[8] B. N. Harmon and A. J. Freeman, Phys. Rev. B 10, 1979, 4849 (1974); B. N. Harmon, private communication.

[9] Spin-orbit and exchange operators do not commute; hence spin is no longer a good quantum number. Spin-orbit effects generally lift the degeneracies of majority- and minority-spin bands.

[10] See Ref. 2 for a good review of the work by Freeman, Watson, and Dimmock, and other workers. See also the original papers by C. Jackson, Phys. Rev. 178, 949 (1969) (Tb); S. C. Keeton and T. L. Loucks, Phys. Rev. 168, 672 (1968) (Gd, Dy, Er, Lu); R. W. Williams, T. L. Loucks, and A. R. Mackintosh, Phys. Rev. Lett. 16, 168 (1966) (Ho).

Direct Observation of Superlattice Formation in a Semiconductor Heterostructure

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(Received 17 March 1975)

We demonstrate, via low-temperature optical-absorption measurements on ultrathin, coupled potential wells in molecular-beam–grown $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$-GaAs heterostructures, the evolution of resonantly split discrete well states into the lowest band of a one-dimensional superlattice. Both electron and hole superlattices appear to be practical.

The evolution of molecular-beam epitaxy[1] as a technique for the growth of ultrathin layers of high-quality III-V–semiconductor single crystals has allowed access to a new regime of quantum effects in structures approaching atomic dimensions. Quantum states of electrons[2,3] and holes[2] in single potential wells of GaAs bounded by thick $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ barriers have been observed in tunneling[3] as well as in optical absorption[2] and stimulated emission.[4] Coupling between wells through thin penetrable barriers is expected to split the bound quantum states[5] into symmetrical and antisymmetrical combinations.[6] In the limit of superlattice formation, multiple energy gaps occur in the Brillouin zone, and new and useful transport properties are anticipated. Tunneling measurements in AlGaAs-GaAs superlattices have been reported,[7] although effects due to the bound-state

coupled wells is increased from one to ten, we are able to present unequivocal evidence for the tunneling of electrons and holes through the $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ barriers. Structures with ten or more coupled wells appear to approximate the superlattice regime, whereas structures with fewer wells are well described in terms of interacting single wells. The experimental data are interpreted with an exact solution of the Schrödinger equation for transmission through multiple rectangular potential barriers.

A series of structures, with GaAs well widths in the range $50\ \text{Å} < L_z < 200\ \text{Å}$ and $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ barrier widths in the range $12\ \text{Å} < L_B < 18\ \text{Å}$ (8–12 monolayers), were grown by molecular-beam epitaxy on GaAs substrates with use of a previously outlined procedure.[2] Al concentrations in the range from $x = 0.19$ to 0.27 were studied. At

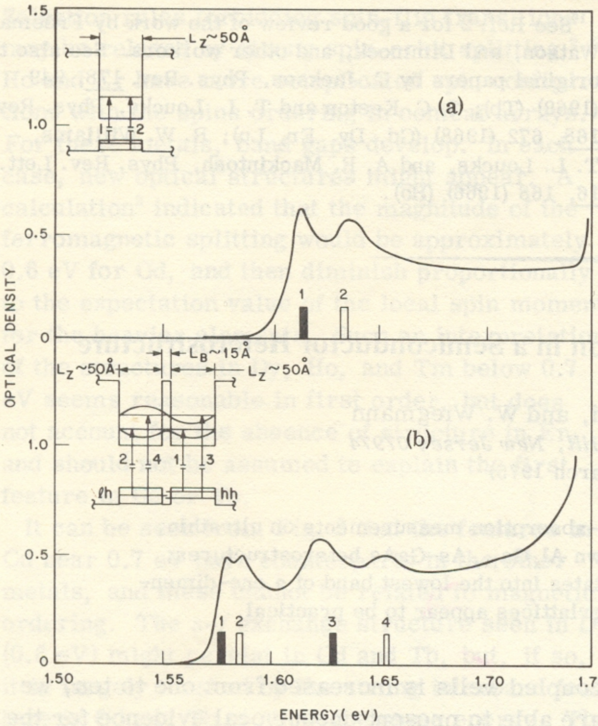

FIG. 1. (a) Optical-density spectrum of a series of eighty single GaAs wells isolated by thick (~ 180 Å) $Al_{0.27}Ga_{0.73}As$ barriers. Peaks 1 and 2 correspond to exciting an electron from $n=1$ heavy-mass and light-mass valence-band bound states, respectively, to the $n=1$ conduction-band bound state, as shown in inset. Calculated predictions of peak positions for rectangular wells, based on growth parameters, are shown on abscissa. Black bars refer to heavy holes, white bars to light holes. (b) Spectrum of a series of sixty double GaAs wells coupled through thin (~ 15 Å) $Al_{0.19}Ga_{0.81}As$ barriers. Both the $n=1$ hole and electron bound states are split by resonant coupling through the penetrable barriers, with symmetric (bonding) combinations of single-well states closest to the well bottoms. Each pair of wells is isolated from the adjacent pairs by 206 Å of $Al_{0.19}Ga_{0.81}As$ barrier. Light- and heavy-hole states extend through both wells.

