Coherent Enhancement of the Hot-Electron Mean Free Path by Superlattice Transmission Resonances

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We report a strong enhancement of the optical responsivity in GaAs/AlGaAs multi-quantum-well infrared detectors. This is explained by a resonant increase in the hot-electron mean free path as a result of the coherent transmission resonances in the continuum states above the energy barriers of the superlattice.

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There has been a great deal of interest recently in studying the continuum resonances above the energy barriers in multi-quantum-well superlattices,[1-5] as well as hot-electron transport[6] and capture into quantum wells.[7-11] We discuss here the direct determination of the hot-electron mean free path L for an electron which is photoexcited out of a doped GaAs quantum well, transports above the superlattice energy barriers, and is subsequently recaptured into a distant well. Using a novel infrared photoconductivity technique, we have observed, for the first time, a striking enhancement in the low-temperature L (relative to the high-temperature L), caused by the coherent resonances in the continuum states.

The two samples used in the study were grown using molecular-beam epitaxy and had the following structures. Sample A had a 50-period superlattice consisting of 45-Å GaAs wells (doped $n=1\times10^{18}$ cm^{-3}) and 320 Å of undoped $Al_{0.28}Ga_{0.72}As$ barriers; these were sandwiched between GaAs contact layers ($n=1\times10^{18}$ cm^{-3}). Sample B was similar except the wells were 38 Å and the barriers were 328 Å of $Al_{0.36}Ga_{0.64}As$. These were processed into 200-μm-diam mesas and the substrate polished at 45° to allow (as required by the quantum-mechanical selection rules) a component of the optical electric field perpendicular to the superlattice.[12,13]

The measured[12] room-temperature-normalized absorption coefficient α for photoexcitation of an electron from the single bound state in the well to the lowest continuum state above the barrier,[4,14,15] is shown in Fig. 1 for both samples (dashed curve). For comparison we also show the photoconductive responsivity spectra (R) measured at $T=20$ K (using the apparatus described previously[13-15]). The shift of the responsivity curves to higher energy is a result of the lower temperature. For sample A the $T=20$ K responsivity spectrum is also narrower than the $T=300$ K absorption spectrum as a result of the lower temperature. The fact that this is not the case for crystal B is due to the excited state being higher in the continuum and hence broader in energy.

The low-temperature bias dependence of the respon-

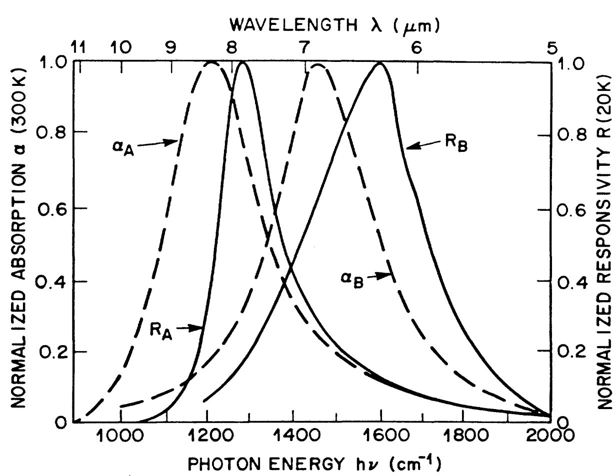

FIG. 1. Normalized absorption α and responsivity R spectra (measured at $T=300$ and 20 K, respectively) for two crystals A and B. The actual peak values are as follows: $\alpha_A=\alpha_B=2000$ cm^{-1}, $R_A=0.55$ A/W, and $R_B=0.44$ A/W.

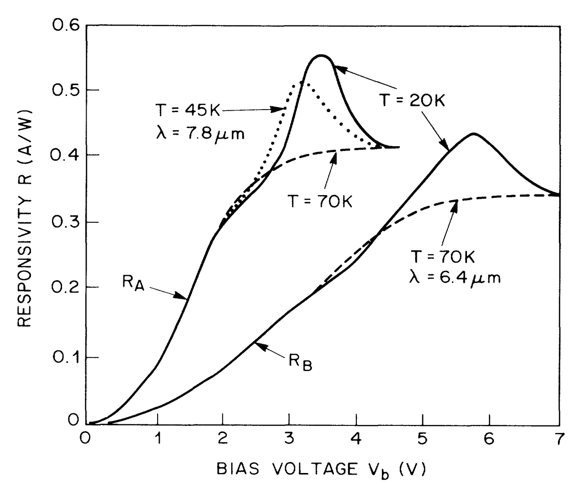

FIG. 2. Responsivity as a function of bias voltage for different temperatures. Sample A was measured at a wavelength $\lambda=7.8$ μm while B was measured at $\lambda=6.4$ μm.

sivities clearly shows a striking resonance enhancement (Fig. 2). These data were taken using infrared radiation having a wavelength near the peak of the responsivity spectra ($\lambda = 7.8$ and 6.4 μm for samples A and B, respectively). Note that as the temperature is increased the resonance enhancement decreases in amplitude and disappears by $T = 70$ K and that the $T = 20$ and 70 K curves cross. Thus, surprisingly there is a bias for which the high-temperature responsivity is actually larger than the low-temperature responsivity. We will show later that this large enhancement is due to a coherent resonance of the photoexcited electron transmission coefficient in the continuum state above the $Al_xGa_{1-x}As$ energy barriers.

