\phi_i(\tilde{\ell}_e; y - Y_e)e^{ikz} \tag{3}$$

and

$$\varepsilon^w_{jk'} = E^w_0 + (j + \tfrac{1}{2})\tilde{\omega}_w + \frac{k'^2}{2m_w} \tag{4}$$

$$\Psi^w_{jk'}(x, y, z) = \chi_w(x)\phi_j(\tilde{\ell}_w; y - Y'_w)e^{ik'z} \tag{5}$$

where $\tilde{\omega}_\alpha$ [$= (\omega_\alpha^2 + \omega_c^2)^{1/2}$] is the subband spacing, ω_c ($= eB/m$) is the cyclotron energy, $m_\alpha = m(\tilde{\omega}_\alpha/\omega_\alpha)^2$, $\tilde{\ell}_\alpha = (m\tilde{\omega}_\alpha)^{-1/2}$, and $\phi_n(\ell; y)$ is the simple harmonic-oscillator-like solution ($n = 0, 1, 2, \ldots$).[10] The orbit center Y_α is related to the wave vector k: $Y_\alpha = -[\omega_c/(m\tilde{\omega}_\alpha^2)]k$. E^α_0 [$\chi_\alpha(x)$] is the solution of the one-dimensional Schrödinger equation for x motion:

$$\left\{\frac{p_x^2}{2m} + V_\alpha(x)\right\}\chi_\alpha(x) = E^\alpha_0\chi_\alpha(x). \tag{6}$$

Note that we consider only the lowest level for this equation.

In this bias direction, the strong peaks associated with optical phonons are observed (see Fig. 2(a) and Ref. 6). This may imply charge build-up in the well.[11] The tunneling current is written as

$$I = 4\pi e \sum_{ik} \sum_{jk'} \frac{\Gamma^c_{jk'}}{\Gamma^e_{jk'} + \Gamma^c_{jk'}} |T_{ik,jk'}|^2 \delta(\varepsilon^e_{ik} + eV^* - \varepsilon^w_{jk'}) n_{ik} \tag{7}$$

where the final state in the collector is assumed to be empty. V^* is the effective applied voltage between the emitter and the well [see Fig. 1(b)], which is considered to be proportional to the total applied voltage V between the emitter and the collector.[2] n_{ik} is the electron occupation number in the emitter, and $T_{ik,jk'}$ is the tunneling matrix element between (ik) states in the emitter and (jk') states in the well. The decay rates of the resonant level in the well due to elastic coupling to the emitter and the collector, $\Gamma^e_{jk'}$ and $\Gamma^c_{jk'}$, are defined by

$$\Gamma^e_{jk'} = \pi \sum_{ik} |T_{ik,jk'}|^2 \delta(\varepsilon^e_{ik} + eV^* - \varepsilon^w_{jk'}), \tag{8}$$

$$\Gamma^c_{jk'} = \pi \sum_{\mathbf{K}} |T_{jk',\mathbf{K}}|^2 \delta(\varepsilon^w_{jk'} + e(V - V^*) - \varepsilon^c_{\mathbf{K}}), \tag{9}$$

where the electronic states in the collector are assumed to be three dimensional (3D), and $\mathbf{K}$ is the 3D wave vector and $T_{jk',\mathbf{K}}$ is the tunneling matrix element between the well and the collector. For 3D collector states, $\Gamma^c_{jk'}$ is a weak function of j and k', and we assume constant $\Gamma^c = \frac{1}{2}\nu_w D_c$ for simplicity with ν_w being the attempt rate for an electron in the well and D_c the collector barrier transmission coefficient.[12] Using the eigenfunctions given by Eqs. (3) and (5), $T_{ik,jk'}$ is written as

$$|T_{ik,jk'}|^2 = \nu_e \nu_w D_e M^2_{ij}(k) \delta_{k,k'} \tag{10}$$

with

$$M_{ij}(k) = \int \phi_j(\widetilde{\ell}_w; y - Y_w) \phi_i(\widetilde{\ell}_e; y - Y_e) dy, \tag{11}$$

where ν_e is the attempt rate in the emitter and D_e is the emitter barrier transmission coefficient. Replacing the δ function in Eq. (7) by a Lorentzian function of characteristic width Γ in order to take into account broadening effects approximately, we have