$\Delta E = E_g^{Al_xGa_{1-x}As} - E_g^{GaAs} = \Delta E_{CB} + \Delta E_{VB}$. Values of $\Delta E_{CB}/\Delta E = 0.85 \pm 0.03$ and $\Delta E_{VB}/\Delta E = 0.15 \pm 0.03$ were used for the conduction-band and valence-band potential-barrier heights.[2] At the Al concentrations used, ΔE is proportional to x and equal to 250 meV at $x = 0.20$.

In Fig. 1(a) we present optical-absorption spectra (2 K) from a series of eighty rectangular GaAs wells of width $L_z = 50 \pm 2$ Å, interleaved by $Al_xGa_{1-x}As$ layers ~ 180 Å thick. At the operating temperature (2 K), tunneling of either bound of remarkably reproducible thin potential wells and barriers, essentially rectangular and uniform to the order of a monolayer, can be created with molecular-beam epitaxy. The coupling behavior of the wells proves that synthetic superlattices can indeed be created. The molecular-beam-epitaxy technique for fabrication and the optical technique for energy-level determination should be applicable to additional configurations and compositions of interest for both basic and applied studies.

We wish to acknowledge L. Kopf for the determination of the Al content of the layers and for expert technical assistance. We thank C. H. Henry and M. B. Panish for helpful discussions.

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[10] We wish to thank G. A. Baraff for providing the computer program which we used to solve the Schrödinger equation for penetration of arbitrary well and barrier profiles.

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ARTICLES

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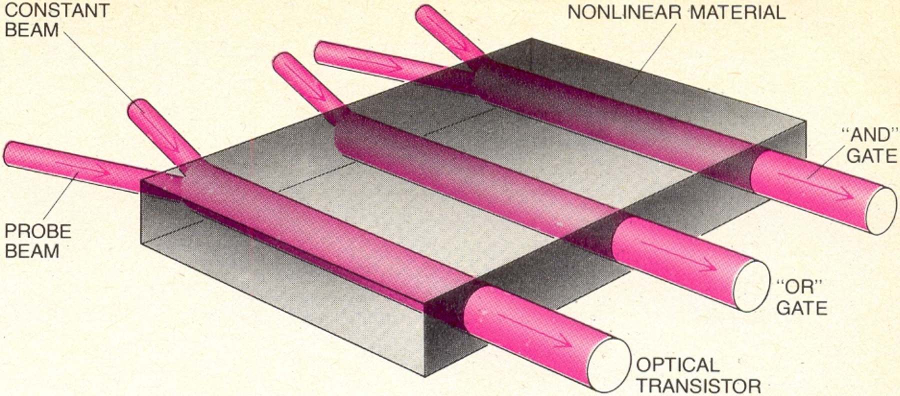
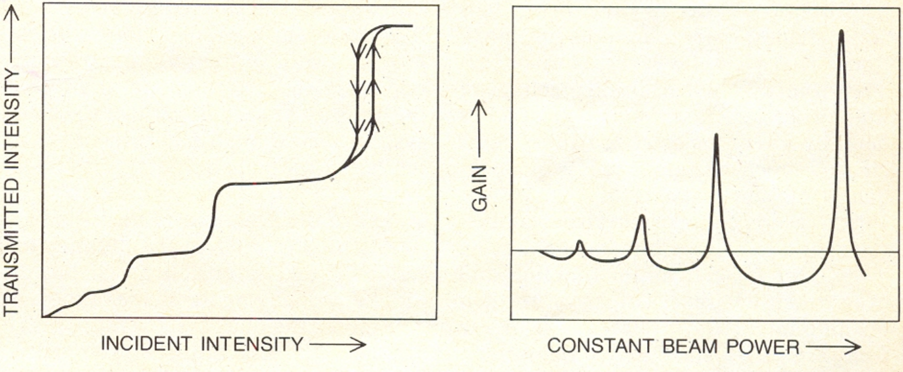


PARALLEL PROCESSING is a capacity of optical switches that could lead to new designs and new capabilities in the computer. Multiple laser beams can be focused so that they remain separate as they pass through a crystal of nonlinear material. Each radiation path could serve as the site of a separate operation, and the operations carried out on the various beams could be different. For example, if there were three paths, one could be an AND gate, one an OR gate and one a transistor. If three incident beams came from one beam that had been split, three operations at once could be done on the original signal. Such a capacity would require a new form of information processing; in electronic computers one operation at a time is done on a signal.



