High-frequency modulation and suppression of chirp in semiconductor lasers

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We propose a new method for modulating laser radiation by controlling simultaneously the pumping current and the optical gain in the active region. The latter can be independently varied by modulating the effective carrier temperature. The method allows to eliminate the relaxation oscillations and enhance the modulation frequency to 50 GHz. It also allows to suppress the wavelength chirping in optical communication systems operating at pulse repetition rates of 10 Gb/s.

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 ($f \lesssim 10$ GHz). Second, for $f \gtrsim 1$ 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, e.g., by varying the effective carrier temperature T_e in the laser active region. Two high-frequency electron-heating mechanisms have been considered: (i) driving an electric current through the active region[1] and (ii) inducing intersubband absorption in quantum wells.[2] High-frequency modulation of T_e by several tens of degrees has been demonstrated experimentally.[3] 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 this work, 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.

We consider a semiconductor laser, subject to a time-dependent pumping current $I(t)$ and an external electron-heating power $P(t)$. The nonlinear system can be described by standard rate equations[4] for the carrier density n and the photon density S per unit volume as follows:

$$\frac{dn}{dt} = J - Sg - Bn^2, \tag{1a}$$

$$\frac{dS}{dt} = (\Gamma g - \tau_{\mathrm{ph}}^{-1})S + \beta Bn^2, \tag{1b}$$

coupled with an energy balance equation for T_e

$$\frac{dT_e}{dt} = P(t) - \frac{T_e - T}{\tau_\epsilon}. \tag{1c}$$

In these equations, J is the electron flux per unit volume of the active layer, $B \approx 10^{-10}$ cm^3/s is the radiative recombination coefficient, $\beta \approx 10^{-4}$ the spontaneous emission factor, $g \equiv g_0 \bar{c}$ the optical gain in the active layer, $\bar{c}$ the speed of light in the medium, Γ the confinement factor for the radiation intensity, τ_{ph} the photon lifetime in the laser cavity, and τ_ϵ the carrier energy relaxation time.

We shall use the notation, $X(t) \equiv \bar{X} + \hat{X}e^{i\omega t}$, for harmonically varying quantities $X(t)$. Let us carry out a small-signal analysis of system (1), linearizing it about a steady state well above the lasing threshold

$$i\omega\hat{n} = \hat{J} - \bar{S}\hat{g} - \bar{g}\hat{S} - 2B\bar{n}\hat{n}, \tag{2a}$$

$$i\omega\hat{S} = (\Gamma g - \tau_{\mathrm{ph}}^{-1})\hat{S} + \Gamma\bar{S}\hat{g} + 2\beta B\bar{n}\hat{n}, \tag{2b}$$

$$i\omega\hat{T}_e = \hat{P} - \tau_\epsilon^{-1}\hat{T}_e. \tag{2c}$$

Depending on a carrier heating mechanism H, the quantity P (defined as a power input per single carrier) may itself be a function of n and T_e, viz., $P = P[n,T_e;H(t)]$. For instance, in heating by an electric field F, we have $(n+p)P = eF^2[n\mu_n(T_e) + p\mu_p(T_e)]$. We assume, however, that these dependences are weak, $\hat{P} = P'_n\hat{n} + P'_T\hat{T}_e + P'_H\hat{H} \approx P'_H\hat{H}$. In this case, Eq. (2c) reduces to

$$\hat{T}_e = \frac{P'_H\hat{H}\tau_\epsilon}{1 + i\omega\tau_\epsilon} \tag{3}$$

and decouples from the rest of the system.[5] We shall regard $\hat{T}_e$ as an independent variable in Eqs. (2a) and (2b). 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 $\bar{n}$ and temperature $\bar{T}_e$, as well as the optical frequency Ω as follows:

$$\hat{g} = g'_n(\bar{n}, \bar{T}_e, \Omega)\hat{n} + g'_T(\bar{n}, \bar{T}_e, \Omega)\hat{T}_e. \tag{4}$$

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

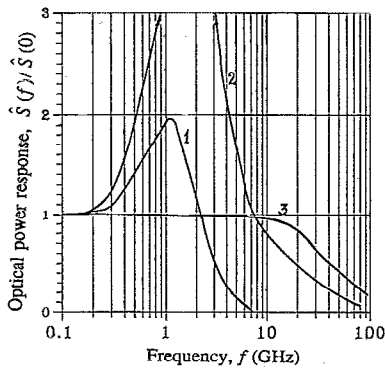$$\bar{g}\bar{S} = \bar{J} - J_{\mathrm{th}}, \quad \Gamma\bar{g}\tau_{\mathrm{ph}} = 1, \tag{5}$$

