FAST DATA CODING USING MODULATION OF INTERBAND OPTICAL PROPERTIES BY INTERSUBBAND ABSORPTION IN QUANTUM WELLS

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ABSTRACT. Intersubband absorption of radiation by a two-dimensional electron gas can be used to control the electron temperature and effect a significant modulation of the interband optical properties of the semiconductor in the quantum well. We discuss the implementation of a fast modulator of infrared radiation for fiber-optical communications as well as the formation of powerful and short single-mode infrared pulses. The method can also be used to modulate laser radiation by controlling simultaneously the pumping current and the optical gain in the active region. The dual modulation method allows to eliminate the relaxation oscillations and suppress the wavelength chirping in optical communication systems operating at high pulse repetition rates.

1. Introduction

It is well known that the interband (IB) absorption coefficient g in a degenerate semiconductor is a strong function of the effective carrier temperature T_e. It has been proposed[1] to use this effect for a rapid control of the semiconductor-laser gain by driving a lateral electron-heating current through the active region. Inasmuch as T_e can be independently controlled and rapidly modulated, one can envisage a number of useful applications in optical communications. We shall review some of the recent ideas for such applications, specializing to the case when the control of T_e is effected by intersubband absorption.[2] The idea is to control the optical transparency of a multiple quantum well (QW) waveguide by modulating the effective temperature of free carriers in the wells.

2. Control of Electron Temperature by Intersubband Absorption

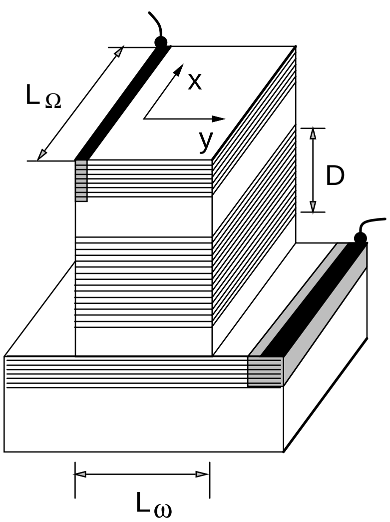
Consider a multiple quantum well (MQW) irradiated by infrared photons $\hbar\omega$ (e.g. from a CO_2 laser, $\hbar\omega = 124\,\mathrm{meV}$) nearly resonant (within $\Delta\omega$) with the intersubband energy. The waveguide structure is illustrated in Fig. 1. The interband (IB) radiation propagates along x and the intersubband (ISB) along y directions. The ISB radiation in the waveguide is assumed TM polarized and its photon flux inside a QW will be denoted by Φ_ω. In a steady state, T_e is determined by the balance equation

$$k\,\Delta T_e = w\,\tau_\varepsilon\,\hbar\omega\;, \tag{1}$$

where τ_ε is the energy relaxation time of QW electrons, $\Delta T_e \equiv T_e - T$, and w is the ISB transition rate. Evaluating w by the Golden rule, we have[3]

$$w = \frac{e^2 Q}{2\gamma}\,\frac{R_0}{m\,\bar{n}}\,\Phi_\omega\,\frac{\gamma^2}{(\Delta\omega)^2+\gamma^2}\;, \tag{2}$$

where $R_0 \equiv 377\,\Omega$ is the vacuum impedance, m the electron effective mass, and $\bar{n}$ the refractive index; $Q \approx 1$ is the oscillator strength and γ the width of the ISB resonance. A typical power flux density $P_\omega = 10\,\mathrm{kW/cm^2}$ of a CO_2 laser corresponds to $\Phi_\omega = P_\omega/\hbar\omega \approx 5\times10^{23}\,\mathrm{cm^{-2}\,s^{-1}}$. We shall be considering the case of an InGaAs QW ($m = 0.041\,m_0$). Taking $\hbar\gamma \approx 4\,\mathrm{meV}$ (from the data[4] in GaAs QW at 300 K) and $\tau_\varepsilon \approx 6\,\mathrm{ps}$ (from the energy loss rate of 7 meV/ps measured[5] in InGaAs/InP QW at $T_e = 500\,\mathrm{K}$ for high carrier densities), Eq. (1) gives $\Delta T_e \approx 260\,\mathrm{K}$ at resonance. Even longer hot electron-hole plasma cooling times (corresponding to $\tau_\varepsilon \approx 20\,\mathrm{ps}$) have been reported[6] in InGaAs/InAlAs MQW at $T_e \lesssim 700\,\mathrm{K}$ and sheet carrier concentrations of $2\times10^{12}\,\mathrm{cm^{-2}}$. The relatively long energy relaxation time is usually explained by phonon bottleneck effects.[7]



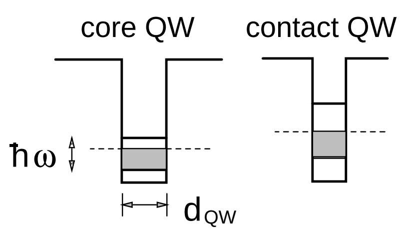


Fig. 1: Schematic diagram of the modulator (after Ref. 2). Dimensions: $L_\Omega = 500\,\mu\mathrm{m}$, $L_\omega = 3\,\mu\mathrm{m}$, $D = 1.5\,\mu\mathrm{m}$. The core (D) contains 50 QW's of width $d_{QW} \approx 100\,\mathrm{\mathring{A}}$ (intersubband separation $\hbar\omega \approx 0.12\,\mathrm{eV}$). Narrow contact QW's do not absorb TM-polarized ISB radiation.