MULTIPLE BISTABILITY is a property of optical switches that could lead to the

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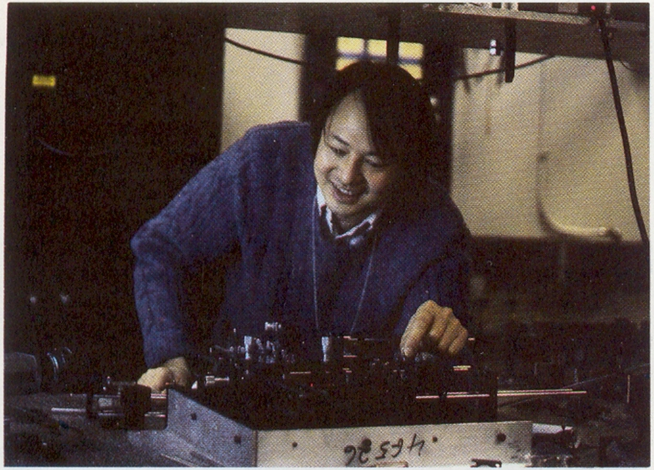
HIDDEN VISUAL PROCESSES

EYE TORSION AND VISUAL TILT ARE M



AT&T BELL LABORATORY

At AT&T Bell Labs, Alan Huang (above) continues his effort to build the first full-blown optical digital computer. Bell researchers David Miller (right) and Jill Henry (far right) are pioneering the self-electrooptic effect devices (SEEDs) that will be the basis for optical gates used in the system.

AT&T BELL LABORATORY

Band-Edge Electroabsorption in Quantum Well Structures: The Quantum-Confined Stark Effect

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(Received 27 April 1984)

We present theory and extended experimental results for the large shift in optical absorption in GaAs-AlGaAs quantum well structures with electric field perpendicular to the layers. In contrast to the Stark effect on atoms or on excitons in bulk semiconductors, the exciton resonances remain resolved even for shifts much larger than the zero-field binding energy and fields > 50 times the classical ionization field. The model explains these results as a consequence of the quantum confinement of carriers.

PACS numbers: 78.20.Jq, 42.80.Ks, 71.35.+z, 73.40.Lq

When semiconductors are fabricated in very thin layers (e.g., ~ 100 Å), the optical absorption spectrum changes radically as a result of the quantum confinement of carriers in the resulting one-dimensional potential wells.[1] In such multiple quantum wells of GaAs, sandwiched between barrier layers of $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ of thickness sufficient to prevent significant coupling between adjacent GaAs layers, the confinement changes the absorption spectrum from the smooth function of bulk material to a series of steps. Additionally, the confinement also increases the binding energy of excitons, resulting in exceptionally clear exciton resonances at room temperature in GaAs-AlGaAs quantum wells.[2-5]

When electric fields are applied to bulk semicon-

layers.[8] The shifts can exceed the exciton binding energy and yet the exciton resonances remain well resolved. Extended room-temperature measurements reported in this paper confirm the existence of exciton resonances up to $\sim 50E_i$ ($\sim 10^5$ V/cm). The purpose of this paper is to explain (i) the large shifts and (ii) the persistence of the exciton peaks to these large fields.

In contrast with the Franz-Keldysh effect which is independent of crystal size, the mechanism which we propose here requires the quantum confinement in the thin semiconductor layers. Large effects are to be expected with moderate fields because the particle-in-a-box envelope functions of electrons and holes and the exciton envelope functions are

For the sake of simplification, we shall take ΔE_c as infinite. (This is a close approximation for barriers $\gtrsim 100$ Å and $x \gtrsim 0.3$.) The wavefunction $u(\mathbf{r})$ of the electrons in this well obeys the time independent Schrödinger equation [2].

$$V(z)u(\mathbf{r}) - \frac{\hbar^2}{2m_c}\left(\frac{\partial^2}{\partial z^2} + \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2}\right)u(\mathbf{r}) = Eu(\mathbf{r}) \qquad (16.1\text{-}1)$$

E is the energy of the electron while $V(z) = E_c(z)$ is the potential energy function confining the electrons in the z direction. We will measure the energy relative to that of an electron at the bottom of the conduction band in the GaAs active region as shown in Figure 16-1. The eigenfunction $u(\mathbf{r})$ can be separated into a product

$$u(\mathbf{r}) = \psi_k(\mathbf{r}_\perp)u(z) \qquad (16.1\text{-}2)$$

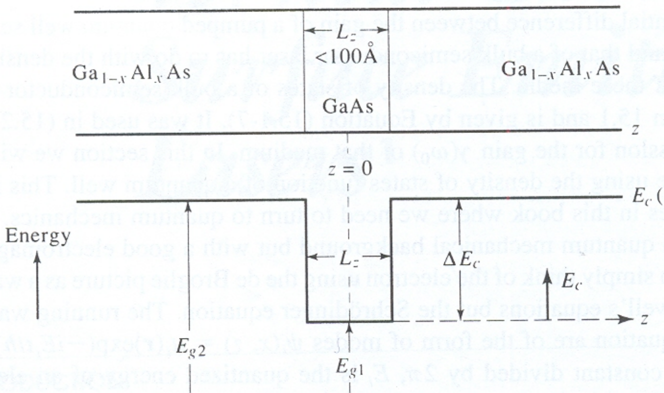
which, when substituted in (16.1-1), leads to

$$\left[V(z) - \frac{\hbar^2}{2m_c}\frac{\partial^2}{\partial z^2}\right]u(z) = E_z u(z) \qquad (16.1\text{-}3)$$

where E_z is a separation constant to be determined. Since we agreed to take the height of $V(z)$ the well region as infinite, $u(z)$ must vanish at $z = \pm L_z/2$.