$$\frac{I}{I_0} = \sum_{ijk} \frac{\gamma^c}{\gamma^e_{jk} + \gamma^c} M^2_{ij}(k) \frac{\Gamma^2}{(\varepsilon^e_{ik} + eV^* - \varepsilon^w_{jk})^2 + \Gamma^2} n_{ik} \tag{12}$$

with

$$\gamma^e_{jk} = \sum_i M^2_{ij}(k) \frac{\Gamma^2}{(\varepsilon^e_{ik} + eV^* - \varepsilon^w_{jk})^2 + \Gamma^2}, \tag{13}$$

where the unit of current is $I_0 = 4e\nu_e\nu_w D_e/\Gamma$ and γ^c is the dimensionless parameter characterizing the charge buildup in the well, defined by $\gamma^c = (\Gamma/\Gamma_r)^2 \propto D_c/D_e$. Here $\Gamma_r = (2\nu_e D_e \Gamma/D_c)^{1/2}$ is the resonance width for charge buildup in the well.[12] Γ is considered to have a contribution from both the intrinsic width and the inhomogeneous broadening in the conduction-band minimum along the length of the wire.

Figure 3 shows the calculated currents as a function of ΔE [$= eV^* - (E^w_0 + \frac{1}{2}\omega_w - E^e_0 - \frac{1}{2}\omega_e)$] in the presence of a magnetic field between $B = 0$ and 6 T. At $\Delta E = 0$ the lowest level in the emitter is aligned to that in the well for $B = 0$. The following parameters were used for the calculation: $\omega_e = 6\,\text{meV}$, $\omega_w = 3.5\,\text{meV}$, $\Gamma = 0.5\,\text{meV}$, and $\Gamma_r = 2.5\,\text{meV}$. The electron density in the emitter is assumed to be independent of B, and the Fermi energy at $B = 0$ is set to 6 meV.

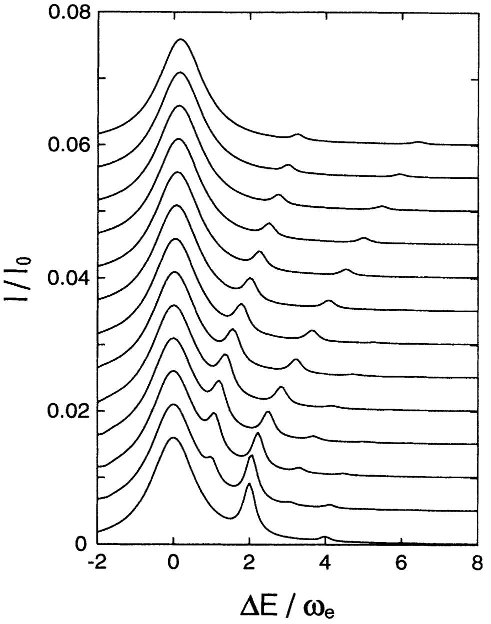

FIG. 3. Calculated results of tunneling currents as a function of effective applied bias ΔE in the presence of a magnetic field between $B = 0\,\text{T}$ (lowest curve) and $B = 6\,\text{T}$ (top curve) in 0.5 T steps. The field is oriented parallel to the current direction. $\Delta E = 0$ corresponds to an applied voltage, which aligns the lowest levels of the emitter and the well for $B = 0$. Parameters are given in the text.

Since k is conserved in the tunneling process, resonant tunneling occurs between states (ik) and (jk) when the voltage across the device is adjusted so that $eV^* = \varepsilon^w_{jk} - \varepsilon^e_{ik}$, and the current is proportional to $M^2_{ij}(k)$. For $B = 0$, electrons are confined electrostatically and the orbit center is located at the middle of the wire regardless of k ($Y_e = Y_w = 0$). For this reason $M_{ij}(k) = 0$ when $(j - i)$ is an odd integer, and tunneling between states with different parity is not permitted. We can therefore associate peaks in $I(V)$ for $B = 0$ with transitions from $i = 0$ to $j = 0, 2, 4, \ldots$. The ratio of square matrix elements M^2_{0j} and M^2_{00} is given by $\eta_j \equiv M^2_{0j}/M^2_{00} = 2^{-2n}(2n)!(n!)^{-2}\{(\omega_w - \omega_e)/(\omega_w + \omega_e)\}^{2n}$ for even integer $j = 2n$ ($n = 1, 2, 3, \ldots$). For the parameters used for Fig. 3 ($\omega_e = 6\,\text{meV}$ and $\omega_w = 3.5\,\text{meV}$), this ratio for $j = 2$ $\eta_2 \approx 0.035$, while the ratio of current peak height $I_2/I_0 \approx 0.5$ (see Fig. 3). This difference is because of the fact that weak structure appears relatively enhanced in the presence of the charge buildup.