A desired variation of T_e can be accomplished by varying either the intensity Φ_ω or the frequency $\Delta\omega$. At a constant CO_2 power and frequency it is possible to vary the detuning $\Delta\omega$ by Stark shifting the ISB resonance. For this purpose, we need means for applying an electric field in the z direction. To minimize the loss of ISB power by free-carrier absorption in the contact layers controlling the electric field, the use of doped bulk regions should be avoided. Instead, the contact layers can themselves be multiple quantum wells, doped or modulation-doped to a high conductivity and narrow enough that the intersubband resonance in the contact QW's is far above the ISB photon energy $\hbar\omega$.

While the relatively long energy relaxation time enhances the amplitude of the T_e response under quasi-static modulation, cf. Eq. (1), it degrades the high-frequency characteristics (above $1/\tau_\varepsilon$). The deviation $\Delta T \equiv T_e - T$ of the carrier temperature T_e from the ambient temperature T satisfies an equation of the form

$$\frac{\mathrm{d}\Delta T}{\mathrm{d}t} = G - \frac{\Delta T}{\tau_\varepsilon}\ , \tag{3}$$

where $G(t) \equiv \hbar\omega\, w(t)$ is the absorbed power. Solution of Eq. (3), satisfying $\Delta T = 0$ at $t = 0$, is given by

$$\Delta T(t) = \mathrm{e}^{-t/\tau_\varepsilon} \int_0^t G(t')\, \mathrm{e}^{t'/\tau_\varepsilon}\, \mathrm{d}t'\ . \tag{4}$$

For a sinusoidal $G(t) = ½\, G_0 (1 - \cos\omega t)$, we have

$$\Delta T(t) = ½\, G_0 \tau_\varepsilon \left[1 - \cos\gamma \cos(\omega t - \gamma) - \mathrm{e}^{-t/\tau_\varepsilon} \sin^2\gamma\right]\ , \tag{5}$$

where γ is defined by $\tan\gamma \equiv \omega\tau_\varepsilon$. The quantity $T_\infty \equiv G_0 \tau_\varepsilon$ represents the temperature by which T_e would be raised by a constant G_0 equal to the signal amplitude. Figure 2 illustrates the time-dependent heating of carriers by a sinusoidal ISB irradiation, $\Phi_\omega \sin^2(\pi f t)$, where we assume $P_\omega \equiv \hbar\omega\Phi_\omega = 10^4\ \mathrm{W/cm^2}$.

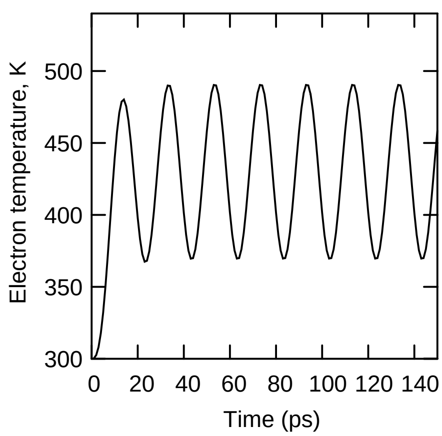


Fig. 2a: Evolution of the electron temperature T_e under CO_2 laser irradiation with $f = 50\,\mathrm{GHz}$. For $t < 0$ we assume $T_e = T_0 = 300\,\mathrm{K}$.

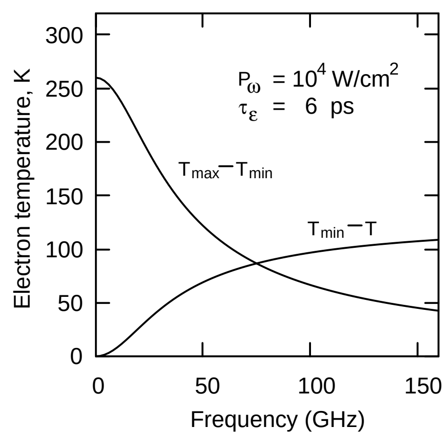


Fig. 2b: Frequency dependence of the T_e modulation amplitude and of the minimum T_e during a period for a sinusoidal modulation.

Next consider an input signal $G(t)$ consisting of a series of Gaussian pulses

$$G = \sum_{n=0}^{\infty} G_0(t - nT)\ , \quad \text{where } G_0(t) = (T_\infty/\tau_\varepsilon)\, \mathrm{e}^{-t^2/2\Delta t^2} \tag{6}$$

and T is the period, $T = 1/f$, with f being the pulse repetition rate. Assuming that $\Delta T(t) = 0$ for $t < 0$, we can use Eq. (4) to find:

$$\Delta T(t) = \sqrt{\frac{\pi}{2}}\, T_\infty \,\frac{\Delta t}{\tau_\varepsilon} e^{\Delta t^2/2\tau_\varepsilon^2} \sum_{n=0}^{\infty} e^{(nT - t)/\tau_\varepsilon}$$
$$\times \left[\Phi\left[\frac{nT + \Delta t^2/\tau_\varepsilon}{\Delta t\sqrt{2}}\right] + \Phi\left[\frac{t - (nT + \Delta t^2/\tau_\varepsilon)}{\Delta t\sqrt{2}}\right]\right], \tag{7}$$

where Φ is the error integral $[\Phi(x) \to 1$ for $x \gtrsim 1$ and $\Phi(-x) = -\Phi(x)]$. The sum of two Φ's in the square brackets decreases with n much faster than $\exp(-nT/\tau_\varepsilon)$, so the overall sum converges very rapidly. Results of calculations with the formula (7) are plotted in Fig. 3.

