17.4 Hot electrons in lasers

Carrier heating in double heterostructure lasers arises mainly due to injection of electrons and holes into the narrow-gap active region "over the cliff" from the wide-gap cladding layers, cf. Fig. 16. Another contribution to heating, which becomes appreciable far above the threshold, is free carrier absorption of the coherent radiation in the laser cavity. The energy balance in the active region is described by Eq. (6a) with the average power input per carrier approximately given by

$$\mathcal{P} = \frac{\alpha S\, \hbar\Omega}{2} + \frac{J\,\Delta}{2n} \quad , \tag{21}$$

where $J \equiv I/eVA$ is the electron flux per unit volume VA of the active layer, I the pumping current, S the photon density in the active layer, Ω the optical frequency, Δ the kinetic energy per carrier injected into the active region, and $\alpha > 0$ the free carrier absorption rate. It is assumed in Eq. (21) that effectively only one type of carriers absorbs radiation and that the densities of electrons and holes are the same in the active layer, $n = p$, hence the factor of 2.

For the free-carrier absorption coefficient Henry et al. (1983) cite $\alpha_0 = 25\,\mathrm{cm}^{-1}$ at $p = 10^{18}\,\mathrm{cm}^{-3}$ and $\lambda = 1.6\,\mu\mathrm{m}$. Our α above is related to α_0 by $\alpha_0 = \alpha p\, \bar{c}$, where $\bar{c} \equiv c/\kappa_g$ is the group velocity of light. The energy Δ depends on the laser structure, e.g., the fraction of carriers that get into the active region by tunneling; in general we can say that $\Delta \lesssim \Delta E_G$, where ΔE_G is the bandgap difference between the cladding and the active layers. For InP cladding ($E_G = 1.35\,\mathrm{eV}$) and $\lambda = 1.55\,\mu\mathrm{m}$ active layer, $\Delta E_G = 0.55\,\mathrm{eV}$.

Taking the average electron energy in the approximate nondegenerate form (4) with $\langle \mathbf{v} \rangle = 0$, we can estimate the carrier overheating near the threshold ($S \approx 0$, $J \approx n/\tau_{sp}$):

$$k\,(T_e - T) = \frac{\tau_\varepsilon}{\tau_{sp}}\,\frac{\Delta}{3} \quad , \tag{22}$$

where τ_{sp} is the lifetime of carriers (including nonradiative processes). With $\tau_{sp} \approx 0.5\,\mathrm{ns}$ and $\tau_\varepsilon \approx 1\,\mathrm{ps}$ (Hall et al. 1992) we find from Eq. (22) an overheating by $(T_e - T) \approx 4\,\mathrm{K}$.

It should be noted that experimental estimates for τ_ε vary widely, because in a high-density bipolar plasma one can expect a strong dependence of τ_ε on the carrier concentration and temperature. Indeed, the dominant energy-loss mechanism in long-wavelength semiconductors is due to the emission of polar optical phonons and typically leads to $\tau_\varepsilon < 1\,\text{ps}$. However, at *high carrier concentrations* the hot-carrier cooling rates may be much slower, especially in quantum wells, as has been found by a number of workers who studied luminescence of photo-excited carriers on a picosecond scale (Westland et al. 1988; Lobetanzer et al. 1989). Similar results were found by Belenky et al. (1994b) by measuring thermionic emission of heated carriers from a quantum well. The anomalously long relaxation times τ_ε, substantially exceeding the characteristic time of optical-phonon emission, can be explained by hot phonon effects (Kumekov and Perel, 1988). Because of the small group velocity, optical phonons do not move appreciably during their lifetime and are likely to be re-absorbed by the free carriers, provided the carrier plasma has a high enough concentration (typically, higher than $10^{18}\,\text{cm}^{-3}$). In this situation, termed the *optical phonon bottleneck*, the optical phonon gas and the carriers are characterized by the same effective temperature T_e and the energy relaxation of the coupled carrier-phonon system is determined by the process of polar optical phonon decay into acoustic phonons. The above estimate of $4\,\text{K}$ overheating at room temperature near the laser threshold is rather modest compared to calculations that include the bottleneck effect (see, e.g., Tsai et al. 1993). Nevertheless, even this small amount of carrier overheating may have important consequences for applications.