$$u_\ell(z) = \begin{cases} \cos \ell \dfrac{\pi}{L_z} z & \ell = 1, 3, 5, \ldots \\[2ex] \sin \ell \dfrac{\pi}{L_z} z & \ell = 2, 4, 6, \ldots \end{cases} \qquad (16.1\text{-}4)$$

$$E_z = \ell^2 \frac{\hbar^2 \pi^2}{2m_c L_z^2} = \ell^2 E_{1c} \equiv E_{\ell c} \qquad \ell = 1, 2, 3, \ldots \qquad (16.1\text{-}5)$$



Using (16.1-3, 16.1-4, and 16.1-5) in (16.1-1) leads to

$$H(\mathbf{r}_\perp)\Psi(\mathbf{r}_\perp) = (E - E_z)\Psi(\mathbf{r}_\perp) \tag{16.1-6}$$

We can take $\psi(\mathbf{r}_\perp)$ as a two-dimensional Bloch wavefunction (see 15.1-1).

$$\psi(\mathbf{r}_\perp) = u_{\mathbf{k}_\perp}(\mathbf{r}_\perp)e^{(i\mathbf{k}_\perp \cdot \mathbf{r}_\perp)} \tag{16.1-7}$$

where $u_{\mathbf{k}_\perp}(\mathbf{r}_\perp)$ possesses the crystal periodicity. The wavefunction $\Psi(\mathbf{r}_\perp)$ obeys the Schrödinger equation

$$H(\mathbf{r}_\perp)\Psi(\mathbf{r}_\perp) = \frac{\hbar^2 k_\perp^2}{2m_c}\Psi(\mathbf{r}_\perp) \tag{16.1-8}$$

and from Equations (16.1-6, 7, 8)

$$E_c(\mathbf{k}_\perp, \ell) = \frac{\hbar^2 k_\perp^2}{2m_c} + \ell^2 \frac{\hbar^2 \pi^2}{2m_c L_z^2} = \frac{\hbar^2 k_\perp^2}{2m_c} + E_{\ell c} \qquad \ell = 1, 2, 3, \ldots \tag{16.1-9}$$

where the zero energy is taken as the bottom of the conduction band.

$u_{\mathbf{k}_\perp}(\mathbf{r}_\perp)$ possesses the lattice (two-dimensional) periodicity. Similar results with $m_c \to m_v$ apply to the holes in the valence band. We recall that the hole energy E_v is measured downward in our electronic energy diagrams so that

$$E_v(\mathbf{k}_\perp, l) = \frac{\hbar^2 k_\perp^2}{2m_v} + l^2 \frac{\hbar^2 \pi^2}{2m_v L_z^2} = \frac{\hbar^2 k_\perp^2}{2m_v} + E_{lv} \qquad l = 1, 2, 3, \ldots \tag{16.1-10}$$

measured (downward) from the top of the valence band. The complete wavefunctions are then

$$\psi_c(\mathbf{r}) = \sqrt{\frac{2}{L_z}}\Psi_{\mathbf{k}_{\perp c}}(\mathbf{r}_\perp)CS\left(\ell \frac{\pi}{L_z} z\right) \tag{16.1-11}$$

for electrons and

$$\psi_v(\mathbf{r}) = \sqrt{\frac{2}{L_z}}\Psi_{\mathbf{k}_{\perp v}}(\mathbf{r}_\perp)CS\left(\ell \frac{\pi}{L_z} z\right)$$

for holes. We defined $CS(x) \equiv \cos(x)$ or $\sin(x)$ in accordance with (16.1-4).

The lowest-lying electron and hole wavefunctions are

$$\psi_c(\mathbf{r})_{\text{ground state}} = \sqrt{\frac{2}{L_z}}\,\psi_{\mathbf{k}_{\perp c}}(\mathbf{r}_\perp)\cos\left(\frac{\pi}{L_z} z\right)$$

$$\psi_v(\mathbf{r})_{\text{ground state}} = \sqrt{\frac{2}{L_z}}\,\psi_{\mathbf{k}_{\perp v}}(\mathbf{r}_\perp)\cos\left(\frac{\pi}{L_z} z\right) \tag{16.1-12}$$

and are shown along with the next higher level in Figure 16-2. In a real semiconducting quantum well, the height ΔE_c of the confining well (see Figure 15-10) is finite, which causes the number of confined states, i.e., states with exponential decay in the z direction outside the well to be finite. The mathematical procedure for solving

$$N_1(m^{-3}) \longrightarrow \frac{\rho_{\mathrm{QW}}(\mathbf{k})}{L_z}\,dk f_v(E_v)[1 - f_c(E_c)] = \frac{k}{\pi L_z}(f_v - f_v f_c)dk$$

$$N_2(m^{-3}) \longrightarrow \frac{\rho_{\mathrm{QW}}(\mathbf{k})}{L_z}\,dk f_c(E_c)[1 - f_v(E_v)] = \frac{k}{\pi L_z}(f_c - f_c f_v)dk$$

where $\rho_{\mathrm{QW}}(\mathbf{k})$ is given by (16.1-15) and is independent of $\mathbf{k}$. The effective inversion population density due to carriers between k and $k + dk$ is thus

$$N_2 - N_1 \longrightarrow \frac{kdk}{\pi L_z}\left[f_c(E_c) - f_v(E_v)\right] \tag{16.2-3}$$

