Wavelength Chirp and Dependence of Carrier Temperature on Current in MQW InGaAsP–InP Lasers

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***Abstract*—In this paper, we derive a relation between the wavelength chirp and carrier temperature in semiconductor lasers. The coefficient relating the change in carrier temperature and chirp is expressed in terms of the temperature derivative of the optical gain, and two parameters describing the variation of refractive index produced by the variation of optical gain due to change of carrier quasi-Fermi level separation or carrier temperature. We have measured these parameters for MQW InGaAsP lasers. Using this data, we estimated the rate of the temperature increase with current above threshold in these devices, which is 0.13 K/mA.**

I. INTRODUCTION

UNDER operation conditions of 1.3-μm InGaAsP–InP lasers, the injection current density is about 15 kA/cm^2. Under these conditions, one can expect a significant difference between the lattice temperature T_L and the temperature of the electron-hole-plasma in the active region $T_{e\text{-}h}$. This difference is determined by the power acquired by carriers in the active layer and their energy relaxation time. Both quantities have been disputed in the literature (see [1] and references therein).

Since the modal optical gain, wavelength chirp, and carrier leakage over the heterobarrier are sensitive to the carrier temperature $T_{e\text{-}h}$, the study of carrier heating is important for understanding the laser physics and improving device design. Experimental studies of the heating in heterostructures using optical excitation and time-resolved spectroscopy were carried out by different groups (see, for example [2]). In semiconductor lasers and light-emitting diodes (LED's), the carrier temperature was usually estimated from the high-energy tail of the spontaneous emission spectra [3]–[6]. Carrier temperatures up to 400 K at room temperature were reported in InGaAsP LED's [3]. In [4], the authors studied InGaAsP–InP buried heterostructure lasers with a bulk active layer. They did not observe carrier heating at room temperature for current densities up to 10 kA/cm^2 within an experimental accuracy of 10 K. In [5], the authors studied photoexcieted InGaAsP–InP broad-area quantum-well (QW) structures and found considerable heating at low temperatures. In AlInAs–GaInAs–InP bulk double heterostructure light-emitting transistors significant carrier heating at cryogenic temperatures was determined from the spontaneous emission profile [6] as well as using an analysis of thermionic emission [7].

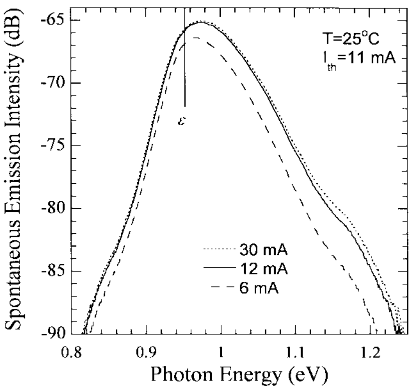

Fig. 1. TSE spectra of DFB laser recorded from the side of the laser chip at different currents: below threshold (6 mA—dashed line), slightly above threshold (12 mA—solid line), and above threshold (30 mA—dotted line). The lasing energy ϵ is also indicated.

The problem with the estimation of carrier temperature from the high energy tail of the spontaneous emission arises from the fact that the density of states is not a well-known function of energy. This makes it difficult to analyze the spontaneous emission, especially in the case of complicated laser structures with high injection levels at room temperature (see, for example, Fig. 1). Study of the thermionic emission requires a special sample geometry and is generally unsuitable for real laser structures.

In this paper, we propose a new experimental technique which allows us to measure the rate of change of the carrier plasma temperature with pumping current above threshold. The new method is based on the relation between carrier heating in the active layer and wavelength chirp. The paper is organized in the following way. In Section II, we derive the formula relating the rate of carrier heating to the wavelength chirp. We also define a parameter α_μ which relates the change of the real and imaginary parts of the refractive index due to the change of carrier temperature at constant quasi-Fermi level separation. This parameter is similar to the linewidth enhancement factor α_T, which is the ratio of the change of the real and imaginary parts of the refractive index due to the change of carrier concentration (or quasi-

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Fermi level separation between electrons and holes) at constant temperature. In Section III, we describe a method to measure α_T and α_μ, which along with the temperature derivative of the modal optical gain determines the relation between the carrier heating and wavelength chirping. In Section IV, we describe the experimental procedures. The results are summarized and discussed in Section V.