A primary example of such applications is the subcarrier-multiplexed optical transmission system used for cable TV distribution. These systems require source lasers with an extremely *linear* functional relationship between the input current $I(t)$ and output light intensity $L(t)$. As functions of time, these quantities encode all the information carried by multichannel cable TV networks. Even a small amount of nonlinearity brings about some intermodulation distortion (e.g., TV channels operating at 180 and 240 MHz produce unwelcome parasitic signals around

60 and 420 MHz). Channel capacity in current systems is about 40 channels, but future cable networks are expected to carry much larger number of channels. Linearity specifications for an 80-100 channel system become extremely stringent.

The unwelcome consequences of hot-carrier effects arise from the dependence of the optical gain g on the carrier temperature. This dependence is typically of the form

$$g\,(T_{\rm e}, n, \Omega) = g_0\,(f_{\rm e} + f_{\rm h} - 1)\ , \tag{23}$$

where $f_{\rm e}$ and $f_{\rm h}$ are the Fermi functions of electrons and holes, respectively, at energies selected by the incident photons $\hbar\Omega$ and g_0 is a constant. Dynamics of the laser can be described by the following system of rate equations (Gorfinkel and Luryi, 1995):

$$\frac{{\rm d}n}{{\rm d}t} = J - g\,S - \frac{n}{\tau_{\rm sp}}\ ; \tag{24a}$$

$$\frac{{\rm d}S}{{\rm d}t} = \Gamma S\,(g - \alpha n) - \frac{S}{\tau_{\rm ph}}\ , \tag{24b}$$

$$\frac{{\rm d}T_{\rm e}}{{\rm d}t} = \frac{2}{3}\mathcal{P} - \frac{T_{\rm e} - T}{\tau_{\varepsilon}}\ , \tag{25}$$

where $\tau_{\rm ph}$ is the lifetime of photons in the laser cavity and $\mathcal{P}$ is the power input given by Eq. (21). Nonlinear effects originate from the terms $g \cdot S$ and $S \cdot n$ in Eqs. (24).

If the carrier temperature were constant, the gain in a given optical mode of the laser would depend only on the concentration n. Since above the threshold the value of n is "pinned" in a static regime by the condition of stable generation, g = loss, the dependence of the output light power ($\propto S$) on the pumping current J is linear to a good approximation. Small nonlinear effects are possible at high modulation frequencies, when the carrier concentration deviates from the pinned value. These effects contribute the so-called "intrinsic" distortion (Darcie et al. 1985). The latter goes as $(f/f_{\rm r})^2$, where $(2\pi f_{\rm r})^2 = S_0\, g_n'/\tau_{\rm ph}$ is the resonant frequency of the laser operated around a steady-state output level S_0 (cf. Eq. 32 below). The intrinsic distortion decreases with S_0 and becomes small at sufficiently high bias currents. However, increasing the steady-state bias introduces additional nonlinearities, associated with carrier heating.

Because of the T_e variation above threshold, carrier concentration is no longer "pinned" even in the static regime. Since normally $g_T' < 0$ and $g_n' > 0$, the increasing T_e with increased pumping current is accompanied by an increase of carrier concentration n in the active region. An important adverse consequence of the T_e variation is a current-dependent carrier leakage out of the active region in laser heterostructure. This leakage contributes an unwelcome intermodulation distortion. It can be minimized in properly designed double-heterostructure lasers (Belenky et al. 1995).

Recently, Gorfinkel and Luryi (1995) considered nonlinear distortions in systems arising from carrier heating effects which cannot be avoided in any laser structure, since they are brought about by the very existence of the modulated optical and electrical signals. These effects are described by Eq. (21) − namely, the free-carrier absorption of the coherent radiation in the cavity and the power flux into the active layer, associated with the input current. It turns out that the resultant intermodulation distortion approaches the tolerance limit for an 80-channel cable TV system − even if all other "parasitic" nonlinearities are eliminated, including the above mentioned carrier leakage. Thus, hot carrier effects present a fundamental limitation for the number of channels available on cable television.

Another fundamental nonlinearity in lasers is associated with carrier heating by the power released in Auger recombination processes. In these processes the potential energy of an electron-hole pair is transferred to free carriers. Gorfinkel and Luryi (1995) showed that even though this power is typically comparable to that due to free carrier absorption, its contribution to the intermodulation distortion is negligible, being of higher order in the optical modulation depth. The rate of Auger recombination goes as $C_A n^3$, see e.g., Agrawal and Dutta (1993), and hence above threshold it is independent of the signal level.

