Dependence of Electron Temperature on Well Width in the $Al_{0.48}In_{0.52}As/Ga_{0.47}In_{0.53}As$ Single-Quantum Well

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Abstract—**The temperature of electrons confined in a single quantum well (SQW) as a function of well width (L_z) from L_z = 14.5 Å to L_z = 130.5 Å have been determined from photoluminescence spectra of a set of undoped (Al,In)As/(Ga,In)As SQW's. From steady state electron temperature variation with L_z it is shown that the electron and longitudinal optical phonon scattering dominates the excess energy loss of the thermalized two dimensional electrons in the SQW's and the scattering rate is independent of L_z within our experimental accuracy. The average energy loss rate per hot-electron was determined to be much smaller than expected.**

RECENTLY, there has been considerable interest in the carrier transport and optical properties of electrons confined in quantum wells. Knowledge of the hot-electron behavior in semiconductor microstructures is needed to develop small and fast optoelectronic devices. It is crucial to understand the relationship between electron temperature and electron–phonon (e-p) interaction in a single quantum well (SQW) with various well widths (L_z) and to determine how they differ from their bulk counterpart. Theoretical studies have been performed related to e-p interaction in two-dimensional (2-D) structures. Hess and his coworkers [1] predicted an enhanced e-p scattering rate over the bulk semiconductor by a factor of $f = 3.5(E_{LO}/E)^{1/2}$ for $E_{LO} < E < 10E_{LO}$. Ridley [2] obtained a $1/L_z$ relationship of 2-D e-p scattering rate by introducing a momentum conservation approximation in well direction (MCA). Leburton [3] has shown that L_z effect has a rather weaker influence on longitudinal optical (LO) phonon and electron interaction than Ridley's result. Although theoretical calculations predict a somewhat enhanced e-p interaction in quasi-2-D structures, the experimental results by many groups [4]–[6] have shown a much slower hot-electron cooling rate. The explanation for this contradiction is not yet clear. This paper reports on measuring the variation of temperature (T_c) of the thermalized hot electrons in the SWQ's at various well widths. The carrier temperatures are described by a model taking into account the high carrier density and slow carrier relaxation.

The T_c was determined using the method introduced by Shah [7], [8], namely, by fitting the high-energy tails of photoluminescence (*PL*) spectra of electron system in a set of undoped (Al,In)As/(Ga,In)As SQW's for the different well width from L_z = 14.5 Å to L_z = 130.5 Å. The observed variation of steady state electron temperature with L_z provides direct information about the quasi-2-D electron and LO-phonon scattering rate which we found to be independent of L_z within our experimental accuracy. The average energy loss rate per hot-electron was determined to be much smaller than 3-D electron gas [8]. The intrinsic 2-D electron degeneracy may cause further lowering of e-p interaction than the 3-D electron case when carrier density is high enough such that $E_{fe} - E_{cn} = k_B T_c ln(n/N_c) > k_B T_c$ where E_{fe} is the quasi-Fermi level for electrons and E_{cn} is the energy level of the subconduction band edge. However, Shah *et al.* [6] have recently attributed this reduction in electron-energy-loss rate primarily to hot-phonon effect while hole-energy-loss rate is not significantly altered.

The epitaxial layers used in this study were grown by a Varian MBE machine [9]. The structures were grown on a thermally cleaned InP substrate consisted of a 2500 Å thick AlInAs layer followed by the GaInAs SQW's and 100 Å layer of AlInAs. The thickness of the GaInAs SQW was varied from 14.5 Å to 130.5 Å. Both of GaInAs and AlInAs were not intentionally doped. Each SQW was pumped at a wavelength of 5145 Å ($h\omega_e$ = 2.41 eV) with a low-power density excitation fixed at about 10 W/cm^2. The sample was suspended in a gas transfer liquid helium cryostat which kept the sample temperature (T_L) at 3.5 K. The *PL* spectra were measured by a double dispersive spectrometer. The signal was detected by Ge and S-1 photodetectors.