II. ANALYSIS OF THRESHOLD CONDITION AND WAVELENGTH CHIRP

Wavelength chirp results from the dependence of the real part of refractive index on current. The real and imaginary parts of refractive index are related through the Kramers–Krönig relation. Under lasing condition the distortion of spectral profile of the optical gain (and therefore the imaginary part of refractive index) is caused by spectral hole burning [8]–[10] and carrier heating [11]. At the lasing wavelength, this change of the optical gain is compensated by the increase of the carrier concentration; however, pinning of the optical gain does not occur at all wavelengths. This results in dependence of refractive index on current above threshold and in the wavelength chirp.

In order to characterize the spectral hole burning, Linke and Koch [10] introduced a phenomenological gain compression parameter $\in$. They described the deviation of the lasing frequency $\Delta\nu$ in terms of this parameter:

$$\Delta\nu \approx -\frac{\alpha}{4\pi}\left[\frac{d}{dt}\ln P(t) + \frac{2\Gamma \in}{V_{\text{act}}\eta h\nu} P(t)\right] \tag{1}$$

where α is the linewidth enhancement factor [12], $P(t)$ the output optical power, Γ the optical confinement factor, V_{act} the active layer volume, η the quantum efficiency, and h the Plank constant. The first term characterizes the chirp caused by the dynamics of electron-photon resonance. The second term describes the adiabatic chirp, which we addressed in this work. We can rewrite (1) for adiabatic chirp

$$\frac{d\nu}{dI} \approx -\frac{\alpha}{4\pi}\frac{\Gamma \in}{V_{\text{act}} q} \tag{2}$$

where q is the electron charge. In [10], the authors attributed $\in$ to the spectral hole burning, which produces a small symmetric dip in the gain about the laser line. However, the experimental studies of fast dynamics of the optical gain in semiconductor laser structures have shown that the electron-hole plasma thermalization time is much shorter than the time needed for thermal equilibrium between the plasma and crystal lattice [13]. Therefore, in our treatment of the wavelength chirp, we consider the carrier heating to be dominant effect and neglect the spectral hole burning effect. In this case, the energy distribution of electrons and holes is described by Fermi functions with the same temperature $T_{e\text{-}h}$. This temperature may differ from the lattice temperature, also electrons and holes have different quasi-Fermi level energies when voltage is applied to a laser diode.

The threshold condition for a semiconductor laser is

$$g(n, T) = \Gamma G(n, T) - a_{\text{tot}}(\lambda) = 0 \tag{3}$$

where g is the net modal optical gain in the cavity, G the material gain in the active layer, and a_{tot} the total optical loss, which includes internal loss and mirror loss. Optical gain is a function of the carrier concentration in the active layer n and carrier temperature T (subscript e-h is dropped for simplicity). An alternative way is to express the optical gain as a function of the carrier temperature and quasi-Fermi level separation between electrons and holes μ. This is possible because only two variables out of three: n, μ, and T are independent in case of Fermi distribution. We prefer using μ instead of n because it is easier to measure. The value of the total optical loss is only weakly dependent on either μ or T [14]. On the other hand, $a_{\text{tot}}(\lambda)$ is a strong function of wavelength in DFB lasers. If the laser is biased above threshold, any variation of one of the above parameters results in a change of one or two others so that the threshold condition of (3) remains valid. Thus, we can write in differential form, using variables μ and T

$$\left.\frac{\partial G}{\partial \mu}\right|_T d\mu + \left.\frac{\partial G}{\partial T}\right|_\mu dT = \frac{1}{\Gamma}\left.\frac{\partial a_{\text{tot}}}{\partial \lambda}\right|_{T,\mu} d\lambda. \tag{4}$$