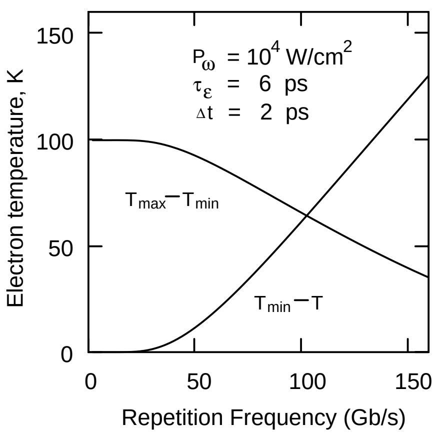


Fig. 3: Frequency dependence of the T_e modulation amplitude and the minimum T_e during a period for a sequence of Gaussian pulses of width $\Delta t = 2$ ps.

It is evident from Figs. 2b and 3 that the modulation amplitude drops at high frequencies and the mean T_e rises. We see that the coding speed of a T_e modulator is limited by ~ 50 Gb/s. At higher bit rates, the temperature of QW electrons does not have time to relax to the lattice temperature T_0 in between two heating pulses.

It can be shown that when the pulse is sufficiently narrow $\Delta t \ll \tau_\varepsilon$ then the difference $T_{max} - T_{min}$ does not depend on the repetition rate. This is easy to demonstrate as follows: for $\Delta t \ll \tau_\varepsilon$, we can replace the Gaussian by a δ function:

$$e^{-\frac{(t-nT)^2}{2\Delta t^2}} \approx \sqrt{2\pi}\,\Delta t\,\delta(t - nT)\,.$$

For a single δ-function pulse, the response is of the form

$$\Delta T_n(t) = T_\infty \sqrt{2\pi}\,\frac{\Delta t}{\tau_\varepsilon}\, e^{(nT-t)/\tau_\varepsilon}\,\Theta(t - nT)\,, \tag{8}$$

where $\Theta(x)$ is a unit step function. We must add responses to all pulses preceding t. For $t \gg T/\tau_\varepsilon$, the summation can be extended to ∞, with the result

$$\Delta T(t) = T_\infty \sqrt{2\pi}\,\frac{\Delta t}{\tau_\varepsilon}\,\frac{e^{-t_T/\tau_\varepsilon}}{1 - e^{-T/\tau_\varepsilon}}\,, \tag{9}$$

where t_T should be understood as the "residual time" past the last pulse, $t_T = t\,[\mathrm{mod}\,T\,]$. Equation (9) can also be derived directly from (7). Any function of t_T is obviously periodic in t. The minimum occurs for $t_T = 0$ and the maximum for

$t_T = T - \varepsilon$, where ε is an infinitesimal. From Eq. (9) we find that $T_{\min} \approx T$ for low frequency, $T \gtrsim \tau_\varepsilon$, and $T_{\min} \propto f$ for $T = 1/f \ll \tau_\varepsilon$. The $T_{\rm e}$ response amplitude equals $T_{\max} - T_{\min} = T_\infty \sqrt{2\pi}\, \Delta t/\tau_\varepsilon$ and does not depend on the repetition rate.

3. Control of the Gain in Steady State

Device length L_ω in the y direction should be chosen short enough that Φ_ω be approximately uniform, $\alpha L_\omega < 1$, where α is the ISB absorption coefficient,

$$\alpha = \frac{r\,\Gamma_\omega\, n_{\rm S}}{d_{\rm QW}}\;\frac{w}{\Phi_\omega} = \frac{\Gamma_\omega\, n_{\rm S}}{d}\;\frac{e^2 Q}{2\gamma}\;\frac{R_0}{m\,\overline{n}}\;\frac{\gamma^2}{(\Delta\omega)^2 + \gamma^2}\ , \tag{10}$$

$d = d_{\rm QW} + d_{\rm B}$ is the MQW period, $d_{\rm QW} \equiv r\,d$ the QW thickness, $n_{\rm S}$ the electron sheet concentration per period, and Γ_ω the confinement factor for the ISB radiation intensity. For an efficient operation of the modulator it is important that the steady-state density $p_{\rm S}$ of holes generated by the IB radiation be small compared to $n_{\rm S} = n_0 + p_{\rm S}$. The equilibrium electron density n_0 (introduced by doping) is limited by the requirement that the Fermi level be less than the ISB separation, $E_{\rm F} < \hbar\omega$. At $n_0 = 2\times10^{12}\,{\rm cm}^{-2}$ in the range of 300 K to 500 K one has $E_{\rm F} \approx 115\,{\rm meV}$. Higher n_0 would result in a diminishing efficiency of carrier heating due to increasing population of the second subband.