Four typical *PL* spectra from well GaInAs for different well size are displayed in Fig. 1 as circles. The emission peak energy increases with decrease of well width L_z due to carrier confinement and the emission extends to much higher energies than their bulk counterpart. More important feature is the extent of high-energy tail increases as

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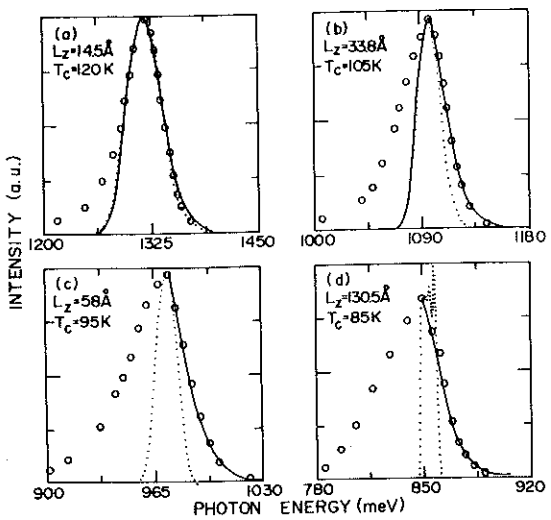

Fig. 1. The PL spectra (circles) from the SQW's of various well-widths with excitation power density about 10 W/cm^2 at lattice temperature T_c = 3.5 K. The dotted lines are the theoretical fits using the *maximum extent* of well-width fluctuations (σ = 1.2 Å) with assumption that carrier temperatures are same as lattice temperature. The solid lines in (a) and (b) are the exact fits using convolution expression (10) with σ = 1 Å and various electron temperatures displayed in the figure. The solid lines in (c) and (d) are the high-energy-tail fits using (7). The edges of dotted lines in low energy sides marks the boundary between extrinsic and intrinsic emissions.

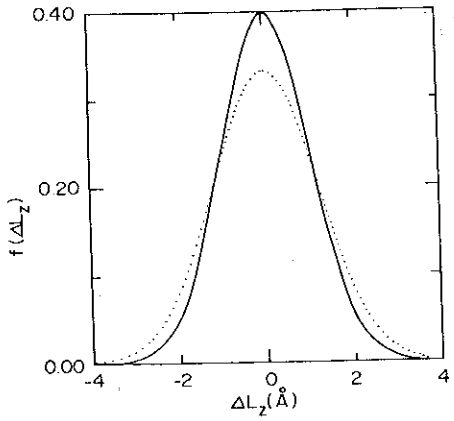

Fig. 2. Probability density functions versus well-width change ΔL_z with σ = 1 Å for solid curve and σ = 1.2 Å for dotted curve, respectively.

well width decreases. For example, at L_z = 14.5 Å and 130.5 Å, the high-energy tails above half-maximum of luminescence extend to 39 meV and 25 meV, respectively. This implies that the fractional number of energetic electrons in the n = 1 subband of the SQW is increased as well width decreases.

The question should be addressed here is whether or not this broad high-energy tail could be explained only by well width fluctuations. The mechanisms of the broading the full line-width at half-maximum (FWHM) of the luminescence from SQW's have been discussed by Welch *et al.* [10]. It was found that the dominant *PL* spectra broading was primarily due to a transfer of electrons from AlInAs into GaInAs. The broad FWHM of the spectra could not interpreted only by well width fluctuations even under the worst case of well width fluctuations. For the present investigation, we focus on the spectral characteristics of high-energy tail where the emission intensity decreases exponentially with photon energy to explain the broadening of the high-energy tails. It is well known that inhomogeneous broadening resulting from well-width fluctuations within a well can be quite large in thin well and the fact that the high energy tail decreases exponentially is not sufficient evidence for assignment of a hot electron temperature which differs from lattice temperature. We must deconvolve the broadening out of the high energy tail before deducing electron temperature from it. The strategy for this purpose is illustrated in following section.