FIG. 1. Frequency dependence of the small-signal modulation depth $\hat{S}(f)$ of an InGaAs quantum well laser ($\lambda = 1.52\ \mu m$). Modulation by a pure pumping current variation $\hat{I}$ (curve 1), pure carrier temperature variation $\hat{T}_e$ (curve 2), and their coherent combination (curve 3). The steady-state pumping current is assumed $I_0 = 2I_{th}$. Laser parameters: active layer dimensions $500 \times 3 \times 0.01\ \mu m^3$, $I_{th} = 5$ mA, $\Gamma = 0.05$.

where we have set $J_{th} = B\bar{n}^2$ in system (1), neglecting the term $\beta B\bar{n}^2$. We are interested in solutions of Eqs. (2) for which the variations $\hat{n}$ and $\hat{T}_e$ have a definite "target" relationship,

$$\hat{n} = \gamma \hat{T}_e. \tag{6}$$

Substituting Eqs. (5) and (6) in Eqs. (2a) and (2b), we find

$$\hat{S} = \frac{\hat{J}}{\bar{g}} \frac{1}{1 - \omega^2 \tau_{ph}\tau + i\omega\tau_{ph}(1 + \tau/\tau_{sp})}, \tag{7}$$

where

$$\tau \equiv \frac{\gamma}{(g_T' + \gamma g_n')\bar{S}}, \qquad \tau_{sp} \equiv \frac{1}{2B\bar{n}}. \tag{8}$$

For $\gamma \to \infty$, Eq. (7) goes over into the "classical" dependence $\hat{S}(\hat{J})$ corresponding to a pure modulation by the pumping current. The response function in Eq. (7) contains the usual pole corresponding to the electron-photon resonance, and the laser signal power ($\propto \hat{S}$) decays as $1/\omega^2$ at high enough frequencies.

In general, the choice of a target relation (6) depends on the engineering problem at hand. 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. (7) reduces to

$$\hat{S} = \frac{\hat{J}}{\bar{g}} \frac{1}{1 + i\omega\tau_{ph}}, \qquad |\hat{S}| = \frac{|\hat{J}|}{\bar{g}} \frac{1}{[1 + (\omega\tau_{ph})^2]^{1/2}}. \tag{9}$$

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

$$\hat{T}_e = \frac{\hat{J}}{\bar{S}g_T'} \frac{i\omega\tau_{ph}}{1 + i\omega\tau_{ph}}. \tag{10}$$

When Eq. (10) 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. 1 for three types of
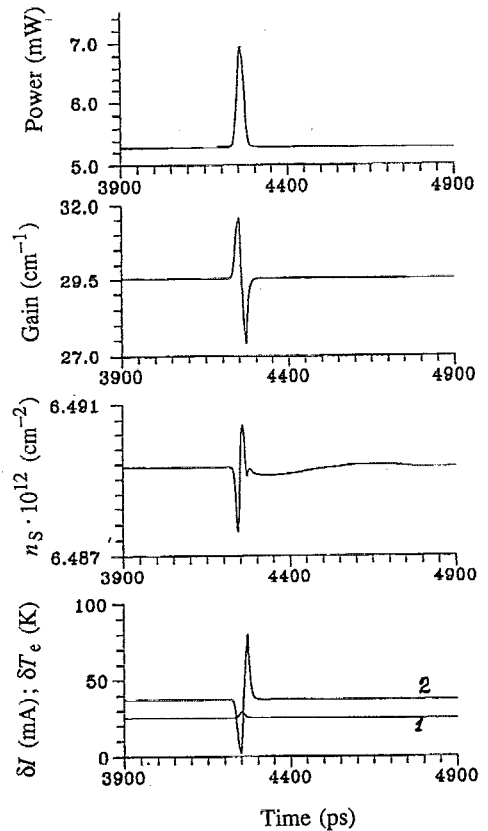

FIG. 2. Response in the output power, the gain, and the sheet carrier concentration n_s in the active region to a dual pulse of the pumping current (curve 1) and the carrier temperature (curve 2). Assumed laser parameters are as in Fig. 1.

modulation: (1) purely by the pumping current, (2) purely by the electron temperature, and (3) by their coherent combination as in Eq. (10).

