Observations of Magnetoquantized Interface States by Electron Tunneling in Single-Barrier n^-(InGa)As-InP-n^+(InGa)As Heterostructures

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A new type of tunneling phenomenon in forward-biased single-barrier n^+(InGa)As-n^-(InGa)As-InP-n^+(InGa)As heterostructures has been observed when a quantizing magnetic field is applied in the plane of the InP barrier. The occurrence of two distinct series of resonances in the voltage or field dependence of the current is interpreted in terms of electron tunneling from a 2D electron gas in the n^- layer into interfacial Landau states in the n^+(InGa)As. These states correspond to classical skipping orbits of an electron along the tunnel-barrier interface.

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Recent developments in the growth of high-quality semiconductor heterostructures have made a wide range of quantum-mechanical phenomena accessible to experimental study. In particular, the physics of tunneling through almost defect-free heterostructure barriers with nearly flat interfaces is currently of great interest.[1,2] In this Letter we report the observation of a new type of tunneling phenomenon in forward-biased, single-barrier n^+(InGa)As-n^-(InGa)As-InP-n^+(InGa)As structures, which occurs when a quantizing magnetic field ($\mathbf{B} \| Oz$) is applied in the plane of the barrier, perpendicular to the direction of current flow ($\mathbf{J} \| Ox$). As shown in Fig. 1, a two-dimensional electron gas (2DEG) accumulates on the left-hand side (LHS) of the InP barrier and tunneling can occur into the right-hand side (RHS) n^+(InGa)As layer. This picture is consistent with previously reported studies[3] with $\mathbf{B} \| \mathbf{J}$. Most of the applied voltage V is dropped across the accumulation layer and barrier since magnetoconductivity measurements on control samples of n^-(InGa)As show that the voltage drop between the left-hand n^+ contact and the 2DEG is small (< 3 mV at 12 T). The difference between the Fermi levels of the 2DEG and RHS n^+ layer is therefore eV, as in Fig. 1, to a good approximation.

A remarkable feature of the results for $\mathbf{B} \perp \mathbf{J}$ is that, in contrast[3,4] to $\mathbf{B} \| \mathbf{J}$, the tunnel current falls rapidly to zero with increasing magnetic field, as shown in Fig. 2 for a fixed bias voltage of 100 mV. Furthermore, weak oscillatory structure, superimposed on this falloff, is clearly revealed in $\partial^2 I/\partial B^2$. For a wide range of bias voltages there are two distinct series of peaks in $\partial^2 I/\partial B^2$, each periodic in $1/B$, occurring at fields given by $1/B = (n+\phi)/B_f$, where n is an integer, ϕ a field-independent phase factor, and $1/B_f$ is the fundamental periodicity. Note that although double differentiation introduces a phase change of 180° in the oscillatory structure, it does not change the period in $1/B$. As shown in Fig. 3, for both series (labeled plus and minus), the fundamental fields $B_{f\pm}$ increase steadily with voltage. These data

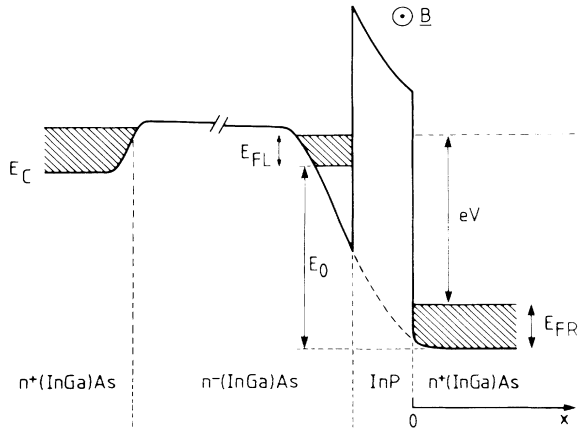

FIG. 1. Spatial variation of potential energy $E_c(x)$ of an electron at the conduction-band edge under forward bias.

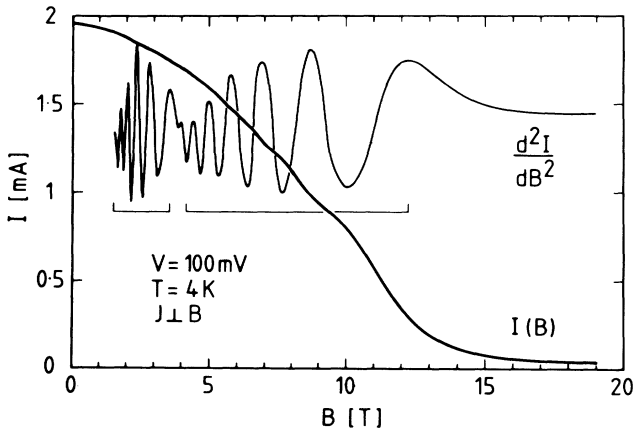

FIG. 2. Variation of current I and $\partial^2 I/\partial B^2$ with B for forward-bias voltage $V=100$ mV. Mesa diameter $=200$ μm. The two series are distinguished by horizontal brackets.