For an ultrathin SQW, the FWHM of its *PL* spectrum is dominated by well-width fluctuations. We can fit (see next paragraph for discussion on fitting) the line-shape of the luminescence spectrum of the 14.5 Å well to obtain a *maximum extent* for well-width fluctuations as if the electron system were in equilibrium with the lattice. Since the interfacial quality of the structures are independent of well-width, we can calculate the line-shapes of *PL* spectra for the thicker wells using the *maximum extent* of well-width fluctuations with assumption that the electron temperatures are same as the lattice temperature. The calculated results are shown in Fig. 1(a)–(d) as dotted curves for various well-widths. It is clearly demonstrated that the change in the high-energy tail with change in the well-width is only a partial result of the increased inhomogeneous broadening as well-width decreases. Therefore, an illusive explanation of high-energy tail broadening entirely by well-width fluctuation is ruled out. Introducing a hot electron temperature to explain the broadening of high-energy tail is numerically justified.

The theoretical calculation of *PL* spectra is outlined in this section. The probability density function for a random variable (r.v.) L_z in "large-island approximation" [11] is assumed to have a Gaussian shape given by

$$f(L_z) = \frac{1}{\sqrt{2\pi}\,\sigma} \exp\left[-(L_z - L_{zo})^2/2\sigma^2\right] \qquad (1)$$

where L_{zo} and σ^2 are the mean well-width and the variance of r.v. L_z, respectively. The statistics of r.v. L_z is then characterized by its variance σ^2 and independent of mean L_{zo}. Two Gaussian shapes are depicted in Fig. 2. The dotted curve with obtained *maximum extent* σ = 1.2 Å gives a value of 2.8 Å for FWHM of the curve, while solid curve corresponds σ = 1 Å having FWHM = 2.4 Å. The electrons and holes within the well will experience the well-width fluctuations resulting in fluctuation of their confinement energies. The effective energy gap is defined as

$$E_{ge} = E_{1e} + E_{1h} + E_g \qquad (2)$$

where, E_{1e} and E_{1h} denote the lowest energy edges of electron and hole subbands, and E_g is the band gap of well material (GaInAs). For a finite SQW, E_{1e} can be expressed by

$$E_{1e} = E_o(L_{zo}/L_z)^{\alpha} \qquad (3)$$

with

$$E_o = \frac{h^2}{8m^* L_{zo}^{\alpha}} \tag{4}$$

where m^* is the electron effective mass and α can be determined from an exact equation given in [12]. Using a standard method [13] of transformation of a r.v., the probability density function of r.v. E_{ge} with neglect of E_{1h} fluctuation and excitonic effect [14] were found to be [15]

$$f(E_{ge}) = g(E_{ge}) \exp\left(-(\eta/2)\left[(E_o/(E_{ge} - E_g - E_{1h}))^{1/\alpha} - 1\right]^2\right) \tag{5}$$

with

$$g(E_{ge}) = \sqrt{(\eta/2\pi)}\, \frac{E_o^{\alpha}}{\alpha (E_{ge} - E_g - E_{1h})^{(1+\alpha)/\alpha}} \tag{6}$$

where $\eta = (L_{zo}/\sigma)^2$. This expression gives *PL* spectrum as if all the carriers were located at random moving subband edges.

Under condition that well-width is certain the intrinsic *PL* spectra arise from recombination of thermalized electrons and holes within subbands were calculated using allowed density of states (ρ) and Fermi–Dirac distributions (f_e and f_h). This is given by

$$PL^{(\text{intri})} \sim \rho_e f_e \rho_h f_h \tag{7}$$

with

$$\rho_i = \frac{m_i^*}{\pi h^2} \tag{8}$$

and

$$f_i = \frac{1}{1 + \exp\left[(\epsilon_i - \epsilon_{fi})/kT_c\right]} \tag{9}$$

where ϵ_i and ϵ_{fi} are the kinetic energy and quasi-Fermi energy measured from the subband edge, $i = e, h$ denote the electron and the hole. In our calculation, momentum conservation and energy dependent electron effective mass in the manner given in [16] were considered.