For lower magnetic field such that $\omega_c < \omega_w$ (or $B < 2\,\text{T}$ for parameters used for Fig. 3) electrostatic confinement dominates, and the energy levels are almost independent of B. The $j = 2$ resonance is therefore only weakly affected by B for this magnetic field region.

We can understand how the magnetic field violates parity conservation by considering its effect on the electron trajectories. At zero field the electrostatically confined motion in the emitter and the well are both symmetrically positioned with respect to the center of the wire. For $B > 0$ the initial and final orbit centers are displaced, *i.e.*, $Y_e \neq Y_w$. The $j = 1, 3, \ldots$ resonances are therefore observed only for $B > 0$. These peaks have maxima for intermediate magnetic field around $\omega_w < \omega_c < \omega_e$ ($2\,\mathrm{T} < B < 3.5\,\mathrm{T}$).

For a further increase in field $\omega_c > \omega_e$, $\widetilde{\omega}_e \approx \widetilde{\omega}_w \approx \omega_c$, and the voltage separation of the jth resonance from the main resonance is linearly dependent on B as $j\omega_c/e$. In the high magnetic field limit such that $\omega_c \gg \omega_e$, magnetic confinement dominates and $M_{ij}(k) = \delta_{ij}$. $I(V)$ therefore has no peaks other than the main resonance. In the calculated $I(V)$ additional peaks become weaker for high magnetic field as expected. This is, however, not observed in the experimental $I(V)$ — for example the $j = 1$ peak becomes progressively stronger at high field [see Fig. 2(a)]. This behavior is similar to that observed for a large area diode[8] in which the additional peaks are attributed to scattering.

To summarize in this paper we have studied the tunneling current through a quantum wire cavity under a magnetic field applied parallel to the current direction. The current has been calculated within the transfer Hamiltonian formalism. We consider only the case that the lateral confining potential in the emitter, ω_e, is stronger than that in the well, ω_w, and electrons occupy only the lowest level in the emitter.

For $B = 0$ only parity conserving transitions are allowed and resonant peaks have a voltage separation of $2\omega_w/e$ in V^*. In the weak magnetic field for $\omega_c < \omega_w$ (or $B < 2\,\mathrm{T}$ for parameters used for the present work), the voltage position of jth resonance is only weakly affected by B because the electrostatic confinement dominates in this magnetic field region.

For finite magnetic field, parity conservation is broken and the $j = 1, 3, \ldots$ resonances are observed because of the combined effect of the magnetic field and the difference in the confining potentials in the emitter and the well. For high magnetic field $\omega_c > \omega_e$ ($B > 3.5\,\mathrm{T}$) magnetic confinement dominates in the emitter and the well and the voltage separation of the jth peak from the main resonance is linearly dependent on B as $j\omega_c/e$. This is accompanied by vanishing amplitude for $j \geq 1$ resonance.

The magnetic field dependence of the voltage position of jth resonance is in good agreement with the experimental results, while the dependence of the amplitude is different from that observed at higher magnetic field. This may be due to our ignoring scattering in the present study. The characteristics magnetic field scale for interface roughness assisted tunneling is much higher than the magnetic field scale considered in the present study ($\approx$ a few T). For example, if the interface roughness of the lateral decay length $\Lambda \leq 100\,\text{Å}$ is taken into account, the $j = 1$ peaks will have a maximum at $B = 2\hbar/(e\Lambda^2) \geq 13\,\mathrm{T}$.[13] Therefore, we have to take into account interface roughness (or another scattering mechanism), which is not considered in the present work, in order to obtain a quantitative agreement with experiments at high magnetic field.

ACKNOWLEDGMENTS

N.M. and L.E. are grateful to the British Council and EPSRC, respectively, for financial support. This work was supported by the UK EPSRC.

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[10] $\phi_n(\ell; y) = (\sqrt{\pi}2^n n!\ell)^{-1/2} H_n(y/\ell)\exp[-\frac{1}{2}(x/\ell)^2]$ with H_n being the Hermit polynomial.

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