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Periodic negative conductance by sequential resonant tunneling through an expanding high-field superlattice domain

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Strikingly periodic negative-differential-resistance oscillations are observed in doped GaAs/$Al_xGa_{1-x}As$ superlattices. These are produced by well-to-well sequential resonant tunneling through weakly coupled localized states.

Extensive recent research on tunneling[1] through GaAs/$Al_xGa_{1-x}As$ single,[2] double,[3-9] and triple[10] barrier structures as well as p-n junction superlattices[11] has clarified the mechanisms of current transport and negative differential resistance. However, only a few experiments[12,13] have been performed on tunneling through unipolar multi-quantum-well superlattices. Esaki and Chang's[12] pioneering tunneling investigation demonstrated miniband conduction through thin-barrier (40 Å AlAs) superlattices. Because of these thin barriers, there is strong coupling between the wells producing a large ground-state bandwidth ($\Delta E_1 = 5$ meV) and thus electron transport by miniband conduction. In this strong-coupling regime they observed a novel oscillatory negative conductance due to the formation of an expanding high-field domain produced by the electric-field-induced breaking of the miniband conduction. The voltage drop across this domain aligns the ground state of one well with the first excited state of the neighboring well allowing resonant tunneling to occur.

In contrast to this strongly coupled miniband conduction regime, we have investigated wells which are weakly coupled through relatively thick barriers. The ground-state bandwidth ($\Delta E_1 = 0.4$ meV) for the quantum-well superlattice we have studied is over an order of magnitude narrower than those investigated by Esaki and Chang.[12] Thus, the quantum-well states will be localized[14] (by well-width fluctuations and also by the applied voltage) and therefore electron transport will be dominated by sequential resonant tunneling[14,15] from well to well. In spite of this important difference in transport mechanism, we demonstrate, for the first time, that strikingly regular negative conductance oscillations can be produced by sequential resonant tunneling through an expanding high-field superlattice domain.

The crystals were grown on a computer-controlled vacuum generator molecular-beam-epitaxy system using a substrate temperature of 600 °C. The superlattice consists of 49 periods of 76 ± 2 Å GaAs wells separated by 88 ± 2 Å $Al_{0.27}Ga_{0.73}As$ tunneling barriers. These thicknesses were determined by high-resolution x-ray scattering.[16] To reduce the possibility of interface states, only the center 56 Å of the wells are doped with Si ($n = 3 \times 10^{17}$ cm^{-3}). Contact layers n^+ doped to 8×10^{17} cm^{-3} were grown under (1 μm) and over (0.5 μm) the superlattice. The device is fabricated using standard photolithographic techniques. At first, a mesa of area 2×10^{-5} cm^2 is etched to expose the n^+-type GaAs underneath the superlattice, then Au-Ge alloy contact is evaporated on the top of the mesa and the area adjacent to the mesa to form a two-terminal device. The current-voltage (I-V) characteristics were measured with a HP4145 parameter analyzer from 5 to 300 K (with the substrate biased negatively).

At room temperature, there are no obvious structures in the I-V characteristics. However, below 220 K, negative differential resistance emerges and at lower temperatures, the current peaks can be clearly identified. Figure 1 shows the transport characteristics at 20 K. From this figure, 48 negative-conductance features can be identified with strikingly regular voltage spacings of 85 ± 18 mV. In order to understand the origin of these structures, we first investigated the transport mechanism at lower bias where the I-V curve is approximately linear. Figure 2 shows the temperature (T) dependence of the current at an applied voltage (V_a) across the entire superlattice equal to 0.1 V. The current first decreases with decreasing T and then remains approximately constant below 220 K. Above 220 K the T dependence is consistent with both thermionic emission and optical-phonon-assisted tunneling.[17] However, the latter predicts a mobility μ of 1.5×10^{-3} cm^2/V s with the

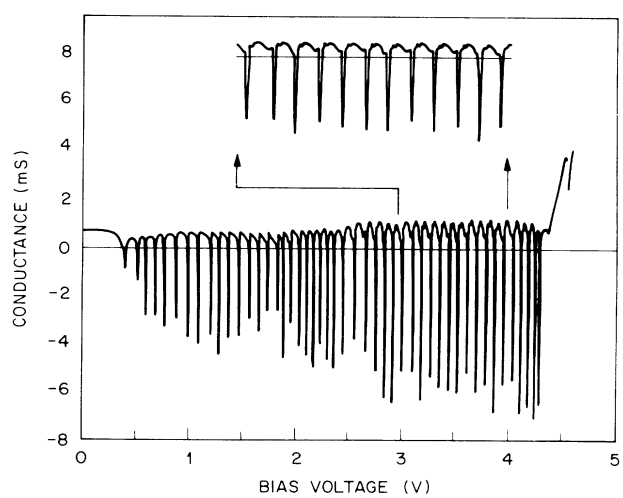

FIG. 1. Conductance vs applied voltage for a 49-period superlattice. The inset is an expanded view of the 3–4-V range, showing the regular periodicity of the negative conductance peaks.