The exact *PL* spectra can be calculated by taking into account both well-width fluctuations and hot electron temperatures using the following convolution expression.

$$PL^{(\text{exact})}(E) = \int_{-\infty}^{E} PL^{(\text{intri})}(E - E_{ge})\, f(E_{ge})\, dE_{ge}. \tag{10}$$

The calculated results for $L_z = 14.5$ Å and 33.8 Å using $\sigma = 1$ Å are shown in Fig. 1(a) and (b) as solid curves. The values of 120 K and 105 K for electron temperatures for the two wells were obtained.

The main conclusions from such a fitting process are summerized as follows:

i) Using $\sigma = 1$ Å the high-energy tails of *PL* spectra from various well-widths can be fitted with various electron temperatures which differs from the lattice temperature T_L.

ii) For $L_z > 40$ Å, the high-energy tail can be fitted only using (7). The fitting using (10) just adds some fine structures [17] on high-energy tails. But the envelope of high-energy tails remain unchanged. Two solid curves for 58 Å well and 130.5 Å-well were calculated using (7). The electron temperatures were determined 95 K and 85 K, respectively.

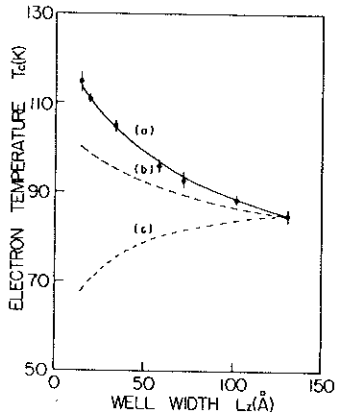

Fig. 3. Electron temperature (T_c) is plotted versus well width L_z at fixed power density excitation of about 10 W/cm^2. The solid curve (a) is a plot of (14). The dash curves (c) and (b) are the plots of (16) (Ridley) and (17) (Leburton), respectively.

iii) It was found that $kT_c/\epsilon_{fe} \approx 0.69$ and $\epsilon_{fh} \sim 5$ meV lead to the best fits of the measured spectra. The values of ϵ_{fe} deduced from spectra are related to the excess carrier density. For example, $\epsilon_{fe} = 10.6$ meV for 130.5 Å well corresponds to a sheet carrier density n_s of 2.2×10^{11} cm^{-2}.

iv) Another observation from Fig. 1 is that the widths of the low energy side of emission peaks are about same ($\sim$60 meV) for all the wells investigated. This confirms that those emissions are extrinsic. They either arise from phonon-assisted recombinations or deep impurities. In general, it is very complicated to analyze. Therefore, no attempts were made to fit those emissions in this paper.

The experimental data of T_c versus L_z extracted from *PL* spectra of SQW's are plotted by the dots in Fig. 3. There are two salient features. First, electrons in a deep SQW are much hotter than their bulk counterpart under the same excitation conditions [18]. This hot feature is due to the high density in the well via the trapping of the carriers from the barriers. The relatively large amount of electrons have to be distributed in high-energy tail in order to form a steady state distribution because the allowed density of states is independent of energy within one subband for a 2-D electron system. Second, the electron temperature increases as well width becomes smaller. This dependence arises from the increase in power transfer from photogenerated energetic electrons to thermalized electrons [7], which is also due to the increasing thermalized electron density in the well as L_z decreases.

The steady state variation T_c with L_z shown in Fig. 3 discussed above is further quantitatively explained using the following model. Under steady state condition the electron temperature T_c is determined by the balance between the rate at which the electron gas receives energy from the photogenerated electron and the rate at which the electron gas loses energy to the lattice. The excitation

laser beam with photon energy $h\omega_e$ creates excess energetic electrons (N_{ex}) into the conduction band of the SQW. It is those electrons via drift and diffusion that are effectively trapped into the SQW and thermalized among themselves at T_c via electron–electron collision. It should be noted that a substantial electron population may exist at high T_L as a result of transfer from unintentionally doped AlInAs even without photopumping. However, at very low T_c, deep impurities cannot be totally ionized. This is convinced by the relative *PL* intensity variation with T_L between well and barriers [12].