In this paper, we study DFB lasers that have HR- and AR-coated mirror facets. The position of the lasing mode in the stop-band is defined by the random cleave position of the HR-coated facet with respect to the grating. Lasers with different position of the laser mode in the stopband have different $\partial a_{\text{tot}}/\partial\lambda$. To simplify the problem, we consider and study experimentally a laser operating at the middle of the stopband. Such a device can be regarded as a folded quarter-wave-shifted laser. In this case, the DFB is maximum and so the optical loss is at its minimum, and therefore $\partial a_{\text{tot}}/\partial\lambda = 0$. We express the modal optical gain in the laser cavity through the imaginary part of the refractive index using

$$N'' = -\frac{G}{\varepsilon}\frac{hc}{4\pi}$$

and

$$\delta N'' = -\frac{1}{\Gamma}\frac{\delta g}{4\pi}\lambda \tag{5}$$

where ε is photon energy, and c the light velocity. Then (4) becomes

$$\left.\frac{\partial N''}{\partial \mu}\right|_T d\mu + \left.\frac{\partial N''}{\partial T}\right|_\mu dT = 0. \tag{6}$$

We can also express the full change of the real part of the index N' due to the same small variations $d\mu$ and dT, as in (6)

$$\left.\frac{\partial N'}{\partial \mu}\right|_T d\mu + \left.\frac{\partial N'}{\partial T}\right|_\mu dT = dN'. \tag{7}$$

The real part of the refractive index N' is related to the imaginary part N'' via the Kramers–Krönig relation. Henry has introduced the linewidth enhancement factor [12], which relates small variations of the real and imaginary parts of the refractive index due to a change of the carrier concentration

$$\alpha_T = \frac{\left.\dfrac{\partial N'}{\partial n}\right|_T}{\left.\dfrac{\partial N''}{\partial n}\right|_T} = \left.\frac{\partial N'}{\partial N''}\right|_T. \tag{8}$$

We label this parameter α_T, which reflects the fact that variations are due to a change of carrier concentration at constant carrier temperature. This parameter can also be defined as a derivative with respect to the quasi-Fermi level separation

$$\alpha_T = \frac{\left.\dfrac{\partial N'}{\partial \mu}\right|_T}{\left.\dfrac{\partial N''}{\partial \mu}\right|_T}. \tag{9}$$

Obviously the Kramers–Krönig relation suggests that there is an analogous parameter relating small variations of the real and imaginary parts of the refractive index due to the change of the carrier temperature at a constant value of the quasi-Fermi level separation

$$\alpha_\mu = \frac{\left.\dfrac{\partial N'}{\partial T}\right|_\mu}{\left.\dfrac{\partial N''}{\partial T}\right|_\mu} = \left.\frac{\partial N'}{\partial N''}\right|_\mu. \tag{10}$$

This definition is very close to one by Hultgren and Ippen [15]. In order to avoid ambiguity, we use subscript symbol to indicate explicitly the parameter held constant. Combining (6)–(10), we can relate the change of the real part of the refractive index and change of the carrier temperature

$$dN' = -(\alpha_T - \alpha_\mu)\cdot \left.\frac{\partial N''}{\partial T}\right|_\mu \cdot dT. \tag{11}$$

We can now rewrite this formula in terms of laser parameters that can be measured experimentally. The change of the real part of the refractive index in the active layer is approximately proportional to the change of the effective group index, which in turn is proportional to the change of the wavelength

$$\Gamma\frac{dN'}{N'} \approx \frac{dN_{\mathrm{eff}}}{N_{\mathrm{eff}}} = \frac{d\lambda}{\lambda}. \tag{12}$$