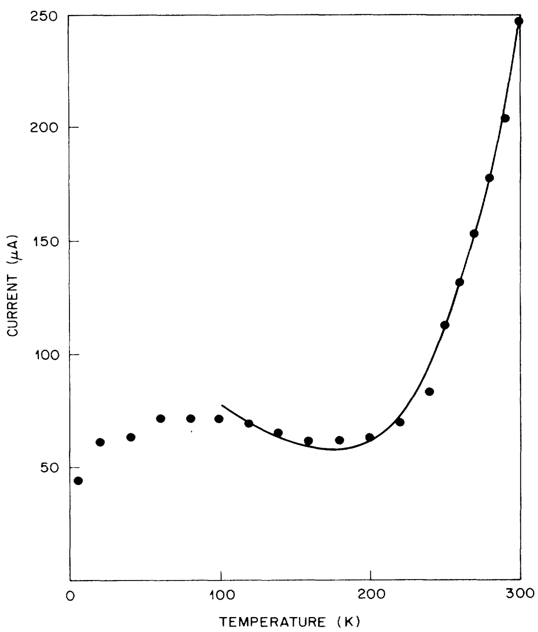

FIG. 2. Current (measured at an applied voltage of $V_a = 0.1$ V) plotted against temperature. The points are obtained from experiment, while the solid line is theory [Eq. (1)].

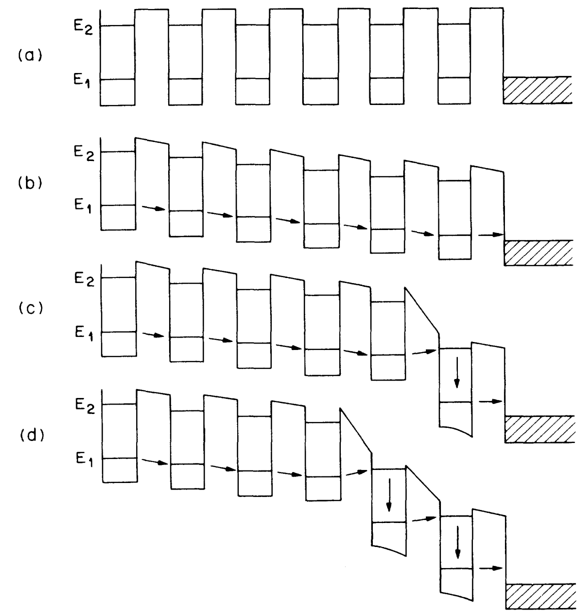

FIG. 3. Schematic band diagram of the superlattice for several values of the average potential drop per period V. (a) Zero bias. (b) Sequential resonant tunneling through the ground state E_1 for $V < \eta\hbar/\tau_1$; arrows indicate electron transport. (c) Formation of first high-field domain for V slightly greater than $\eta\hbar/\tau_1$. Sequential resonant tunneling occurs through E_1 in the low-field region and through the first excited state E_2 in the high-field region. (d) Expansion of the high-field domain by one additional quantum well for a voltage increase of ΔV.

present device parameters at room temperature, which is two orders of magnitude smaller than the measured value of $\mu = 0.37$ cm^2/V s. Therefore, optical-phonon-assisted tunneling can be neglected in the analysis. Below 220 K the current is insensitive to T. However, it is not due to band conduction since the applied field $F = 1.2$ kV/cm is five times the localization field.[14,18] In fact, Ohmic conduction continues until $F = 4.3$ kV/cm, 18 times the localization field. Instead, the current transport is due to ground-state resonant tunneling.[15] Further evidence for this is that our low-temperature ($T = 5$ K) mobility $\mu = 0.12$ cm^2/V s is nearly two orders of magnitude smaller than the miniband conduction value[12] $\mu = 50$ cm^2/V s.