Elimination of the resonance also enables the generation of ultrashort (≈ 50 ps) optical pulses, Fig. 2. Moreover, the suppression of relaxation oscillations of the carrier density makes possible a high repetition rate coding of information with short pulses. We have calculated the laser response to a 10 Gb/s series of dual current-temperature pulses, as in Fig. 2, 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. (10). In this context, one should take into account the delay τ_ϵ, by which T_e lags an external heating pulse H, cf. Eq. (3). We remark that in InGaAs for $n \gtrsim 10^{18}$ cm^{-3} the energy relaxation time can be rather long $\tau_e \gtrsim 5$ ps due to phonon bottleneck effects.[6]

Another interesting possibility offered by the dual modulation method, consists in suppression of the wavelength chirping at high modulation frequencies. Recall that this unwelcome phenomenon[4] originates from the relaxation oscillations, which lead to variations $\delta\eta$ in the real part η of the refractive index $\eta_c = \eta + i\kappa$ in the active re-

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gion. Variations of n affect η in two ways: by changing the free-carrier absorption and by changing the optical gain, Eq. (4). In InGaAs lasers both effects give similar contributions to $\delta\eta$, shifting the lasing mode wavelength by as much as several angstroms. The small-signal index change $\hat{\eta}_{fc} = -A\hat{n}$, where for InGaAs[7] $A \approx 10^{-20}$ cm^3. The contribution $\hat{\eta}_g$ from variations in the interband optical gain can be evaluated by the Kramers–Kronig relation

$$\hat{\eta}_g = -\frac{\bar{\eta}}{\pi} \mathscr{P} \int_0^\infty \frac{\hat{g}(\Omega')d\Omega'}{\Omega'^2 - \Omega^2}, \tag{11}$$

where $\mathscr{P}$ denotes the integral principal value. Keeping in mind that in our present model the gain variation $\hat{g}$ has two contributions, Eq. (4), we should distinguish the corresponding two contributions, $\hat{\eta}_{g_n}$ and $\hat{\eta}_{g_T}$ in the refractive index variation, each given by Eq. (11) with $\hat{g}$ replaced by $g'_n \hat{n}$ and $g'_T \hat{T}_e$, respectively. We now see that it is possible to target a complete suppression of the total index oscillation,

$$\hat{\eta} = \hat{\eta}_{fc} + \hat{\eta}_{g_n} + \hat{\eta}_{g_T} = 0, \tag{12}$$

by judiciously choosing a relationship between $\hat{T}_e$ and $\hat{J}$ in such a way that

$$\gamma = \frac{\bar{\eta}\mathscr{P} \int g'_T(\Omega')(\Omega'^2 - \Omega^2)^{-1} d\Omega'}{\pi A - \bar{\eta}\mathscr{P} \int g'_n(\Omega')(\Omega'^2 - \Omega^2)^{-1} d\Omega'}. \tag{13}$$

It should be emphasized that having targeted the complete elimination of chirp, we pay a penalty in the modulation frequency, since for a non-vanishing γ the electron-photon resonance is not 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 δT_e is usually weaker than that due to δn. If we compare $\hat{\eta}_{g_T}$ at the point of complete elimination of relaxation oscillations ($\gamma = 0$) with the amount of chirp $\hat{\eta}_n = \hat{\eta}_{fc} + \hat{\eta}_{g_n}$ 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 $g'_n(\Omega)$.

We have carried out calculations based on Eq. (11) for an InGaAs quantum-well laser with parameters as in Fig. 1. An optical power pulse of 1 mW corresponds to the variation of either $\delta T_e \approx 80$ K (300 K → 380 K) or $\delta n \approx 10^{11}$ cm^{-2}. These variations produce an approximately equal but opposite in sign contributions to $\delta\eta$, viz. $\delta\eta_{g_T}$ and $\delta\eta_n$, which compensate each other. For a much higher pulse power, the amplitude of δn is also much larger and the required dual pulse δT_e necessary to achieve the compensation $\delta\eta_{g_T} + \delta\eta_n \approx 0$ becomes too large. In this situation, it may be more convenient to maintain the regime $\delta n = 0$, Eq. (10), for which $\delta\eta = \delta\eta_{g_T}$ independent of the optical pulse power. For the situation described in Fig. 2, where $\hat{T}_e = 40$ K (300 K $< T_e <$ 380 K), the variation $\hat{\eta}_{g_T} \approx 0.001$ and corresponds to a wavelength chirping of about 0.1 nm.

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