Using (5), (11), and (12), we express the change of the lasing wavelength due to the change of the carrier temperature in the active layer

$$d\lambda = \frac{\lambda^2}{4\pi \cdot N_{\mathrm{eff}}}(\alpha_T - \alpha_\mu)\cdot \left.\frac{\partial g}{\partial T}\right|_\mu \cdot dT. \tag{13}$$

Carrier heating lifts the pinned carrier concentration (and quasi-Fermi level separation) above threshold and results in wavelength chirp. The carrier heating affects the lasing wavelength in two ways described by the two terms in brackets $(\alpha_T - \alpha_\mu)$ in (13). The first term corresponds to the change of the real part of the refractive index due to change of the quasi-Fermi level separation. The second term describes the change in real part of the refractive index due to the change of the gain profile produced by the carrier heating.

To determine accurately the change of wavelength with current in a small-signal approach, one can measure the frequency chirp. This can be done using with Fabry–Perot etalon or in a gated delayed self-homodyne technique [16]. Rewriting (13) for frequency chirp due to current modulation, we obtain

$$\begin{aligned}\frac{dT}{dI} &= \frac{4\pi \cdot N_{\mathrm{eff}}}{(\alpha_T - \alpha_\mu)\cdot\lambda^2}\,\frac{1}{\left.\dfrac{\partial g}{\partial T}\right|_\mu}\cdot\frac{d\lambda}{dI}\\ &= -\frac{4\pi \cdot N_{\mathrm{eff}}}{(\alpha_T - \alpha_\mu)\cdot c}\,\frac{1}{\left.\dfrac{\partial g}{\partial T}\right|_\mu}\,\beta\end{aligned} \tag{14}$$

where $\beta = d\nu/dI$ is the chirp parameter.

III. MEASUREMENT OF α_T AND α_μ

We start with the Kramers–Krönig relation for small changes of refractive index

$$\delta N' = \frac{2}{\pi}P\int_0^\infty \frac{\delta N''(\varepsilon')\cdot\varepsilon'}{\varepsilon'^2 - \varepsilon^2}\cdot d\varepsilon' \tag{15}$$

where P indicates taking the principal part integral. Equation (5) can be used to relate δG and $\delta N''$ as

$$\delta N' = -\frac{hc}{2\pi^2}P\int_0^\infty \frac{\delta G(\varepsilon')}{\varepsilon'^2 - \varepsilon^2}\cdot d\varepsilon'. \tag{16}$$

We can now determine the coefficients α_T and α_μ by substituting (16) into (9) and (10), respectively. After some simple algebra

$$\begin{aligned}\alpha_T(\varepsilon) &= \frac{\left.\dfrac{\partial N'}{\partial \mu}\right|_T}{\left.\dfrac{\partial N''}{\partial \mu}\right|_T}\\ &= \frac{2}{\pi}\,\frac{\varepsilon\cdot P\displaystyle\int_0^\infty \delta G(\varepsilon')|_T\,\frac{d\varepsilon'}{\varepsilon'^2-\varepsilon^2}}{\delta G(\varepsilon)|_T}\end{aligned} \tag{17}$$

and

$$\begin{aligned}\alpha_\mu(\varepsilon) &= \frac{\left.\dfrac{\partial N'}{\partial T}\right|_\mu}{\left.\dfrac{\partial N''}{\partial T}\right|_\mu}\\ &= \frac{2}{\pi}\,\frac{\varepsilon\cdot P\displaystyle\int_0^\infty \delta G(\varepsilon')|_\mu\,\frac{d\varepsilon'}{\varepsilon'^2-\varepsilon^2}}{\delta G(\varepsilon)|_\mu}.\end{aligned} \tag{18}$$