The modulator length L_Ω in the x direction should be chosen so as to achieve a desired modulation depth of the transmitted IB beam ${\rm e}^{r\,\Gamma_\Omega\, g\, L_\Omega}$, where Γ_Ω is the waveguide confinement factor for the IB radiation intensity. The gain function g is of the form

$$g\,(T_{\rm e}, T_{\rm h}, n_{\rm S}, p_{\rm S}, \hbar\Omega) = g_{\max}\,(f_{\rm e} + f_{\rm h} - 1)\ , \tag{11}$$

where $f_{\rm e}$ and $f_{\rm h}$ are the Fermi functions of electrons and holes, respectively, at energies selected by the incident photons $\hbar\Omega$ inducing transitions between the heavy-hole and the lowest electron subbands. The value of $g_{\max}$ in an InGaAs QW is, typically,[8] $g_{\max} \approx 10^3\,{\rm cm}^{-1}$. For transitions at the fundamental absorption edge in the QW, the Fermi factors are given by

$$f_{\rm e}\,(n_{\rm S}, T_{\rm e}) = 1 - {\rm e}^{-\pi\hbar^2 n_{\rm S}/m\,kT_{\rm e}}\ ; \tag{12a}$$

$$f_{\rm h}\,(p_{\rm S}, T_{\rm h}) = 1 - {\rm e}^{-\pi\hbar^2 p_{\rm S}/m_{\rm h}\,kT_{\rm h}}\ , \tag{12b}$$

where $m_{\rm h} \approx 0.5\,m_0$ and $T_{\rm h}$ are the heavy-hole effective mass and temperature, respectively. Equations (12) are derived from the exact expression for the Fermi integral,

$$n_{\rm S} = \frac{m\,kT_{\rm e}}{\pi\hbar^2}\,\ln\left[1 + {\rm e}^{E_{\rm F}/kT_{\rm e}}\right]\ , \tag{13}$$

in a 2D electron gas with one subband occupied. The electron quasi-Fermi level is

$$E_{\rm F}\,(n_{\rm S}, T_{\rm e}) = kT_{\rm e}\,\ln\left[{\rm e}^{\pi\hbar^2 n_{\rm S}/m\,kT_{\rm e}} - 1\right]\ , \tag{14}$$

and the occupation probability at the subband bottom is given by

$$f_e(n_S, T_e) \equiv \left[1 + e^{-E_F/kT_e}\right]^{-1} = 1 - e^{-T_n/T_h} . \tag{15}$$

For electrons in InGaAs ($m_e = 0.041\, m_0$) the temperature scale in the exponent equals

$$T_n \equiv \frac{\pi \hbar^2 n_S}{m_e k} = 676.62\,\mathrm{K} \times \frac{n_S\,[\mathrm{cm}^{-2}]}{10^{12}} \tag{16}$$

Similarly, for holes ($m_h = 0.5\, m_0$) one has $f_h(p_S, T_h) = 1 - \exp(-T_p/T_h)$, where

$$T_p \equiv \frac{\pi \hbar^2 p_S}{m_h k} = 55.5\,\mathrm{K} \times \frac{p_S\,[\mathrm{cm}^{-2}]}{10^{12}} \tag{17}$$

Except at cryogenic temperatures, the 2D hole gas is nondegenerate even at $10^{12}\,\mathrm{cm}^{-2}$.

The modulator can be expected to perform up to frequencies limited by the inverse energy relaxation time τ_ε, provided the slower processes associated with carrier generation by the IB radiation make negligible contribution to g. At a given value of n_0 the latter requirement puts a limit on the IB flux that is modulated.

To estimate this limit and calculate the temperature dependence of g in a steady-state, we consider the rate equations:

$$\frac{\mathrm{d}p_S}{\mathrm{d}t} = -\,\bar{c}\, g\, S - R_S\, n_S\, p_S\ ; \tag{18a}$$

$$\frac{\mathrm{d}S}{\mathrm{d}t} = (r\Gamma_\Omega)(\bar{c}\, g)\, S - \tau_{ph}^{-1}\,(S - S_0)\ . \tag{18b}$$

Here S is the photon density per unit area in a single QW, $S_0 \equiv \Phi_\Omega\, d_{QW}/\bar{c}$, where Φ_Ω is the incident IB photon flux, $\bar{c} \equiv c/\bar{n}$ is the speed of light, and $\tau_{ph} = L_\Omega/\bar{c}$. The radiative recombination coefficient $R_S \equiv B/d_{QW}$, where $B \approx 10^{-10}\,\mathrm{cm}^3/\mathrm{s}$.

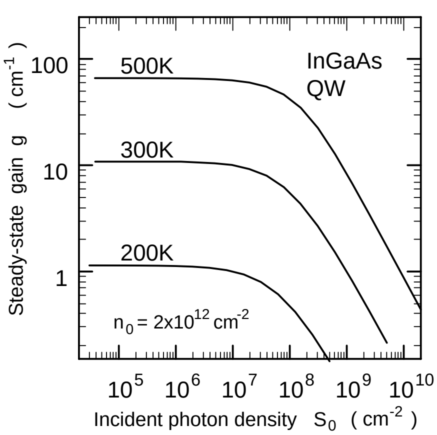


Fig. 4: Dependence of the steady-state gain on the photon density S_0 of incident interband radiation at different carrier temperatures. The assumed parameters: $g_{max} = 10^3\,\mathrm{cm}^{-1}$, $r\,\Gamma_\Omega = 0.3$, and $R_S = 10^{-4}\,\mathrm{cm}^2/\mathrm{s}$.