Carrier heating by Auger recombination processes may play a very important role in long wavelength semiconductors, especially in single QW lasers where, due to a small value of the

optical confinement factor Γ the carrier concentration at the lasing threshold reaches rather high values, typically much exceeding $10^{12}\,\mathrm{cm}^{-2}$. At high injection current I, this may lead to a substantial carrier heating. Inclusion of this effect in the energy balance equation is accomplished by adding a term $\sim C_A n^2 E_G$ into Eq. (21). The increasing carrier temperature $T_e(I)$ suppresses the optical gain g and may lead to the appearance of a maximum g_{max} in the dependence $g(I)$ of gain on the pumping current.[9] If the total losses in the laser cavity exceed g_{max} then the structure will not lase at any pumping current. Since the Auger coefficient C_A itself increases with the temperature, one can expect a sharp disappearance of lasing at some critical temperature. Similar nonmonotonic $g(I)$ dependence with a maximum, which results from electron heating and is accompanied by a sharp disappearance of lasing with increasing temperature, is important in the operation of the new *unipolar* laser and will be discussed in the next section. Note that for a constant T_e the dependence $g(I)$ is always monotonic. If the losses do not exceed g_{max}, then the laser generation regime can be reached, but the negative slope of $g(I)$ characteristic results in peculiar instabilities for currents exceeding $I_{cr} \equiv I(g_{max})$. For the same value of g one can have two regimes that differ in n and T_e and, most importantly, in the output radiation power. The high-T_e regime corresponds to higher n and lower $\mathcal{P}$. This regime is metastable. Illumination of the laser by a sufficiently powerful pulse of external light temporarily suppresses the Auger recombination and switches the laser into the stable regime with a high $\mathcal{P}$.

17.4.1 Quantum cascade lasers. The QCL is a new mid-infrared laser, based on unipolar transitions of electrons between energy levels created by quantum confinement. Since the first demonstration of QCL (Faist et al. 1994), its design has continuously improved through a series of elegant innovations by the Bell Labs group (Faist et al. 1995a,b) culminating in their

9. This discussion follows an unpublished work by V. B. Gorfinkel and S. Luryi (1992).

recent report of a high power room-temperature operation (Faist et al. 1996).

Kinetics of the carrier transport in the QCL and the device structure are schematically illustrated in Fig. 17. The device has multiple periods (25 in reported structures) and each electron performs a cascade of transitions down the periodic ladder. Under the operating conditions, each period must have net zero charge: by Gauss' law this is necessary in order that subsequent periods could replicate each others electrostatic state. To avoid the space charge accumulation, associated with current flow, one therefore needs a reservoir of positive fixed charge, compensating the negative mobile charge in each period. Introduction of such a reservoir, implemented as a doped superlattice region is the key design innovation that led to the successful implementation of a unipolar laser. Subsequent QCL designs entrusted the reservoir with an additional mission to suppress the unwelcome tunneling from the upper level into the continuum. For this purpose, the superlattice is implemented as an electronic Bragg filter with a stop band in the range of energies of upper level states (Faist et al. 1995). Thus, most carriers enter the QW only via the upper and leave only via the lower subband, while leakage from the upper subband directly into the reservoir is negligible.

Theoretical analysis of the QCL operation (Gorfinkel et al. 1996) suggests that it is dominated hot-electron effects. Non-equilibrium electron distributions arise from the power $\mathcal{P} \approx J \cdot \hbar\Omega$ per unit area, dissipated in each cascade period. The energy stored in the *transverse degrees of freedom*, corresponding to in-plane motion of carriers, fundamentally changes the lineshape of intersubband resonance and the spectral characteristics of gain. Depending on the laser design, there may be different scenarios of how this energy is distributed among various electrons and dissipated into the lattice. The simplest (but not the most advantageous) scenario arises at high carrier concentrations, when the electron-electron interaction is sufficiently fast to equilibrate the input power among all free carriers in a QCL period. In this case, the kinetic energy of two-dimensional electrons can be characterized by an effective temperature T_e, which

is moreover the same in both subbands. This temperature can be found from the balance equation,

$$\frac{k\,\mathrm{d}T_\mathrm{e}}{\mathrm{d}t} = \mathcal{P} - \frac{n_\mathrm{D} k\,(T_\mathrm{e} - T)}{\tau_\varepsilon} \quad , \tag{26}$$

where $n_\mathrm{D} \sim 10^{11}\ \mathrm{cm}^{-2}$ is the doping concentration per period. In the steady state the overheating is directly proportional to the current: $k\,(T_\mathrm{e} - T) = J\,\tau_\varepsilon \hbar\Omega / n_\mathrm{D}$.