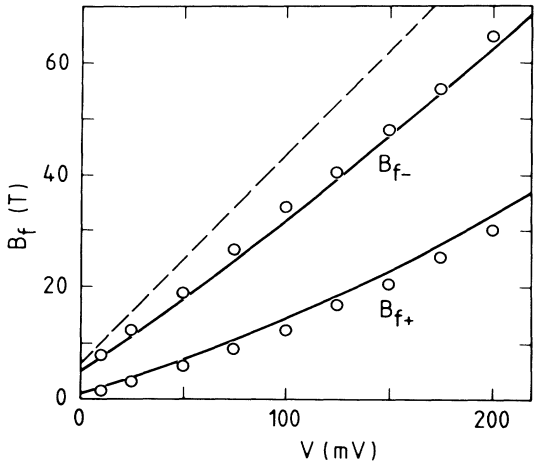

FIG. 3. Variation of B_f with bias voltage V for each series. Experiment: circles. Theory: bulk Landau states, dashed line; interfacial Landau states, solid lines.

were obtained at 4 K for a metalorganic chemical-vapor deposition-grown device with barrier thickness 170 Å and n^-(InGa)As layer of thickness 0.8 μm, doped in the low 10^{15}-cm^{-3} range. Both n^+ layers were doped to 10^{17} cm^{-3}. We have obtained similar results for a range of such molecular-beam epitaxy- and metalorganic chemical-vapor deposition-grown structures with mesa diameters in the range 100–600 μm.

We stress that these results cannot be explained in terms of tunneling into bulk Landau levels in the RHS n^+ layer. If this were the case, a single series of resonances would occur when the Fermi level in the 2DEG coincided with a Landau level on the RHS, which, from Fig. 1, corresponds to

$$eV+E_{\rm FR}=(n+\tfrac{1}{2})\hbar\omega_c, \quad n=0,1,2,\ldots, \tag{1}$$

where $\omega_c=eB/m^*$ and $E_{\rm FR}=eV_{\rm FR}$ is the Fermi level in the n^+ layer. This indeed predicts structure periodic in $1/B$ but gives $B_f=(V+V_{\rm FR})m^*/\hbar$, which does not give a fit to either series for any reasonable choice of effective mass m^*. The dashed line in Fig. 3 corresponds to $m^*=0.043m_e$ which is appropriate to an energy $\simeq E_{\rm FR}=17$ meV in the conduction band. Taking account of nonparabolicity[5] causes this line to deviate further from the data since m^* increases with energy eV.

To understand the origin of the two series we consider the effect of the magnetic field on the electron wave function ψ near the tunnel barrier. Using the vector potential in the Landau gauge $\mathbf{A}=(0,Bx,0)$, we may write $\psi(x,y,z)=\exp[i(k_y y+k_z z)]\psi(x)$. The function $\psi(x)$ obeys a one-dimensional Schrödinger equation[6] containing the potential energy $E_c(x)$ at the conduction-band edge shown in Fig. 1, and a simple harmonic oscillator (SHO) potential $\frac{1}{2}m^*\omega_c^2(x-X)^2$, centered at $X=-\hbar k_y/eB$. The dependence of center coordinate X on transverse momentum $\hbar k_y$ is due to the action of the Lorentz force on the electron motion. In the limit of a very high impenetrable barrier the eigenstates on either side can be considered separately. In the RHS n^+ region we take $E_c(x)$ to be constant.[3] If we take the origin $x=0$ to be the RHS of the barrier, $\psi(x)$ is an eigenstate of the SHO potential with $\psi(0)=0$. We write the energy levels

$$E_{\rm R}(k_y,k_z)=E_n(k_y)+\hbar^2k_z^2/2m^*.$$

For large X, when the electron is far from the interface, $\psi(x)$ is a standard SHO eigenstate centered at $x=X$ and we have bulk Landau levels, $E_n=(n+\frac{1}{2})\hbar\omega_c$. For small X, when the electron can interact with the interface, the SHO states are distorted and E_n increases as X decreases (k_y increases) as shown in Fig. 4. These are interfacial Landau states. In a semiclassical picture, bulk states correspond in the xy plane to closed circular orbits, interfacial states to skipping orbits which intersect the interface and X is the orbit center. If v is the orbital speed, then $E_n=\frac{1}{2}m^*v^2$ and the orbit radius is $r=v/\omega_c=(2E_n/m^*\omega_c^2)^{1/2}$. The condition for a bulk state, $X\geq r$, is thus

$$X\geq\left(\frac{2E_n}{m^*\omega_c^2}\right)^{1/2}, \quad \text{or} \quad -\hbar k_y\geq(2m^*E_n)^{1/2}. \tag{2}$$