Figure 4 shows the stationary gain as a function of incident power, calculated from Eqs. (18) at several carrier temperatures, assuming $n_S = n_0 + p_S$ with $n_0 = 2 \times 10^{12}\,\mathrm{cm}^{-2}$ and $\tau_{ph} = 5\,\mathrm{ps}$ ($L_\Omega \approx 500\,\mu\mathrm{m}$). We take $T_h = T_e$, rather arbitrarily. Results of our calculation are not very sensitive to the choice of the hole temperature in the range $T \lesssim T_h \lesssim T_e$. It is evident from Fig. 4 that for $S_0 \lesssim 10^8\,\mathrm{cm}^{-2}$ carrier generation effects can be neglected. The total modulated power $P_\Omega = \Gamma_\Omega^{-1} A\, \Phi_\Omega\, \hbar\Omega$ is related to S_0 by

$$r\,\Gamma_\Omega\, P_\Omega = N\, L_\omega\, \hbar\Omega\, \bar{c}\, S_0\ ,$$

where $A = D \times L_\omega$ is the MQW cross-sectional area, D the core thickness, and $N = D/d$ the number of periods. Taking $N = 50$, $r\,\Gamma_\Omega \approx 0.3$, and $L_\omega \approx 3\,\mu\text{m}$, we find that $S_0 = 10^8\,\text{cm}^{-2}$ corresponds to a maximum modulated power of $P00 \approx 6\,\text{mW}$. For $P_\Omega \approx 1\,\text{mW}$, the steady-state g varies from $g \approx -9\,\text{cm}^{-1}$ at $T_e = 300\,\text{K}$ to $g \approx -61\,\text{cm}^{-1}$ at $T_e = 500\,\text{K}$. For $L_\Omega = 0.5\,\text{mm}$ and $r\,\Gamma_\Omega \approx 0.3$ this corresponds to a 3.5 dB modulation.

At a higher power (or lower n_0) the modulator efficiency suffers from self-induced transparency effects associated with the accumulation of electrons and holes. In the next Section, we show that these effects can be used advantageously for the formation of short high-power IB radiation pulses. For this purpose, the MQW need not be doped.

4. Formation of short pulses

Consider the situation arising at a high S_0 in the presence of ISB absorption. In what follows, we assume that the MQW is undoped, $n_0 = 0$. In the steady state there is a large number of electrons and holes $p_S(S_0, T_e)$, readily evaluated from Eqs. (18). In this state the gain has a small negative value $g(S_0, T_e)$. If the carrier heating is now abruptly terminated (by chopping Φ_ω or by shifting the ISB resonance with an external electric field) then T_e rapidly goes down to the ambient temperature and the gain function becomes temporarily positive. The excess carriers undergo stimulated recombination accompanied by a large pulse in the IB photon density. An example of such a pulse is shown in Fig. 5 The pulse shape is calculated from Eqs. (18), assuming that the initial electron heating is stopped at $t = 10\,\text{ps}$ and the carrier temperature relaxes from $T_e = 500\,\text{K}$ to $T_e = 300\,\text{K}$, according to $\Delta T_e(t) = \Delta T_e\, e^{-t/\tau_\varepsilon}$ with $\tau_\varepsilon = 6\,\text{ps}$. At longer times t the carrier density and the gain approach their new steady state values $p_S(S_0, T)$ and $g(S_0, T)$, respectively.

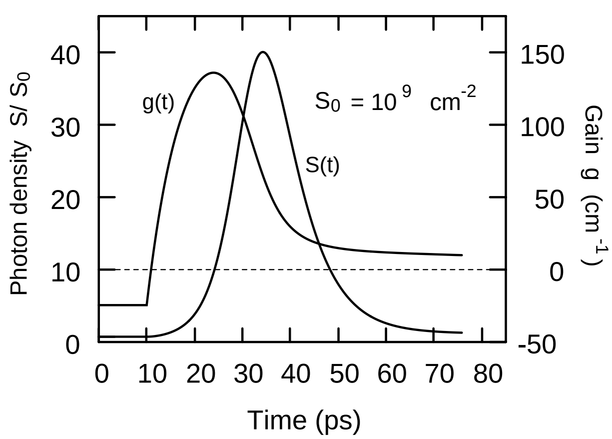


Fig. 5: Formation of short pulses by abrupt termination of carrier heating. In the presence of both IB and ISB radiations, the device is allowed to reach a steady-state with assumed carrier temperature $T_e = 500\,\text{K}$. At $t = 10\,\text{ps}$ the ISB absorption terminates abruptly and T_e relaxes to the ambient temperature $T = 300\,\text{K}$. Full width at half maximum of the pulse S is 15 ps.

The pulse shown in Fig. 5 has a full width at half maximum $\Delta t_{\text{fwhm}} = 15\,\text{ps}$ and a peak photon density of $4 \times 10^{10}\,\text{cm}^{-2}$ – corresponding to a power of 2.4 W (the pumping level is only 60 mW). It is interesting to note that the peak power varies only by 30 % over the decade variation in the pump power $10^9 \lesssim S_0 \lesssim 10^{10}\ \mathrm{cm}^{-2}$. The width Δt_{fwhm} is practically constant over the same range. Increasing S_0 in this range mainly leads to a faster device recovery in preparation for the next pulse.

Indeed, when the carrier heating is turned on, the negative gain function temporarily increases in magnitude. This results in an enhanced absorption of IB radiation and the number of carriers increases back toward the steady-state value $p_S(S_0, T_e)$. During this relatively slow process the power is stored for the next pulse. The storage time depends on the incident power and scales approximately as $1/S_0$ for $S_0 \gtrsim 10^9\ \mathrm{cm}^{-2}$. For a given value of S_0 and a given swing in the carrier temperature, the peak power is inversely proportional to its duration.[2]

5. Fast Data Coding

High-frequency characteristics of the modulator for fast data coding depend on the inertia of the response of T_e to the heating signal, discussed in Sect. 2, and on the inertia associated with carrier accumulation and depletion in the MQW, discussed in Sect. 3. For error-free information coding, it is essential that the response to an isolated pulse should be similar to that to a sequence of pulses following each other at a given bit rate. As we have seen in Fig. 3, the speed of encoding the electron temperature is limited by $\sim$50 Gb/s – even without a consideration of the optical response. For not too high optical IB power, this turns out to be also the coding speed limitation of a T_e controlled optical modulator.