The division of ρ_{QW} by L_z is due to the need, in deriving the gain constant to use the *volumetric* density of inverted population consistent with the definition of N_1 and N_2 in (15.2-4). E_c and E_v are, respectively, the upper and lower energies of the carriers involved in a transition. We use (15.2-4) and (16.2-3) to write the contribution to the gain due to electrons within dk and in a single, say $\ell = 1$, sub-band as

$$d\gamma(\omega_0) = \frac{kdk}{\pi L_z}\left[f_c(E_c) - f_v(E_v)\right]\frac{\lambda_0^2}{4n^2\tau}\frac{T_2}{\pi[1 + (\omega - \omega_0)^2T_2^2]} \tag{16.2-4}$$

where T_2 is the coherence collision time of the electrons and τ is the electron-hole recombination lifetime assumed to be a constant. We find it more convenient to transform from the k variable to the transition frequency ω (see Equation 16.2-1). From (16.2-1) it follows that

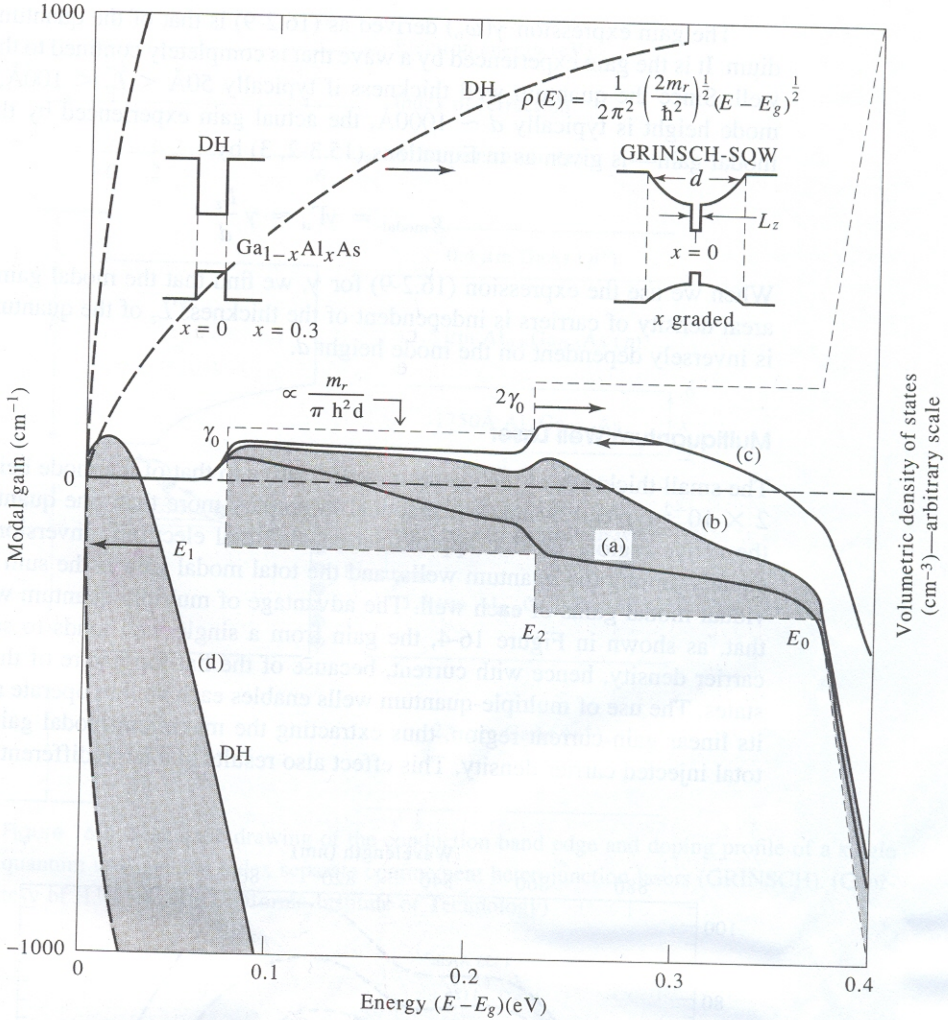
$$dk = \frac{m_r^*}{\hbar k}\,d\omega$$

so that (16.2-4) becomes

$$\gamma(\omega_0)_{n=1} = \frac{m_r^*\lambda_0^2}{4\pi\hbar L_z n^2\tau}\int_0^\infty [f_c(\hbar\omega) - f_v(\hbar\omega)]\frac{T_2 d\omega}{\pi[1 + (\omega - \omega_0)^2T_2^2]} \tag{16.2-5}$$

where we used the convention that $f_c(\hbar\omega)$ is the Fermi function at the upper transition (electron) energy E_c, while $f_v(\hbar\omega)$ is the valence band Fermi function at the lower transition energy. To include, as we should, the contributions from all other sub-bands ($\ell = 2, 3, \ldots$) we replace, using (16.1-16)

$$\frac{m_r^*}{\pi\hbar^2} \longrightarrow \frac{m_r^*}{\pi\hbar^2}\sum_{\ell=1}^{\infty} H(\omega - \omega_\ell) \tag{16.2-6}$$