The power per electron given to the thermalized electrons from photogenerated energetic electrons is given by [7]

$$P_e = P_a(E_k/h\omega_e)\, f(1/C_{tr}N_{ex}) \tag{11}$$

where f is the fraction of absorbed laser power (P_a) by the sample given to the thermalized electron system by the energetic electrons; $A \sim 4 \times 10^{-4}$ cm^2 is the laser beam area at the sample; E_k is the average excess energy per electron and the factor ($E_k/h\omega_e$) is about 0.3; and C_{tr} is a phenomenological factor that takes trapping process into account. The value of f is given by $f = n_w/(n_w + n_c^*)$ where $n_c^* \sim 5 \times 10^{17}$ cm^{-3} is the density of thermalized electrons at which the photogenerated energetic electrons lose the same amount of power as directly to the lattice by emission of the energetic *LO* phonons. The $n_w = C_{tr}N_{ex}/AL_z$ is the thermalized electron gas density in the well. The fractional number f can be rewritten as follows:

$$f = 1/(1 + L_z/4L_c) \tag{12}$$

by defining a characteristic length $L_c = C_{tr}N_{ex}/4An_c^*$. Since the well width L_z is much less than absorption length (L_a) which is about 1 μm for both GaInAs and AlInAs, P_a and E_k can be treated as constant. The factor C_{tr} is inversely proportional to the average trapping time τ_{tr} which may be expressed by [14] $\tau_{tr} \sim (L_a - L_z)/v_d \sim L_a/v_d$ (for $L_a >> L_z$) where v_d is an average drift velocity. The v_d may be very close to the maximum value of 10^7 cm/s (1 μm/10 ps) due to large well depth [12] for our samples. Therefore, the total number $C_{tr}N_{ex}$ of electrons that are transferred from the AlInAs can be also assumed as a constant [10]. The photogenerated carrier density is estimated to be about 1.1×10^{15} cm^{-3}. This yields a value of $L_c = 11$ Å. Therefore, from (12), the narrower well width, the more efficient is the energy transfer from energetic electrons to the thermalized electrons.

When the energy loss of thermalized electrons in the SQW is still dominated by the polar *LO* phonons, the average power per electron transferred from electrons in the SQW to the lattice is given by [7], [20]

$$P_L = E_{LO}/\tau_{e\text{-}p} \exp\,[-E_{LO}/k_B T_c(L_z)] \tag{13}$$

where $E_{LO} = 32$ meV is the *LO* phonon energy. The value of $E_{LO} \exp[-E_{LO}/k_B T_c(L_z)]$ is the average energy loss per hot-electron. The exponential factor is the fractional number of the thermalized electrons in the SQW that can emit *LO* phonons, and the $\tau_{e\text{-}p}$ is an approximate *e-p* scattering time (not averaged over a Fermi–Dirac distribution). Under steady state equilibrium condition, the electron temperature T_c can be determined from $P_e(L_z) = P_L(T_c)$, that is