Equation (2) defines the boundary between interfacial and bulk states and is shown in Fig. 4 as a dashed curve. Using the WKB approximation, we have calculated $E_n(k_y)$ for the interfacial states and, with an additional algebraic approximation (accurate to within 10%), find the following analytic expression:

$$\frac{1}{2}\left(1-\frac{\hbar k_y}{(2m^*E_n)^{1/2}}\right)E_n(k_y)=(n+\tfrac{3}{4})\hbar\omega_c, \tag{3}$$

$$n=0,1,2,\ldots.$$

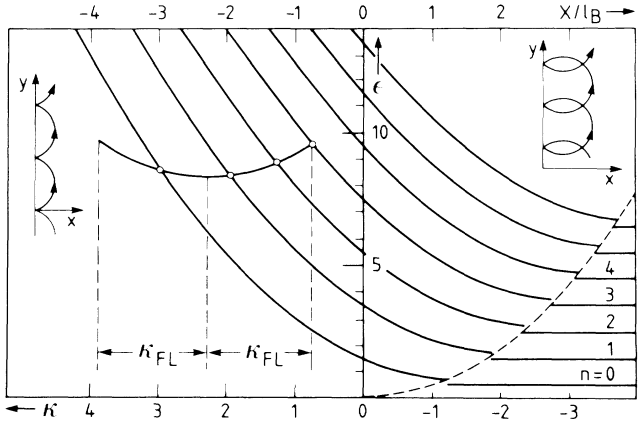

FIG. 4. Dispersion relations for interfacial Landau states in terms of scaled variables $\epsilon = E_n/\hbar\omega_c$ and $\kappa = k_y l_B = -X/l_B$, where $l_B = (\hbar/eB)^{1/2}$. Insets: Classical skipping orbits for interface states with $k_y < 0$ and $k_y > 0$. The parabola corresponds to $E = E_0 + \hbar^2(k_y - k_0)^2/2m^*$, $-k_{FL} < k_y - k_0 < k_{FL}$, and is drawn for $V = 100$ mV and $B = 5$ T. Intersection points $k_y(n)$ are shown as open circles. The dashed line marks the boundary between bulk and interfacial Landau states.

In the 2DEG (LHS of barrier), the magnetic field only slightly perturbs[6] $\psi(x)$, which is a bound state of the accumulation-layer potential. The energy of this state, including transverse motion parallel to the barrier, is

$$E_L(k_y,k_z) = E_0 + \hbar^2(k_y'^2 + k_z^2)/2m^*.$$

E_0 is the bound-state energy, which is only weakly field dependent,[6] $k_y' = k_y - k_0$, $k_0 = eB(b + a_0)/\hbar$, b is the barrier width, and a_0 is the mean distance of a bound electron from the LHS of the barrier. The states of transverse motion are occupied up to the Fermi energy $E_{FL} = \hbar^2 k_{FL}^2/2m^*$. The shift $\hbar k_0$ is just the momentum change due to the Lorentz force as an electron traverses the barrier region.

For a barrier of finite height the eigenstates on the two sides are no longer independent since tunneling can occur. In the transfer-Hamiltonian approach,[7] tunneling is described in terms of transitions between LH and RH states. Under forward bias a net current flows due to transitions from the occupied states of the 2DEG into empty Landau states in the n^+ region. Current continuity through the device is maintained by scattering processes which allow the electron to diffuse to and from these states. For a plane barrier the transverse components of momentum k_y, k_z and energy are conserved in tunneling.[7] Hence $E_L(k_y,k_z) = E_R(n,k_y,k_z)$, which gives