1000
0
−1000
Modal gain (cm^{-1})
DH
$\rho(E) = \frac{1}{2\pi^2}\left(\frac{2m_r}{\hbar^2}\right)^{\frac{1}{2}}(E - E_g)^{\frac{1}{2}}$
DH
GRINSCH-SQW
d
L_z
x = 0
x graded
$Ga_{1-x}Al_xAs$
x = 0
x = 0.3
$\propto \frac{m_r}{\pi \hbar^2 d}$
$2\gamma_0$
γ_0
(c)
(b)
(a)
E_1
E_2
E_0
(d)
DH
Volumetric density of states (cm^{-3})—arbitrary scale
0
0.1
0.2
0.3
0.4
Energy $(E - E_g)$(eV)

Multiquantum Well Laser

The small thickness of the quantum well relative to that of the mode height ($L_z/d \approx 2 \times 10^{-2}$ typically) makes it practical to employ more than one quantum well as the active region. To first approximation, the total electronic inversion is divided equally among the quantum wells, and the total modal gain is the sum of the individual modal gains of each well. The advantage of multiple-quantum well lasers is that, as shown in Figure 16-4, the gain from a single well tends to saturate with carrier density, hence with current, because of the flat-top nature of the density of states. The use of multiple-quantum wells enables each well to operate much within its linear gain-current region, thus extracting the maximum modal gain at a given total injected carrier density. This effect also results in a large differential gain $A \equiv$

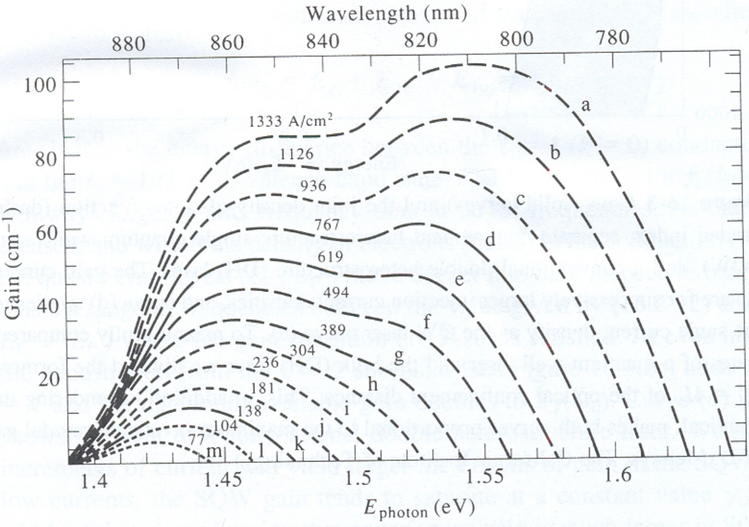


Figure 16-4 A theoretical plot of the exponential (modal) gain constant vs. wavelength of a quantum well laser. (Courtesy of Michael Mittelstein, The California Institute of Technology)

Phys. T. 10/92

ESAKI LEAVES IBM TO BECOME PRESIDENT OF JAPAN'S TSUKUBA UNIVERSITY

When Leo Esaki, the creator of the tunnel (or Esaki) diode, was elected president of the University of Tsukuba, it made front-page news in the Japanese press—not too surprising, given that he is the country's only living physics Nobel laureate. Besides, the circumstances of Esaki's election were most unusual: He is the first president of any of Japan's 97 national universities to come from outside academia, and, what's more, he has spent most of the past 32 years working and living outside Japan. "In physics terms," Esaki says, "it was almost a forbidden transition."

Esaki assumed his new position in April after retiring from IBM, with whom he was a fellow at the T. J. Watson Research Center in Yorktown Heights, New York. A native of Japan, he earned a BS from the University of Tokyo in 1947, and then, while working for Sony Corp, he received a PhD in physics from Tokyo in 1959. Shortly afterward he moved to the US to join IBM, where his research included work on semiconductor superlattices. (The article by Leroy Chang and Esaki on page 36 describes this work.)

Leo Esaki, right, listens to a talk given at his retirement symposium held near IBM's Watson Research Center on 1 May. The event was sponsored by IBM, Esaki's employer for 32 years before he became president of Tsukuba University in April. Seated next to Esaki is C. N. Yang.

Instant status

United States Patent

[11] 3,626,257

[72] Inventors Leo Esaki
Chappaqua;
Rudolf Ludeke, Katonah; Raphael Tsu,
Yorktown Heights, all of N.Y.

IBM

[21] Appl. No. 811,871
[22] Filed Apr. 1, 1969
[45] Patented Dec. 7, 1971
[73] Assignee International Business Machines
Corporation
Armonk, N.Y.