In order to obtain the hot-electron mean free path L from these data we first express the responsivity to a good approximation as[14,15] $R = (e/h\nu)\eta g$, where $\eta = (1 - e^{-2\alpha l})/2 = 25\%$ is the quantum efficiency (including the reflection off the top metal contact) and g is the optical gain derived below. The probability p of an electron generated at x_0 traveling distance x is $p = (1/L) \times e^{(x_0 - x)/L}$, where L is the mean free path. Thus, the transport distance Δx for an electron generated between 0 and l and collected at l is

$$\Delta x = \int_{x_0}^{l} (x - x_0) p\, dx + \int_{l}^{\infty} (l - x_0) p\, dx . \qquad (1)$$

For uniformly generated electrons, the average transport distance Δx_{av} is given by

$$\Delta x_{\text{av}} = \frac{1}{l} \int_0^l \Delta x\, dx_0 . \qquad (2)$$

Thus, the normalized transport or optical gain $g = \Delta x_{\text{av}}/l$ is

$$g = \left(\frac{L}{l}\right) \left[1 - \left(\frac{L}{l}\right) (1 - e^{-l/L}) \right] , \qquad (3)$$

where for our sample the superlattice length is $l = 1.8$ μm. Using Eq. (3) together with the data from Fig. 2 we can obtain the bias dependence of the hot-electron mean free path L (shown in Fig. 3). Note that the striking resonance in L is even more pronounced than the corresponding enhancement in the responsivity. This is due to the exponential relation between g and L.

We now analyze the origin of this resonance in L. As discussed previously,[15,16] the hot-electron mean free path is very long (i.e., many periods) in these structures and thus the transport of the photoexcited electron above the top of the energy barriers is in a quasiequilibrium condition. That is, on the average the energy gained from the potential drop through the superlattice is approximately balanced by the energy loss to LO phonons (having an energy $\hbar\omega_0$). Thus, an electron will propagate with an average energy $\Delta E \simeq \frac{1}{2} \hbar\omega_0$ above the top of the energy barriers (as schematically illustrated in Fig. 4). This is possible since the short electron LO-phonon emission time[17-19] $\tau_{\text{LO}} = 170$ fs is comparable to the transit time across a period. One can calculate the bias-dependent mean free path $L(V)$ for this transport in terms of the hot-electron superlattice transmission coefficient $T(V)$ using

$$L(V) = v(V) \tau_w(V) T(V) , \qquad (4)$$

$$v(V) = \mu E / [1 + (\mu E / v_s)^2]^{1/2} , \qquad (5)$$

$$\tau_w(V) = (\tau - \tau_0)(v/v_s) + \tau_0 , \qquad (6)$$

where $v(V)$ is the velocity of the hot electrons (given in terms of the mobility[20] $\mu = 500$ cm^2/V s, the electric field $E = V/l$, and the saturation velocity[21] $v_s = 1 \times 10^7$ cm/s, and $\tau_w(V)$ is the well recapture time [given in terms of the high- (τ) and low-bias (τ_0) values[22]]. This dependence on v accounts for the expected increase in τ_w as the transit time and hence capture rate decreases. The linear dependence of L on $T(V)$ [shown in Eq. (2)] is a simplification obtained from the expected linearity of the velocity on $T(V)$. That is, the velocity is approximately proportional to the superlattice energy bandwidth δE

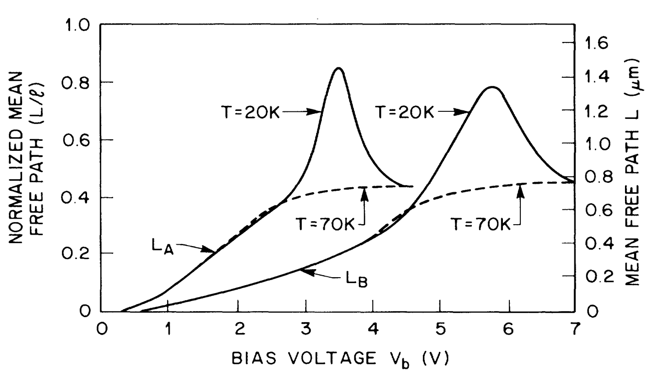

FIG. 3. Experimental hot-electron mean free path L (for both samples A and B) as a function of bias voltage, for $T = 20$ and 70 K.

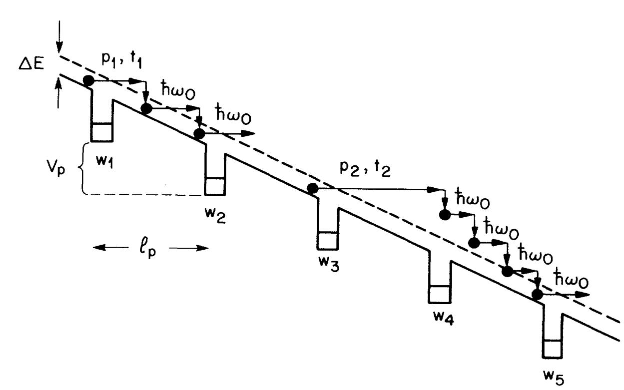

FIG. 4. Schematic superlattice band diagram showing electron transport and phonon emission.