In principle, due to the two-dimensional (2D) nature of the electron gas in the well, resonant tunneling is possible only when the energy levels in each well coincide, a condition generally not fulfilled in the presence of a field. However, Kazarinov and Suris[14] showed that in the presence of acoustic phonons and impurity scattering within each well, conservation of energy and momentum is relaxed and resonant tunneling is possible provided that $eV \ll \hbar/\tau_1$, where V is the potential difference between the adjacent wells and τ_1 is the ground-state scattering time. Therefore, at small bias, the electrons are able to conduct by ground-state resonant tunneling through the ground states of each well [Fig. 3(b)]. The first negative-differential-resistance peak occurring at $V_a = 0.35$ V then indicates the disruption of the resonant tunneling when $eV = \eta\hbar/\tau_1$, where η is a constant of order unity [Fig. 3(c)]. Assuming the potential drop across the device is uniform at small bias, the corresponding V ($\simeq V_a/50$) is 7 meV which gives τ_1 an approximate value of 10^{-13} s. Thus, for $V_a < 0.35$ V when the ground-state resonant tunneling condition is still satisfied, the total current can be written as

$$I = \frac{eA}{\hbar L^2} D_0 kT \ln\left(\frac{1+e^{E_F/kT}}{1+e^{(E_F-eV)/kT}}\right) + \frac{e^2 m^*}{\pi\hbar^2}\frac{A}{L} v_D t V \exp[-(H-E_1-E_F)/kT] \ , \quad (1)$$

where the first term is due to tunneling and the second term to thermionic emission. In the above equation,

$$D_0 = \exp\left\{\frac{-4L_b}{3eV\hbar}(2m_b^*)^{1/2} \times [(H-E_1)^{3/2} - (H-E_1-eV)^{3/2}]\right\} \quad (2)$$

is the tunneling coefficient of individual trapezoidal barriers in the WKB approximation, A is the area of the device, e is the electronic charge, L is the well thickness, E_F is the T-dependent Fermi level, E_1 is the ground-state energy, m^* is the effective mass of GaAs, v_D is the drift velocity near the top of the barrier, t is the quantum-mechanical transmission coefficient of an electron near the top of the barrier, V is the average potential drop per period, H is the barrier height, and in Eq. (2) m_b^* is the effective-mass ratio of the barrier to the free electron, and L_b is the barrier width. Using H, L_b, and $v_D t$ as fitting parameters, we achieve reasonable agreement with the experimental current versus temperature curve (measured using $V_a=0.1$ V) as shown in Fig. 2. The small discrepancy appearing at $T<100$ K is due to carrier freeze out which we have neglected. The freeze-out temperature observed is in agreement with theory.[19] When the data are taken at an increasing T as in the present case, carrier freeze-out would be significant when kT is less than $2R_0$ (~ 120 K) for our device parameters, where R_0 is the effective Rydberg. From the fit, we obtain $H=209\pm 10$ meV, $L_b=90\pm 2$ Å, and $v_D t=1.8\times 10^6$ cm/s which agree quite well with $H=202$ meV from the 60% band discontinuity rule[20] and $L_b=88$ Å from the x-ray analysis. We have also performed a computer calculation of t using our device parameters, obtaining $t\simeq 6\%$ near the top of the barrier. This gives $v_D\approx 10^7$ cm/s at $F=1.2$ kV/cm.

With the information obtained from small fields, we are now in a position to explain the structures at high fields. As mentioned earlier, when the potential drop V across a period is larger than $\eta\hbar/\tau_1$, ground-state resonant tunneling is not possible and as a result negative differential resistance occurs. As each period breaks off from the resonant condition, the resistance across this period becomes much larger and a high-field domain forms. Any subsequent increase in the bias will appear across this domain until the ground level rises to within $\eta\hbar/\tau_2$ of the first excited level (E_2) of the next well where upon the resonant tunneling condition is restored [Fig. 3(c)]. Further increases in bias will cause another well to break off from the resonant condition and the I-V characteristic repeats [Fig. 3(d)]. Due to the screening effect of the space-charge buildup[21] at large bias, the domain formation is not a random process but occurs first at the anode and then extends one after the other towards the cathode. As a result, one would expect there to be $p-1$ negative conductance peaks for a device with p periods. Indeed in Fig. 1, 48 oscillations can be identified, in agreement with this interpretation.