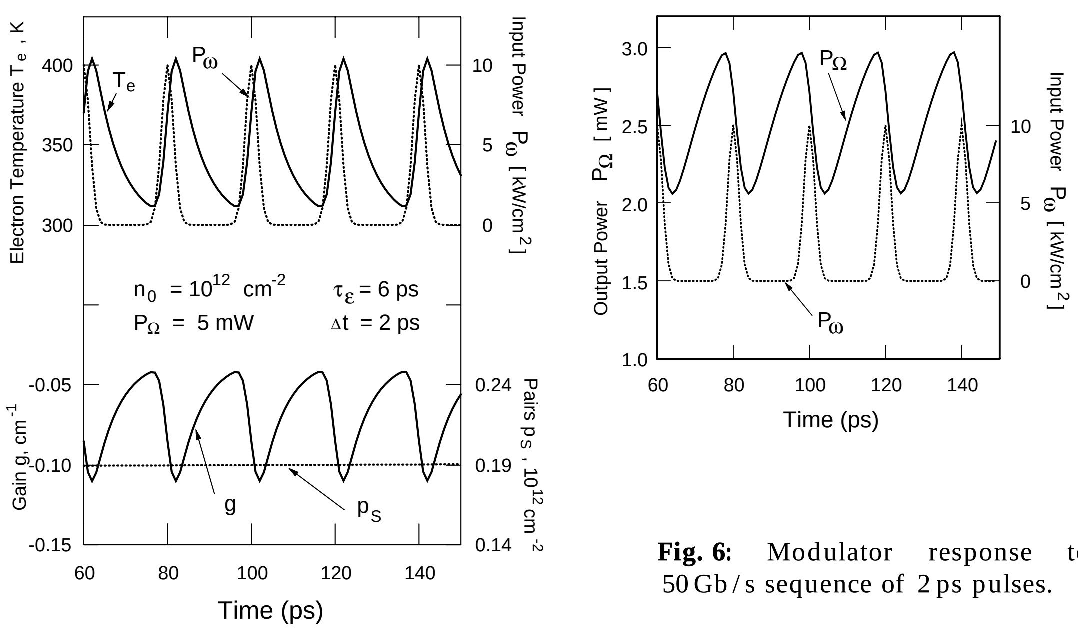


Fig. 6: Modulator response to a 50 Gb/s sequence of 2 ps pulses.

Figures 6 show the calculated room-temperature response of the modulator – the carrier temperature $T_e(t)$, the density $p_S(t)$ of photogenerated pairs, the optical gain $g(t)$, and the transmitted IB light power $P_{\mathrm{out}}(t)$ – subject to a periodic sequence of heating pulses of the form $P_\omega \exp(-t^2/\Delta^2)$, with $P_\omega = 10^4\ \mathrm{W/cm^2}$ and $\Delta = 2$ ps, repeated at the rate of 50 GHz. At this repetition frequency, for not too high incident IB power ($\lesssim 35\,\mathrm{mW}$), the response is nearly identical to that for an isolated pulse of the same form. This means that the effect of accumulation of photogenerated pairs, produced during the the passage of a heating pulse, is relatively small ($p_S \ll n_S = n_0 + p_S$), where n_0 is the electron concentration in the QW's in the absence of IB radiation. The optimum value of n_0 was found to be about $10^{12}\,\mathrm{cm}^{-2}$. For the incident power above $\sim 35\,\mathrm{mW}$, the accumulation of pairs becomes important, so that the maximum coding rate becomes limited by the recombination processes and drops below $10\,\mathrm{Gb/s}$.

6. High-frequency modulation and suppression of chirp in semiconductor lasers

The common method of modulating the laser radiation amplitude by varying the pumping current suffers from two drawbacks. First, it is limited to relatively low frequencies ($f \lesssim 10\,\mathrm{GHz}$). Second, for $f \gtrsim 1\,\mathrm{GHz}$, it is plagued by oscillations in the wavelength of the dominant mode (chirp). Both of these problems arise from the relaxation oscillations due to an intrinsic resonance in the nonlinear laser system (the electron-photon resonance).

An alternative principle for modulating the laser output is to directly control by external means the gain coefficient g_0 of the active medium by varying the effective carrier temperature T_e in the laser active region.[1] High-frequency modulation of T_e by several tens of degrees has been demonstrated experimentally,[9] by driving an electric current through the active region. Although this method in principle allows a faster laser modulation, by itself it eliminates neither the relaxation oscillations nor the frequency chirp. The new approach, proposed in our recent work,[10] allows an enhancement of the coding frequency, suppressing the chirp at the same time. The key idea consists in a coherent combination of *two* independent means of controlling the output radiation: the pumping current I and the effective carrier temperature T_e.

Consider a semiconductor laser, subject to a time-dependent pumping current $I(t)$ and an external electron-heating power $G(t)$. The nonlinear system can be described by the standard rate equations[8] for the carrier density n and the photon density S per unit volume:

$$\frac{\mathrm{d}n}{\mathrm{d}t} = J - S\tilde{g} - Bn^2 \quad ; \tag{19a}$$

$$\frac{\mathrm{d}S}{\mathrm{d}t} = (\Gamma\tilde{g} - \tau_{\mathrm{ph}}^{-1})S + \beta B n^2 \quad , \tag{19b}$$

coupled with an energy balance equation (3) for T_e. In these equations, J is the electron flux per unit volume of the active layer, $B \approx 10^{-10}\,\mathrm{cm^3/s}$ is the radiative recombination coefficient, $\beta \approx 10^{-4}$ the spontaneous emission factor, and $\tilde{g} \equiv g\,\bar{c}$ the optical gain $[\mathrm{sec}^{-1}]$ in the active layer.