(i.e., $v \propto \delta E$),[23] which is in turn proportional to the transmission coefficient [i.e., $\delta E \propto T(V)$].[24] In contrast, τ_w is not expected to depend strongly on $T(V)$ since it is primarily a single-quantum-well property describing the capture of the hot electron from the lowest continuum state above the top of the well into a quantum-well bound state.

Since $v(V)$ and $\tau_w(V)$ are nearly constant at high bias they do not significantly affect the shape of the resonance in L, which is produced by the barrier transmission resonances. To calculate $T(V)$, consider an electron being transported through the superlattice and having an average energy ΔE above the top of the barriers. Since, on the average, the electron will lose one or more phonons of energy per period, it will usually (with probability p_1) traverse a single well (having a transmission coefficient t_1), emit at least one phonon, and then incoherently traverse the next well (as schematically illustrated in Fig. 4 for the electron near well w_1). However, some of the time (with probability p_2) an electron can coherently traverse two quantum wells (transmission coefficient t_2) with no energy loss (i.e., ballistically) and then emit sufficient phonons to return to an average energy ΔE (schematically illustrated in Fig. 4 for the electron near well w_3). It is this ballistic transport across a pair of wells which leads to the coherent transmission resonance in the mean free path. An exact calculation of the average superlattice transmission coefficient $T(V)$ in terms of p_1, p_2, t_1, and t_2 has been done.[25] However, one can use the much simpler expression (which agrees numerically to within a few percent) given below

$$T(V)=[(1-p_2)+p_2t_2(V)]^N, \tag{7}$$

where N is the number of *pairs* of quantum wells traversed before recapture and where we have simplified the above expression using the calculated result $t_1 \simeq 1$. This expression for $T(V)$ is just the sum of the probability-weighted transmission coefficients multiplied together for N-independent events. The coherent transmission coefficient for a pair of wells, $t_2(V)$, is calculated[26] using the transfer-matrix method in an analogous way to that used for the transmission of a pair of barriers.[27] By now combining Eqs. (4)-(7) and solving self-consistently at each bias voltage using $L(V) = 2l_pN(V)$ (where $l_p = 365$ Å is the superlattice period) we can calculate the bias dependence of the mean free path. The result is shown in Fig. 5 for two values of the probability for coherent transport $p_2 = 0.3$ and 0.05 which have been adjusted (together with τ) to give the best fit with experiment. The good agreement with the measurements shown in Fig. 3 is evident with the position, shape, and temperature dependence of the hot-electron mean free path $L(V)$ being well explained. Note that, as previously mentioned, the high-field well recapture time τ is significantly larger than the low-field value τ_0 and is, in fact, comparable to the electron transit time $\tau_T = l/v$ across the superlattice. This is the reason that the hot-electron mean free path is so long and thus the optical gain close to its maximum value of $g = \frac{1}{2}$.

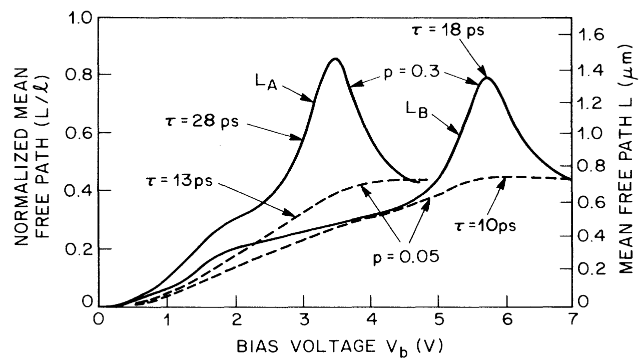

FIG. 5. Calculated mean free path L as a function of bias voltage (for both samples A and B). The well recapture times τ and coherent probabilities p_2 are indicated.

The *ballistic* mean free path[28,29] is defined in terms of the coherent probability p_2 using $p_2 = \exp(-l_p/\Lambda_b)$. Substituting $p_2 = 0.3$ at $T = 20$ K and $p_2 = 0.05$ at $T = 70$ K we determine $\Lambda_b = 300$ and 120 Å at these temperatures which is in satisfactory agreement with other determinations.[28,29] It is this coherent ballistic transport and the associated transmission resonances that is responsible for the unusual crossing of the different temperature and wavelength curves mentioned previously (Figs. 2 and 3).

In conclusion, we have presented the first direct measurement of a striking enhancement in the hot-electron recapture mean free path, and shown that it is caused by the coherent double-well transmission resonances in the superlattice continuum states. The strength of this resonance demonstrates that the hot electrons remain in the Γ valley as a results of the short phonon emission time ($\tau_{LO} = 170$ fs) and resulting rapid phonon energy loss. This new understanding of the transport physics can lead to improved quantum-well infrared photodetectors.[14,15,30,31]

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