In order to further substantiate this argument, we fabricated several devices having varying periods by etching mesas of different depths. As expected, the devices with fewer periods had proportionately fewer negative resistance peaks, which corresponded exactly to the initial oscillations in the original superlattice. This demonstrates that the domain formation is not random, but forms at the anode and moves progressively toward the cathode. For fields large enough so that the high-field domain encompasses the entire superlattice, we reach the limit of sequential resonant tunneling through the whole superlattice which was discussed in detail by Capasso, Mohammed, and Cho.[15]

From the above discussion the voltage difference between the oscillations is given by $\Delta V=(E_2-E_1-\eta\hbar/\tau_1-\eta\hbar/\tau_2)=85\pm 18$ meV. Using the value $\eta\hbar/\tau_1=7$ meV obtained previously and $\eta\hbar/\tau_2=11$ meV (determined below), we deduce an experimental value for the intersubband spacing of $(E_2-E_1)=103\pm 18$ meV. In order to compare with this measurement we have done a calculation of E_2-E_1 using a band-offset parameter[20] $\Delta E_c/\Delta E_g=0.60$ and a five-band nonparabolicity parameter[20] $\gamma=4.9\times 10^{-15}$ cm^{-2} for the well ($\gamma=0$ for the barrier). We obtain a theoretical value of $E_2-E_1=105$ meV in good agreement with our experiments.

As an additional check on this value of E_2-E_1, we performed infrared absorption experiments[22,23] to directly measure the intersubband energy, and to give an estimate of the excited-state lifetime τ_2. Because of the relatively low doping and consequently weak optical absorption we used a multipass waveguide geometry[23] to enhance the absorption (see inset in Fig. 4). The intersubband peak is clearly seen at $E_2-E_1=820\pm 10$ cm$^{-1}=102\pm 1$ meV in excellent agreement with both the calculated and tunneling determinations. As shown in Fig. 4, the line shape is well fitted by a Lorentzian (the deviation on the low-energy side is due to substrate absorption). The full width at half maximum Lorentzian linewidth $\Delta\nu=2\hbar/\tau_2=86$ cm$^{-1}=11$ meV corresponds to a lifetime of $\tau_2=1.2\times 10^{-13}$ s. This is the value used above, to determine the intersubband spacing from the I-V oscillations.[24] A further independent confirmation of this value of E_2-E_1 was obtained from photoluminescence excitation spectroscopy experiments which measured the transition energies from the heavy-hole states to the $n=1$ and $n=2$ electron levels. This yielded $(E_2-E_1)=106\pm 2$ meV (Ref. 25) in good agreement with the other experimental and theoretical determinations. We can also determine the oscillator strength $f=4\pi m^*\nu\langle z\rangle^2/\hbar$ and transition dipole moment $\langle z\rangle$ of this intersubband transition from the integrated absorption strength;[22,23] the results are $f=0.5\pm 0.2$ and $\langle z\rangle=17\pm 4$ Å. The large oscillator strength, of order unity, indicates that these structures may be useful for infrared detectors.[26-30] We have also theoretically calculated $\langle z\rangle$ using the same quantum-well parameters and find $\langle z\rangle=22$ Å in good agreement with experiment. These large values for f and $\langle z\rangle$ are consistent with previous work.[22,23]

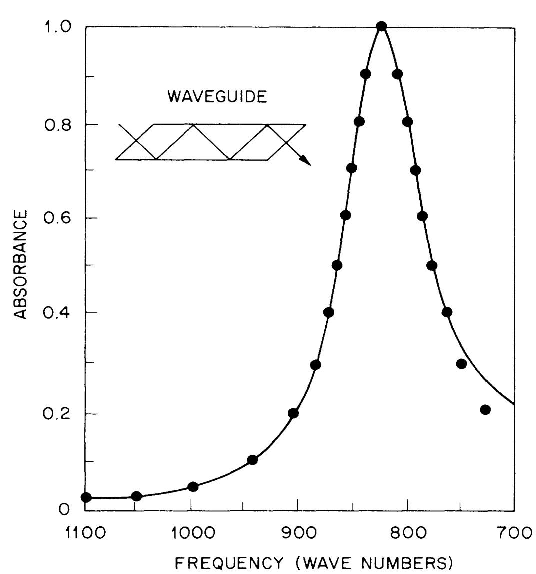

FIG. 4. Infrared optical absorbance spectrum measured (smooth curve) in the multipass waveguide geometry shown in the inset. The dots are a Lorentzian fit with a full width half maximum linewidth of $\Delta\nu=86$ cm$^{-1}=11$ meV.

In conclusion, we have observed strikingly regular negative-conductance oscillations due to resonant tunneling through an expanding high-field superlattice domain. These periodic negative-resistance peaks are produced by well-to-well sequential resonant tunneling through thick barriers into localized states. This is in contrast to previously seen oscillations which are observed in the opposite strongly coupled miniband conduction regime for superlattices having thin barriers.

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