We shall use the notation, $X(t) \equiv \overline{X} + \hat{X}\,e^{i\omega t}$, for harmonically varying quantities $X(t)$. Let us carry out a small-signal analysis of the system (19), linearizing it about a steady state well above the lasing threshold:

$$\mathrm{i}\,\omega\,\hat{n} \;=\; \hat{J} - \overline{S}\;\hat{g} - \overline{g}\;\hat{S} - 2B\,\overline{n}\;\hat{n} \quad ; \tag{20a}$$

$$\mathrm{i}\,\omega\,\hat{S} \;=\; (\Gamma\tilde{g} - \tau_{\mathrm{ph}}^{-1})\,\hat{S} + \Gamma\overline{S}\;\hat{g} + 2\beta B\,\overline{n}\;\hat{n} \quad . \tag{20b}$$

Assume that the absorbed power $G(t)$ per carrier depends only on the ISB absorption rate w, Eq. (2), the variation of T_{e} is of the form

$$\hat{T}_{\mathrm{e}} \;=\; \frac{G_w'\;\hat{w}\;\tau_\varepsilon}{1 + \mathrm{i}\,\omega\,\tau_\varepsilon} \tag{21}$$

and decouples from the rest of the system. We are neglecting here possible dependences of the carrier distribution function on the photon density S due to the spectral hole burning and/or direct carrier heating by the lightwave[11, 12]. Such effects can be taken into account in a similar fashion.

$\hat{T}_{\mathrm{e}}$ will be considered an independent variable in Eqs. (20a) and (20b). The small signal variation of the gain can be described by two coefficients g_n' and g_T', each of which depends on the steady-state values of concentration $\overline{n}$ and temperature $\overline{T}_{\mathrm{e}}$, as well as the optical frequency Ω:

$$\hat{g} \;=\; g_n'(\overline{n}, \overline{T}_{\mathrm{e}}, \Omega)\,\hat{n} \;+\; g_T'(\overline{n}, \overline{T}_{\mathrm{e}}, \Omega)\,\hat{T}_{\mathrm{e}} \quad . \tag{22}$$

The steady-state coefficients in the laser generation regime $\overline{J} > J_{\mathrm{th}}$ are related as follows:

$$\overline{g}\;\overline{S} \;=\; \overline{J} - J_{\mathrm{th}} \;; \qquad\qquad \Gamma\overline{g}\;\tau_{\mathrm{ph}} = 1 \;, \tag{23}$$

where we have set $J_{\mathrm{th}} = B\,\overline{n}^{\,2}$, neglecting the term $\beta B\,\overline{n}^{\,2}$.

We are interested in solutions of Eqs. (20) for which the variations $\hat{n}$ and $\hat{T}_{\mathrm{e}}$ have a definite "target" relationship, $\hat{n} = \gamma\,\hat{T}_{\mathrm{e}}$. Substituting this relationship in Eqs. (20a) and (20b) and using Eq. (23), we find

$$\hat{S} \;=\; \frac{\hat{J}}{\overline{g}}\;\frac{1}{1 - \omega^2\,\tau_{\mathrm{ph}}\,\tau_\gamma + \mathrm{i}\,\omega\,\tau_{\mathrm{ph}}\,(1 + \tau_\gamma/\tau_{\mathrm{sp}})} \quad , \tag{24}$$

where

$$\tau_\gamma \equiv \frac{\gamma}{(g_T' + \gamma g_n')\,\overline{S}} \;; \qquad\qquad \tau_{\mathrm{sp}} \equiv \frac{1}{2B\,\overline{n}} \;, \tag{25}$$

For $\gamma \to \infty$, Eq. (24) goes over into the "classical" dependence $\hat{S}(\hat{J})$ corresponding to a pure modulation by the pumping current. The response function in (24) contains the usual pole corresponding to the electron-photon resonance, and the laser signal power decays as $1/\omega^2$ at high enough frequencies. If we are concerned with increasing the frequency of modulation, then the target should be chosen so as to eliminate the relaxation oscillations – which corresponds to setting $\gamma = 0 = \hat{n}$. In this case, Eq. (24) reduces to

$$\hat{S} = \frac{\hat{J}}{\overline{g}}\;\frac{1}{1 + \mathrm{i}\,\omega\,\tau_{\mathrm{ph}}} \;; \qquad\qquad |\,\hat{S}\,| \;=\; \frac{|\,\hat{J}\,|}{\overline{g}}\;\frac{1}{[\,1 + (\omega\,\tau_{\mathrm{ph}})^2\,]^{1/2}} \;, \tag{26}$$

This solution requires that the variations $\hat{J}$ and $\hat{T}_e$ be related to each other in a definite way:

$$\hat{T}_e = \frac{\hat{J}}{\bar{S}\ g'_T} \frac{i\omega\tau_{ph}}{1 + i\omega\tau_{ph}} . \tag{27}$$

When this relation is fulfilled, then there is no electron-photon resonance in the system and the modulation efficiency decays with frequency as $1/\omega$. Small-signal response of the laser output power is plotted in Fig. 7 for three types of modulation: (1) purely by the pumping current, (2) purely by the electron temperature, and (3) by their coherent combination as in Eq. (27).