[54] SEMICONDUCTOR DEVICE WITH SUPERLATTICE REGION

22 Claims, 10 Drawing Figs.

[52] U.S. Cl. .... 317/234 R,
317/235 K, 317/235 AL, 317/235 AK, 317/235 AD
[51] Int. Cl. H011 5/00
[50] Field of Search 317/234
(10), 235 (25), 235 (42), 235 (43), 235; 331/107
G, 115; 307/284

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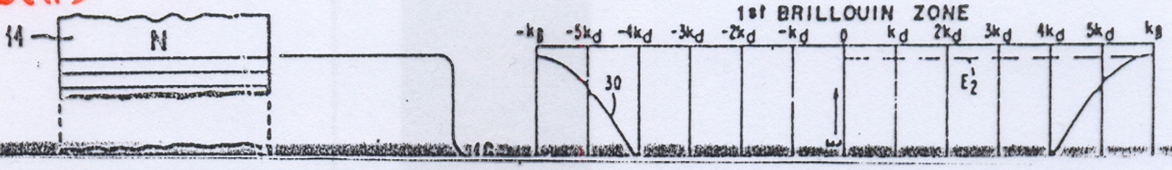
Primary Examiner—John W. Huckert
Assistant Examiner—Martin H. Edlow
Attorneys—Hanifin and Jancin and John E. Dougherty, Jr.

ABSTRACT: The semiconductor device has two highly N-type end portions to which ohmic contacts are made, and a central portion which has a one dimensional spatial periodic variation, in its band-edge energy. This spatial periodic variation, or superlattice, is produced by doping or alloying to form a plurality of successive layers having alternating band-edge energies. The period of the spatial variation is less than the carrier mean free path, and is such as to form in momentum space a plurality of periodic mini-zones which are much smaller than the Brillouin zones. The device exhibits a bulk negative resistance and is used in oscillator and bistable circuits.

[e] 공명터널

Dingle .. (AT&T Bell) ~'74
[o]
엑서톤



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Miller

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SEED

[54] NONLINEAR AND BISTABLE OPTICAL DEVICE

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372/12; 372/45; 372/46
[58] Field of Search .... 250/211 J; 340/870.29;
377/102; 372/12, 45, 46; 350/354

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[57] ABSTRACT

The invention is a nonlinear or bistable optical device having a low switching energy. The invention uses a means responsive to light for generating a photocurrent, a structure having a semiconductor quantum well region, and means responsive to the photocurrent for electrically controlling an optical absorption of the semiconductor quantum well region. The optical absorption of the semiconductor quantum well region varies in response to variations in the photocurrent. A photodiode or phototransistor may be used as the means responsive to light, and may be made integral with the structure having the semiconductor quantum well region. An array of devices may be fabricated on a single chip for parallel logic processing.

26 Claims, 27 Drawing Figures

of the degenerate $|2p\rangle$ states:

$$
\begin{aligned}
|x\rangle &\equiv \frac{1}{\sqrt{2}}(u_{211} + u_{21-1}) = \frac{1}{\sqrt{32\pi}}\left(\frac{Z}{a_0}\right)^{5/2} e^{-Zr/2a_0}\, x \\
|y\rangle &\equiv \frac{1}{\sqrt{2}}(u_{211} - u_{21-1}) = \frac{i}{\sqrt{32\pi}}\left(\frac{Z}{a_0}\right)^{5/2} e^{-Zr/2a_0}\, y \qquad (8.3.11) \\
|z\rangle &= u_{210} = \frac{1}{\sqrt{32\pi}}\left(\frac{Z}{a_0}\right)^{5/2} e^{-Zr/2a_0}\, z
\end{aligned}
$$

Referring to the ground state u_{100} as $|1\rangle$, we have

$$\langle 1|x|x\rangle = \frac{1}{\sqrt{32\pi^2}} \frac{1}{(a_0')^4} \int_0^\infty r^4 e^{-3r/2a_0'}\, dr \int_0^\pi \sin^3\theta\, d\theta \int_0^{2\pi} \cos^2\Phi\, d\Phi$$

where $a_0' \equiv a_0/Z$. Using

$$\int_0^\infty r^4 e^{-3r/2a_0'}\, dr = \frac{\Gamma(5)}{(3/2a_0')^5}$$

leads to

$$\langle 1|x|x\rangle = \langle 1|y|y\rangle = \langle 1|z|z\rangle = 0.7449\, \frac{a_0}{Z}$$

All the other matrix elements connecting the upper and lower states are zero as determined by symmetry consideration, so we can write

$$
\begin{aligned}
|x_{12}|^2 &= |\langle 1|x|x\rangle|^2 = 1.5539 \times 10^{-21}/Z^2 \\
|y_{12}|^2 &= |\langle 1|y|y\rangle|^2 = 1.5539 \times 10^{-21}/Z^2 \\
|z_{12}|^2 &= |\langle 1|z|z\rangle|^2 = 1.5539 \times 10^{-21}/Z^2
\end{aligned}
$$

To finish the calculation, we recall that for a hydrogenic transition

$$\hbar\omega_0(n = 2 \rightarrow n = 1) = \left(\frac{3}{4}\right) \frac{\mu Z^2 e^4}{32\pi^2 \varepsilon_0^2 \hbar^2}$$

$$a_0 = \frac{4\pi\varepsilon_0\hbar^2}{\mu e^2}$$

$$\mu^{-1} = m_e^{-1} + m_n^{-1}$$

which when used in (8.3-9) gives

$$t_{\text{spont}|2p\rangle\rightarrow|1s\rangle} = \frac{1}{W_{\text{spont}}} = \frac{1.595 \times 10^{-9}}{Z^4}\ \text{sec.} \qquad (8.3\text{-}12)$$