The subband concentrations in the QCL are found from the following rate equations:

$$\frac{\partial n_2}{\partial t} = J - \frac{n_2}{\tau_{21}} - \bar{g}S \quad , \tag{27a}$$

$$\frac{\partial n_1}{\partial t} = \frac{n_2}{\tau_{21}} + \bar{g}S - \frac{n_1}{\tau_{1\,\mathrm{out}}} \quad , \tag{27b}$$

where J is the current flux and S is the photon density per unit area in the lasing mode (frequency Ω_L). The gain $\bar{g} = c\,\bar{g}\,(\Omega_\mathrm{L})$ is the intersubband gain (in units of sec^{-1}) at the lasing frequency, where $\bar{c} = c/\sqrt{\kappa_\infty}$ is the speed of light in the lasing mode. The time constants τ_{21} and $\tau_{1\,\mathrm{out}}$ describe, respectively, the rate of nonradiative intersubband transitions and carrier removal from the bottom subband. The ratio of these constants, $\xi_0 \equiv \tau_{1\,\mathrm{out}}/\tau_{21}$, is an important design parameter of an intersubband laser, as it determines the carrier density ratio in a steady state below threshold, $n_1 = \xi_0 n_2$.

To illustrate the importance of hot-electron effects, consider an oversimplified model of the QCL which neglects the *nonparabolicity* in the conduction band. In this case one can use the well-known expression for the gain of a two-level system (Yariv, 1991, Chap. 5)

$$g \equiv \frac{4\pi e^2 |\, z_{12}\,|^2 \Omega}{\hbar c \sqrt{\kappa_\infty}\,\gamma_0}\ \frac{n_2 - n_1}{a} \quad , \tag{28}$$

where z_{12} is the transition matrix element, $\sqrt{\kappa_\infty}$ the refractive index, a the quantum well width and $\hbar\gamma_0$ is the spontaneous linewidth. For high T_e the dominant contribution to the linewidth results from the emission of optic phonons by electrons making transitions within the same subband. One can therefore assume that γ_0 depends on the effective temperature of the electron

ensemble.

From the generation condition $g = \alpha_{\text{Loss}}$ we find that the threshold is described by the following equation for T_e:

$$\frac{\hbar\gamma_0(T_e)}{k\,(T - T_e)} = \mathcal{R} \ , \tag{29}$$

where $\mathcal{R}$ is a dimensionless constant,

$$\mathcal{R} = \frac{4\pi e^2\,|\,z_{12}\,|^2\,\Gamma\, n_D}{\hbar c\,\sqrt{\kappa_\infty}\,a\,\alpha_{\text{Loss}}}\,\frac{\tau_{21} - \tau_{1\,\text{out}}}{\tau_\varepsilon} \ . \tag{30}$$

Consider a graphical solution to Eq. (29), Fig. 18, assuming that the intersubband linewidth γ_0 is an increasing and *concave* function of the effective carrier temperature.[10] The threshold carrier temperature is determined by the first intersection of a line $(T_e - T)\cdot\mathcal{R}$ with the concave upward curve $\hbar\gamma_0(T_e)$. Depending on the value of $\mathcal{R}$ and the ambient temperature, equation (29) may have no solutions at all or − if the value of $\mathcal{R}$ is sufficiently large − *two* solutions. For a given $\mathcal{R}$ the highest lattice temperature $T = T_{\max}$ at which threshold is possible, is determined by the condition that the line is *tangent* to the curve. Finally, there is a minimum value of $\mathcal{R}$ such that $T_{\max} > 0$. If $\mathcal{R}$ is less than this minimum value, then there can be no lasing at any temperature.

From this qualitative example we conclude that both the transverse relaxation of the intersubband resonance and the hot-carrier effects are indispensable for the correct physical understanding of the temperature behavior of quantum cascade lasers. Quantitatively, the two-level model (28) does not adequately describe the intersubband gain function in experimentally realized (Faist et al. 1994-1996) QCL heterosystems. The reason is two-fold. Firstly, one cannot neglect the nonparabolicity: its inclusion suppresses the peak gain value by more than

10. Intersubband photoluminescence experiments indicate that $\gamma_0(T)$ is an increasing concave function of the *lattice* temperature (J. Faist, private communication).