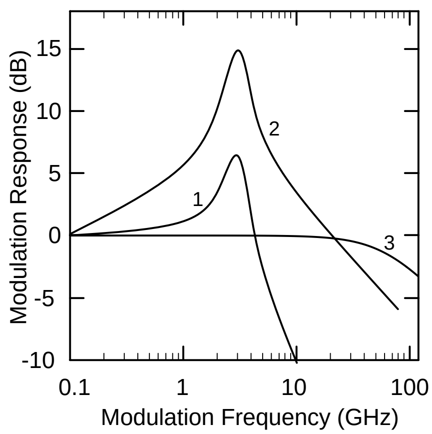


Fig. 7: Frequency dependence of the optical response $\hat{S}(f)/\hat{S}(0)$ to the variation of different parameters in a stripe MQW laser. Assumed laser parameters: five 100 Å QW's, area $360 \times 6\,\mu$m, $\Gamma = 0.2$, $\tau_{ph} = 2.5$ ps, $g'_n = 2.5 \times 10^{-16}\ \text{cm}^{-2}$, $\bar{S} = 3.5 \times 10^{14}\ \text{cm}^{-3}$. Curve 1: modulation by the pumping current, curve 2: modulation by the electron temperature, curve 3: dual modulation as in Eq. (27).

Suppression of relaxation oscillations of the carrier density makes possible a high repetition rate coding of information with short pulses. We have calculated[10] the laser response to a 10 Gb/s series of dual current-temperature pulses, and found that the response is practically undistorted, compared to the single pulse situation, It is clear that small-signal pulses δJ of *any* shape, as well as analog signals, can be transmitted in a regime of constant n, provided the system can Fourier analyze $\delta J(t)$ and form in real time an appropriate complementary pulse $\delta T_e(t)$ from the spectrum given by Eq. (27). Large signal theory of dual modulation is discussed in a forthcoming publication.[13]

In general, the choice of a target relation for dual modulation depends on the engineering problem at hand. Instead of targeting the regime of constant carrier concentration, we may be interested in the suppression of the wavelength chirping at high modulation frequencies. Recall that this unwelcome phenomenon[8] originates from the relaxation oscillations, which lead to variations in the real part η of the refractive index $\eta_c = \eta + i\kappa$ in the active region. Variations of n affect η in two ways: by changing the free-carrier absorption and by changing the optical gain. In InGaAs lasers both effects give similar contributions, shifting the lasing mode wavelength by as much as several Å.

In our present model the gain variation $\hat{g}$ has two contributions, Eq. (22). Therefore, we should distinguish the corresponding two contributions, $\hat{\eta}_{g_n}$ and $\hat{\eta}_{g_T}$ in the refractive index variation, each given by the Kramer-Kronig relation

$$\hat{\eta}_{\mathrm{g}} = -\frac{\overline{\eta}}{\pi} P \int_0^{\infty} \frac{\hat{g}(\Omega')\, d\Omega'}{\Omega'^2 - \Omega^2} \ , \tag{28}$$

where P denotes the integral principal value. We now see that it is possible to target a complete suppression of the total index oscillation,

$$\hat{\eta} = \hat{\eta}_{\mathrm{fc}} + \hat{\eta}_{g_n} + \hat{\eta}_{g_T} = 0 \ , \tag{29}$$

by judiciously choosing a relationship between $\hat{T}_{\mathrm{e}}$ and $\hat{J}$.

It should be emphasized that having targeted the complete elimination of chirp, we pay a penalty in the modulation frequency, since the electron-photon resonance is no longer eliminated and the modulation efficiency $\hat{S}/\hat{J}$ decays as ω^{-2} at sufficiently high frequencies. On the other hand, we note that the variation of η due to variations in T_{e} is usually weaker than that due to variations in n. If we compare $\hat{\eta}_{g_T}$ at the point of complete elimination of relaxation oscillations with the amount of chirp $\hat{\eta}_n = \hat{\eta}_{\mathrm{fc}} + \hat{\eta}_{g_n}$ for a purely pump-current modulation with the same output signal power, we find that the latter is usually larger, because of both the additional contribution from free-carrier absorption and the essentially different frequency dependence of $g'_T(\Omega)$ and $g'_n(\Omega)$.

7. Conclusion

We have discussed the possibility of controlling the electron temperature by intersubband absorption in a semiconductor quantum well. This gives rise to a significant modulation of the interband optical properties which can be used for the implementation of a fast and chirp-free modulator of infrared radiation. The method can also be used to modulate laser radiation by controlling simultaneously the pumping current and the optical gain in the active region. This allows to eliminate the relaxation oscillations and suppress the wavelength chirping in optical communication systems operating at high pulse repetition rates. Rapid control of intersubband absorption is possible at *a constant* CO_2 flux by varying the intersubband separation with an applied electric field. For this purpose it is crucial that means for applying the electric field would themselves be transparent to CO_2 radiation, which can be accomplished with a novel scheme discussed in Sect. 2.

Obviously, the dual modulation method is not limited to the variation of T_{e} by intersubband absorption. In principle, any of the parameters of the system of rate equations (19) may be varied externally, though not necessarily by intersubband techniques. Besides different schemes for the variation of T_{e}, several alternative proposals have been discussed recently, including varying the confinement factor[13] in a stripe laser, the photon lifetime[14] in a vertical cavity laser, and even the spontaneous emission factor[15] in a microresonator. We believe that intersubband physics has a good chance to contribute to this emerging field.

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