The interaction Hamiltonian is (see Appendix 5)

$$\mathcal{H}' = -e\mathbf{E}_l(\mathbf{r}, t) \cdot \mathbf{r} \tag{8.3-1}$$

Let the mode l correspond to a plane-wave propagating along $\mathbf{k}$ with a polarization λ. If we use (5.6-15), the interaction Hamiltonian is

$$\mathcal{H}' = ie\sqrt{\frac{\hbar\omega_{\mathbf{k}}}{2V\varepsilon}}\,[a^{\dagger}_{\mathbf{k},\lambda}e^{-i\mathbf{k}\cdot\mathbf{r}} - a_{\mathbf{k},\lambda}e^{i\mathbf{k}\cdot\mathbf{r}}]\hat{\mathbf{e}}_{\mathbf{k},\lambda} \cdot \mathbf{r} \tag{8.3-2}$$

Consider the transition depicted in Figure 8.3. The initial state is $|2, n_k\rangle$ in which the atom is in level 2 and the mode $(\mathbf{k}, \lambda)$ has $n_{\mathbf{k}}$ quanta. The final state finds the atom in level 1, whereas the mode has gained a quantum and is in the state $(n_{\mathbf{k}} + 1)$. The transition rate of the system from the initial state $|2, n_{\mathbf{k}}\rangle$ to the final $|1, n_{\mathbf{k}} + 1\rangle$ state is obtained from (3.12-14) as

$$W = \frac{2\pi}{\hbar}|\mathcal{H}'_{fi}|^2\delta(E_{\text{initial}} - E_{\text{final}}) \tag{8.3-3}$$

$$= \frac{2\pi e^2}{\hbar}\left(\frac{\hbar\omega_{\mathbf{k}}}{2V\varepsilon}\right)|\langle 1, n_{\mathbf{k}} + 1|a^{\dagger}_{\mathbf{k},\lambda}\hat{\mathbf{e}}_{\mathbf{k},\lambda} \cdot \mathbf{r}|2, n_{\mathbf{k}}\rangle|^2\,\delta(E_2 - E_1 - \hbar\omega_{\mathbf{k}})$$

$$= \frac{\pi e^2\omega_{\mathbf{k}}}{V\varepsilon}|\langle 1|\hat{\mathbf{e}}_{\mathbf{k},\lambda} \cdot \mathbf{r}|2\rangle|^2(n_{\mathbf{k}} + 1)\,\delta(E_2 - E_1 - \hbar\omega_{\mathbf{k}}) \tag{8.3-4}$$

where we used

$$\langle n_{\mathbf{k}} + 1|a^{\dagger}_{\mathbf{k},\lambda}|n_{\mathbf{k}}\rangle = \sqrt{n_{\mathbf{k}} + 1}$$

Plane-Wave Quantization

The discussion just concluded uses a generalized resonator of unspecified shape. We will find it useful to consider the form of the field operators in the case of a plane-wave resonator. Although such a resonator, which requires an infinite cross-sectional area does not exist, most optical resonators that employ curved mirrors as reflectors involve nearly plane-wave propagation.

To be specific, consider the *l*th mode of a resonator of length L along the axis and mode volume V. Let the electric and magnetic field vectors point along the y and x directions, respectively. Equations (5.5-3, 4, 9) are satisfied by

$$
\begin{aligned}
\mathbf{E}_l(\mathbf{r}, t) &= -i\hat{\mathbf{x}}\left(\frac{\hbar\omega_l}{V\varepsilon}\right)^{1/2}[a_l^\dagger(t) - a_l(t)]\sin k_l z \\
\mathbf{H}_l(\mathbf{r}, t) &= \hat{\mathbf{y}}\left(\frac{\hbar\omega_l}{V\mu}\right)^{1/2}[a_l^\dagger(t) + a_l(t)]\cos k_l z
\end{aligned}
\tag{5.6-14}
$$

where V is the mode volume. In the case of plane wave modes, the total field can be written as a summation over all modes $\mathbf{k}$

$$
\begin{aligned}
\mathbf{E} &= \sum_{\mathbf{k},\lambda} - i\,\hat{\mathbf{e}}_{\mathbf{k},\lambda}\sqrt{\frac{\hbar\omega_k}{2\varepsilon V}}\left(a_{\mathbf{k},\lambda}^\dagger\, e^{-i\mathbf{k}\cdot\mathbf{r}} - a_{\mathbf{k},\lambda}e^{i\mathbf{k}\cdot\mathbf{r}}\right) \\
\mathbf{H} &= \sum_{\mathbf{k},\lambda} - i\,\frac{\mathbf{k}}{k}\times\hat{\mathbf{e}}_{\mathbf{k},\lambda}\sqrt{\frac{\hbar\omega_k}{2\mu V}}\left(a_{\mathbf{k},\lambda}^\dagger\, e^{-i\mathbf{k}\cdot\mathbf{r}} + a_{\mathbf{k},\lambda}e^{i\mathbf{k}\cdot\mathbf{r}}\right)
\end{aligned}
\tag{5.6-15}
$$