$$T_c = E_{LO}/k_B \ln\,[g(1 + L_z/4L_c)] \tag{14}$$

where the factor g is given by

$$g = (C_{tr}N_{ex}/0.3P_a)(E_{LO}/\tau_{e\text{-}p}). \tag{15}$$

The value of $g = 20$ is obtained by the least square fit of the experimental data given in Fig. 3. The calculated variation of T_c with L_z given by (14) is plotted in Fig. 3 as a solid curve (a). It should be emphasized that the data were taken under the fixed excitation power since T_c depends on the excitation as shown in [12]. The excellent fit of (14) to the steady state electron temperature data extracted from high-energy tails of the PL spectra implies three important conclusions. First, the electron and polar *LO* phonon scattering dominates the excess energy loss of the thermalized electrons in the SQW's. Second, since the fractional number f given in (12) increases by the same amount as the exponential factor given in (13) does when L_z decreases from 130.5 Å to 14.5 Å, the e-p scattering rate ($R_{2D} = 1/\tau_{e\text{-}p}$) in SQW is independent of the well width L_z within our experimental accuracy, indicating that the influence of reduced dimensionality does not play an important role in carrier-energy-loss rates. This is in contrast to the theories of Ridley [2] and Leburton [3]. (See the next section for details.) Third, because the excess carriers trapped into the well are proportional to the power absorbed by the sample, the value of g given in (15) is inversely proportional to the time constant $\tau_{e\text{-}p}$ for the hot-electron relaxation through phonon emission. This enables us to estimate $\tau_{e\text{-}p}$. Since P_a is about 4 mW, a value for the e-p scattering time $\tau_{e\text{-}p}$ of about 18 ps is calculated from (15). This is comparable to the time determined by Ryan *et al.* [5]. The $\tau_{e\text{-}p}$ in the SQW's is about a factor of 4 longer than its bulk counterpart [8] within the carrier density range from $n_w \sim 1.7 \times 10^{17}$ cm^{-3} to 1.5×10^{18} cm^{-3}. This indicates the intrinsic 2-D electron degeneracy effect may be the main cause for this further slower relaxation than 3-D case since ϵ_{fe} is greater than $k_B T_c$ for the wells measured [21]. Shah *et al.* [6] and others [4], [5] have also observed further slower electron energy relaxation in (AlGa)As/GaAs system. Shah *et al.* [6] concluded that nonequilibrium optical phonon effects was the primary cause for the low electron-energy-loss rate. The hot-phonon effect may explain some experimental results. However, how to establish this nonequilibrium population when $\tau_{e\text{-}p}$ is comparable [5] or longer than phonon lifetime τ_p is still not clear.

From theoretical studies, [1]–[3] the e-p scattering rate R_{2D} should be dependent on L_z and somewhat enhanced as L_z decreases. This implies the electron temperature T_c would increases slowly or even decreases as L_z becomes smaller. When the approximate analytic results for the R_{2D} at the onset of phonon emission calculated by Ridley [2] and Leburton [3] are used in deriving (14), T_c are given

by the following.

$$T_c = E_{LO}/k_B \ln g[(1 + L_z/4L_c)130.5/L_z] \quad (16)$$

and

$$T_c = E_{LO}/k_B \ln g[(1 + L_z/4L_c)/1.63(1 + L_z/4L_o)] \quad (17)$$

respectively, where $L_o = h/(2m^*E_{LO})^{1/2}$. The calculated variation of T_c with L_z given in (16) (Ridley) and (17) (Leburton) are also plotted in Fig. 3 as dash curves (c) and (b). A forced fit at the data point (130.5 Å, 85 K) has taken into account in both (16) and (17). These curves do not fit the data well. However, Leburton's model shows a similar trend for the variation of T_c with L_z, indicating that the phonons that interact with confined electrons do not behave much differently from those of the bulk. The dash curve (c) clearly demonstrates that the MCA fails for the polar *LO* phonon intrasubband scattering in thin quantum wells. However, the discrepancy is partially due to the crude approximation used for calculating the phonon scattering rate in 2-D structures. More exact calculation is still needed in order to take the uncertainty relationship and the fuzzy momentum conservation in well direction [2] into account.

In conclusion, temperature of electrons confined in the SQW's has been determined from high-energy tails of *PL* spectra of a set of undoped SQW's. The high electron temperature in the SQW's is attributed to the high carrier density and slow hot-electron relaxation. The observed variation of the steady state electron temperature with L_z indicates that the phonon scattering rate for electrons confined in SQW's does not strongly depend on the well width L_z within our experimental accuracy.

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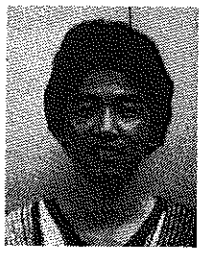

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P. P. Ho, for a photograph and biography, see p. 204 of the January 1986 issue of this JOURNAL.

R. R. Alfano, for a photograph and biography, see p. 204 of the January 1986 issue of this JOURNAL.

D. F. Welch, photograph and biography not available at the time of publication.

G. W. Wicks, photograph and biography not available at the time of publication.

L. F. Eastman (A'53–M'58–SM'65–F'69), photograph and biography not available at the time of publication.