$$E_0 + \hbar^2 k_y'^2/2m^* = E_n(k_y). \quad (4)$$

This condition is most easily interpreted graphically by our looking for intersections in the E-k_y plane of the parabola $E = E_0 + \hbar^2(k_y - k_0)^2/2m^*$ with the set of dispersion curves $E = E_n(k_y)$ of the Landau states. As shown in Fig. 4 this gives a discrete set of k_y values, $k_y(n)$, each of which corresponds to a group of electrons which contribute to the current. As B or V is varied the parabola shifts relative to the dispersion curves and intersection points $k_y(n)$ enter or leave the extremities of the parabola. The consequent changes in the number of contributing electrons account for the periodic structure in $J(V,B)$. Moreover, the two series of different periodicity are naturally associated with the two extremities where $k_y' = \pm k_{FL}$. For these electrons the energy conservation condition (4) gives $E_n(\pm k_F + k_0) = E_0 + E_{FL} = eV + E_{FR}$, the latter equality following directly from Fig. 1. For coupling to bulk levels this relation reduces to Eq. (1). For coupling to interface states we use the WKB energies given by Eq. (3). Substituting $E_n = eV + E_{FR}$ and $k_y = \pm k_{FL} + k_0$ gives

$$(V + V_{FR})(m^*\xi_\pm/\hbar B) = n + \tfrac{3}{4} + \phi', \quad (5)$$

where $\xi_\pm = \frac{1}{2} \mp \frac{1}{2}\hbar k_{FL}/\{2m^*e(V + V_{FR})\}^{1/2}$ and $\phi' = \{2m^*e(V + V_{FR})\}^{1/2}(b + a_0)/4\hbar$ is field independent. Equation (5) is of the required form $B_{f\pm}/B = n + \frac{3}{4} + \phi'$ and the fundamental fields of the $\pm k_{FL}$ series are

$$B_{f\pm} = (\xi_\pm m^*/\hbar)(V + V_{FR}).$$

To calculate $\xi_\pm$, we obtain $k_{FL} = (2\pi n_s)^{1/2}$ from the areal electron density n_s in the 2DEG, which is controlled by the applied voltage and capacitance per unit area $C = n_s e/(V - V_{fb})$, where V_{fb} is the flat-band bias. Using[3] $C = 4250$ μF m^{-2}, $m^* = 0.043 m_e$ and neglecting charge in the barrier ($V_{fb} = -V_{FR}$) gives $\xi_+ = 0.31$, $\xi_- = 0.69$. This approximation results in linear $B_{f\pm}$-V relations whose slopes are both much less than would result from coupling to bulk levels for which $\xi = 1$. However, the experimental results (Fig. 2) show significant deviations from linearity, which are attributable primarily to nonparabolicity of the conduction band and also variation of capacitance with bias voltage. We have therefore used a voltage-dependent capacitance obtained directly from experiment[3] and an effective mass $m^*/m_e = 0.04(1 + \alpha E)$ dependent on electron energy E. As shown in Fig. 3, a good fit to experiment is then obtained with $\alpha = 1.3 \times 10^{-3}$ meV^{-1}. The coefficient α is rather smaller than found from optical measurements[5] ($\alpha = 5 \times 10^{-3}$ meV^{-1}, $E < 20$ meV) but is reasonable[2] considering the wide energy range investigated (40–220 meV). From the minimum field at which oscillations in $\partial^2 J/\partial B^2$ are observed ($\simeq 1.7$ T at 100 mV) we estimate a length $\simeq 0.3$ μm between collisions with interface for the corresponding classical skipping orbit. Thus a detectable proportion of the electrons have this rather long ballistic trajectory. The corresponding time $\simeq 0.3$ ps is comparable with LO-phonon emission times for hot electrons in III-V materials.[8]

It is interesting to note why tunneling into bulk Lan-

dau levels does not occur. To be launched into the closed bulk orbit, which just grazes the interface, an electron must have a negative transverse momentum component (see Fig. 4) $\hbar k_y = -m^* v = -(2m^* E_n)^{1/2}$, in agreement with Eq. (2). Using the extremal value $k_y = -k_{FL} + k_0$, we see that coupling to bulk orbits requires

$$\hbar k_{FL} - \hbar k_0 > (2m^* E_n)^{1/2}.$$

But in typical semiconductor heterostructures the charge in the 2DEG, and hence E_{FL}, is limited by the small capacitance and, as in Fig. 1, $E_{FL} \ll eV + E_{FR} = E_n$ at all biases. Hence $\hbar k_{FL} \ll (2m^* E_n)^{1/2}$ and the above inequality cannot be satisfied even when the momentum shift $\hbar k_0$ due to the Lorentz force is small. By contrast, in a metal-oxide-semiconductor structure, the large and voltage-independent Fermi momentum in the metal allows observation of tunneling into bulk Landau levels in the semiconductor.[9]

In conclusion, we have observed oscillatory structure in the current due to electrons tunneling from a 2DEG into n^+(InGa)As, in a magnetic field parallel to the plane of the tunnel barrier. This is attributed to Landau-level structure in the n^+ material. When conservation of energy *and momentum* are taken into account in the tunneling process, we find that the electrons tunnel into interfacial Landau states (skipping orbits).

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