Equations (17) and (18) allow us to determine α_T and α_μ from experiment provided that the gain spectrum is known. In order to perform this procedure, gain spectra should be obtained in a very broad energy range. However, the commonly used Hakki–Paoli technique [17] for extraction of the gain spectra from amplified spontaneous emission (ASE) from the laser facet does not allow for that. We determined the gain from true spontaneous emission (TSE) spectra recorded from the side of the laser chip using the relation derived by Henry [18]:

$$G(\varepsilon) \propto \frac{I_{\mathrm{sp}}(\varepsilon)}{\varepsilon^2}\cdot\left\{1 - \exp\left(\frac{\varepsilon-\mu}{kT}\right)\right\}, \tag{19}$$

where $I_{\rm sp}(\varepsilon)$ is the TSE intensity. Equation (19) is based on the assumption that carriers have Fermi distribution function. This assumption is supported by experimental data [13] as we discussed above. The quasi-Fermi level separation μ is equal to the transparency energy, which can be measured using the Andrekson technique [19]. Equations (17) and (18) require some caution in treating the singularity in the integrand at $\varepsilon' = \varepsilon$, a procedure described by Henry [20]. It should also be noticed that the TSE intensity can be measured in arbitrary units because only the ratio of the integral of the spontaneous emission to its value at a certain energy enters (17) and (18).

IV. EXPERIMENTAL RESULTS

Lasers used for this work were capped-mesa buried heterostructure MQW DFB lasers [21], [22]. For the gain measurements, we used 300-μm-long uncoated devices. The threshold current was about 11 mA at 25 °C. Shown in Fig. 1 are TSE spectra recorded from the side of the laser chip at 25°C for different values of the current. Below threshold the TSE intensity increases rapidly with current as the carrier concentration increases. Above threshold, the carrier concentration changes very slightly (to offset the decrease of gain due to carrier heating as discussed above). This results in a small increase of the TSE intensity as it is seen in Fig. 1. TSE spectra above threshold do not show a dip, which would be expected in the case of spectral hole burning.

The TSE spectra similar to those shown in Fig. 1 were measured to obtain the optical gain using (19). A typical optical gain spectrum is presented in Fig. 2 (solid line). As mentioned above, (19) does not produce the value of gain in absolute units. The gain curves obtained form TSE were scaled so that the one at 10 mA coincides with the gain curve obtained from the ASE spectrum at the same current using the Hakki–Paoli technique (circles). The ASE spectrum is perturbed in the vicinity of DFB wavelength, as is the extracted gain spectra. Away from the DFB wavelength the gain curves extracted from the TSE and ASE lie very close to each other. As discussed in the previous section, it is not necessary for determination of α_T and α_μ to normalize the optical gain obtained from the TSE spectra. However, it provides a good check for the measurement accuracy and also allows us to determine $(\partial g/\partial T_{e\text{-}h})|_\mu$ from the above results. In these measurements, we changed the heat-sink temperature, thus varying both the crystal lattice and carrier plasma temperature. In order to obtain the value of $(\partial g/\partial T_{e\text{-}h})|_\mu$, the derivative with respect to the carrier temperature, we need to exclude the effect of lattice heating. This can be done by taking into account the fact that the lattice heating causes bandgap energy reduction at the rate of approximately 4 Å/K [24], [25]. We also measured the shift of the electroluminescence inflection point for the test sample and obtained the same result. Therefore, to estimate the derivative from measured gain curves, we shifted the wavelength with the same rate. From the gain curves obtained at heat-sink temperatures of 25 °C and 35 °C and shifted by 4 nm, we determined the value of $(\partial g/\partial T_{e\text{-}h}|_\mu \approx -0.45 \pm 0.05$ (cm^{-1})/K. The material gain and confinement factor do not depend on the value of the total loss of a particular laser, therefore this value can be used for the lasers with different coatings.

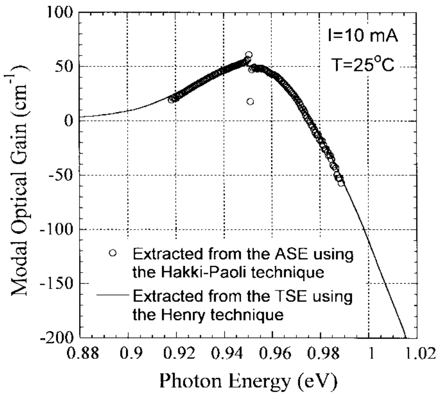

Fig. 2. Modal optical gain spectra of an uncoated DFB laser determined from the TSE spectra at 10 mA using the Henry technique (solid line), and the modal gain spectra extracted from the ASE spectra at the same operating conditions using the Hakki–Paoli technique (empty circles).