an order of magnitude. Secondly, the relaxation rate of intersubband resonance, which is mainly due to transverse intrasubband scattering processes, is strongly dependent on electron kinetic energy, $\gamma = \gamma(E)$. Indeed, the rate of optical phonon scattering has a threshold nature: when $E > \hbar\omega_{op}$ it increases by nearly an order of magnitude due to the onset of intrasubband emission. The generalized expression for the intersubband gain at an optical frequency Ω is of the form (Gelmont et al. 1996):

$$g(\Omega) = \frac{4e^2 |z_{12}|^2 m_2 \Omega}{\hbar^3 a\, c\sqrt{\kappa_\infty}} \int_0^\infty \frac{d\varepsilon\, \gamma(\varepsilon)\,[f_2(\varepsilon) - f_1(\varepsilon_1)]}{[\Omega - \Omega_\varepsilon]^2 + [\gamma(\varepsilon)]^2} \quad , \tag{31}$$

where Ω_ε is the optical transition frequency for the in-plane electron momentum $\hbar k = \sqrt{2m_2\varepsilon}$, viz. $\hbar\Omega_\varepsilon \equiv \hbar\Omega_0 + \varepsilon_2 - \varepsilon_1$, where $\varepsilon_2 \equiv \varepsilon$ and $\varepsilon_1 = \hbar^2 k^2 / 2m_1$ are kinetic energies in the upper and lower subbands, respectively, characterized by the effective masses m_2 and m_1 and the distribution functions f_2 and f_1. The function $\gamma(\varepsilon)$ describes the transverse phase relaxation rate due to intrasubband scattering, both elastic and inelastic.

It is worthwhile to stress that Eq. (31) does not rely on the electron temperature approximation and remains valid for general electron energy distributions within each subband. The most attractive situation in fact arises at low carrier concentrations ($n_D \ll 10^{11}$ cm^{-2}) when electron-electron collisions are not fast enough to establish a common T_e between the two subbands. It turns out (Gorfinkel et al. 1996) that in this regime the lower-subband electron distribution can be approximately characterized by a "negative" temperature, so that states near the subband bottom are mostly unoccupied.[11] In this case it is possible to achieve positive values of the gain at room temperature even in the absence of an overall population inversion between the two subbands.

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11. For high positive temperatures the electron distribution becomes uniform over the entire subband. For $T_e < 0$ the distribution is inverted, the lower energy states having smaller occupation probability; one can say that negative temperatures are higher than $T_e = \infty$,

For higher (but still moderate, $n_D \lesssim 10^{11}$ cm^{-2}) concentrations, the T_e approximation becomes applicable, but it is still possible (at least semi-quantitatively) to regard optical phonon scattering as the dominant phase breaking mechanism. In this range, Gorfinkel et al. (1996) predict a peculiar hot electron effect manifested in the dependence of the lasing wavelength λ on the pump current, including a regime where λ switches "digitally" between two stable values. Still higher concentrations require a special consideration, as it becomes necessary to explicitly include electron-electron scattering in the calculation of $\gamma(\varepsilon)$. Also the energy balance equation in this range should include the dependence of the electron cooling rate on the carrier density and effective temperature due to optical phonon bottleneck. Validity of Eq. (31) itself is not restricted to any particular scattering mechanism responsible for the transverse phase relaxation.

17.4.2 Hot-electron lasers. Due to the dependence $g(T_e)$, Eq. (23), variations in the carrier temperature can produce rapid changes in the optical output of a semiconductor laser (Rivlin, 1985). Possible device applications of this effect stem from the observation (Gorfinkel and Filatov, 1990) that controlling the laser by an independently modulated T_e offers a higher modulation bandwidth than the conventinal pump current control. In general, the modulation response of a semiconductor laser $\mathcal{R}_J(\omega) \equiv \delta S/\delta J$ (at a constant T_e) is limited in bandwidth by an intrinsic resonance in the nonlinear laser system (the electron phonon resonance). Due to this effect, the response $\mathcal{R}_J$ reaches a maximum near the resonance frequency ω_r where

$$\omega_r^2 = \frac{S_0 g_n'}{\tau_{ph}} \quad , \tag{32}$$

and then decreases by 20 dB/dec, i.e. as $1/\omega^2$. The analogous response function for T_e modulation $\mathcal{R}_T(\omega) \equiv \delta S/\delta T_e$ (at a constant J) rolls off at high frequencies only as $1/\omega$, i.e. by 10 dB/decade (Gorfinkel et al. 1991). This result can be demonstrated by a small-signal analysis of rate equation (24) taking both J and $\mathcal{P}$ as externally varied parameters.

Rapid electron heating can be accomplished in a variety of ways. The original suggestion by Gorfinkel et al. (1991) was to employ a four terminal laser structure with a carrier-heating current driven laterally through the active layer. High-frequency modulation of T_e by several degrees was demonstrated in such a structure (Bagaeva et al. 1991). Other proposals include power delivery to the electronic system by intersubband absorption in quantum wells (Noda et al. 1990; Gorfinkel and Luryi, 1992) and by a vertical hot-electron injection in a transistor-like heterostructure (Tolstikhin and Mastrapasqua, 1995).

Modulation of the laser output power via carrier heating at a constant pumping has the potential of offering higher bandwidth of modulation. Still higher potential for applications lies in the possibility of a *synchronous* control of both the T_e and the pumping current (Gorfinkel and Luryi, 1993, 1994). This technique, called the dual modulation, permits in principle a complete elimination of electron-photon oscillations and linearization of the light-current characteristics in a wide frequency band. Indeed, with an independent control of gain via T_e one can impose the condition of constant n even as both J and S vary. Setting $dn/dt = 0$ in Eqs. (24). we find

$$0 = J - J_{th} - gS \quad ; \tag{33a}$$

$$\frac{dS}{dt} = \Gamma(J - J_{th}) - \tau_{ph}^{-1} S \quad , \tag{33b}$$

where $J_{th} \equiv n_{th}^2/\tau_{sp}$ and n_{th} is the pinned carrier concentration above threshold. For the sake of simplicity, we have neglected the free carrier absorption term. Equation (33b) establishes a *linear* relationship between the functions $S(t)$ and $J(t)$:

$$S = = \Gamma e^{-t/\tau_{ph}} \int_0^t [J(t') - J_{th}]\, e^{t'/\tau_{ph}}\, dt' \quad . \tag{34}$$

For this property to hold we must satisfy Eq. (33a) with the help of a simultaneous variation of T_e, viz.

$$g\,[T_e(t),\, n_{th}] = \frac{J(t) - J_{th}}{S(t)} \quad . \tag{35}$$

Having solved Eq. (35) for the time dependence $T_e(t)$, we can determine the required heating power signal from Eq. (25). Linearity of the $S\,[J(t)]$ relationship is of great value for optical communication systems.

Small signal analysis of dual modulation can be done by linearizing Eqs. (24):

$$J(t) = J_0 + \delta J\ e^{i\omega t} \quad ; \tag{36a}$$
$$g(t) = g_0 + \delta g\ e^{i\omega t} \quad , \tag{36b}$$

where $\delta g = g_n' \,\delta n + g_T' \,\delta T_e$. Under the condition $dn/dt = 0$ we find the following response function:

$$\delta S = \frac{\delta J}{g_0 (1 + i\omega\tau_{ph})} \quad , \tag{37a}$$

while the required relation between the dual inputs δg and δJ is of the form

$$\delta g \equiv g_T' \,\delta T_e = \frac{i\omega\tau_{ph}}{1 + i\omega\tau_{ph}} \frac{\delta J}{S_0} \quad . \tag{37b}$$

We remark that the "target condition" (37b) is practically frequency-independent for $\omega\tau_{ph} > 1$ and that the magnitude of the required dual modulation input is inversely proportional to S_0. 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. 19 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. (37b).

Suppression of relaxation oscillations of the carrier density makes possible a high repetition rate coding of information with short pulses. Calculated laser response to a 10 Gb/s series of dual current-temperature pulses (Gorfinkel and Luryi, 1993) 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 $\delta T_e(t)$. The dual

modulation method is not limited to the variation of T_e. Any of the parameters of the system of rate equations (24), such as the modal gain and the photon lifetime, may be varied externally.

Control of the semiconductor injection laser by modulating T_e is not the only way of utilizing hot electrons in lasers. In fact, there is a far more advanced laser application of hot electrons. *Injectionless* far infrared lasers based on *hot hole* effects in germanium have been demonstrated several years ago.[12] In these devices population, the inversion of hot holes arises due to the very different dynamical behavior of heavy and light holes in crossed electrical and magnetic fields. Discussion of these effects would be out of scope of this chapter.[13]

13. For an introductory discussion, the reader is referred to an early review by Andronov (1987) as well as to several papers in the special issue on far-infrared semiconductor lasers, *Opt. Quantum Electron.* **23**, No. 2 (1991).