Fig. 3 shows the modal gain spectra obtained from the TSE spectra for currents of 8 and 10 mA at 25 °C. We determined the quasi-Fermi level separation at each current using the modified Andrekson technique [23]. This method is not sufficiently accurate to ensure the independence of the gain coefficient on the current at high energies. Therefore, a small adjustment for the transparency energy was needed to make the gain curves converge at high energies, a procedure analogous to that used by Henry [18]. This conversion of the gain curves is not seen in Figs. 3 and 4 because the gain curves are plotted in a limited range of wavelength. The modal gain at 10 mA is higher than at 8 mA in the entire energy range, but particularly for energies above the lasing energy. When (17) is used, this results in a positive value of α_T. Using the gain curves plotted in Fig. 3, we determined the value of $\alpha_T \approx 2.1 \pm 0.2$. We also used an alternative way to determine α_T from ASE spectra below threshold. In this technique, the parameter α_T is determined from the ratio of the change of the modal gain and shift of the wavelength of Fabry–Perot peaks with current below threshold [26]. This procedure gave a value $\alpha_T \approx 2.0 \pm 0.2$. These two results are in a very good agreement that validates the use of (18) to determine α_μ. Fig. 4 shows the modal gain spectra obtained from TSE spectra at two different conditions: 10 mA at 25 °C (solid line) and 10.9 mA at 35 °C (dashed line). To measure the parameter α_μ, we need to vary the carrier temperature and current at the same time to maintain the quasi-Fermi level separation constant. Because we varied the heat-sink temperature, we shifted the quasi-Fermi level separation to compensate for the bandgap shrinking as discussed above. We also used the same procedure of fine tuning the quasi-Fermi level separation as in the measurement of α_T. The modal gain at 35 °C is lower than one at 25 °C over a small range of energies below the transparency energy. The curves cross at transparency energy and at higher energies the modal gain at 35 °C is higher. This behavior results from the dependence of Fermi function on temperature. The high-energy part of the gain spectrum is dominant when used in (18), which results

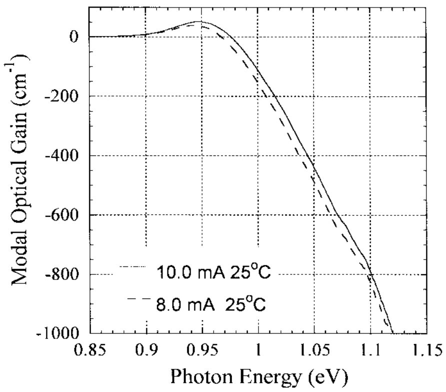

Fig. 3. Modal gain spectra determined from spontaneous emission using the Henry technique at 25 °C for currents of 10.0 (solid line) and 8.0 mA (dashed line).

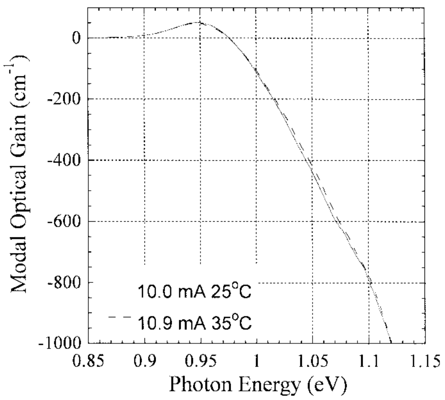

Fig. 4. Modal gain spectra determined from spontaneous emission using the Henry technique for currents of 10.0 mA at 25 °C (solid line) and 10.9 mA at 35 °C (dashed line).

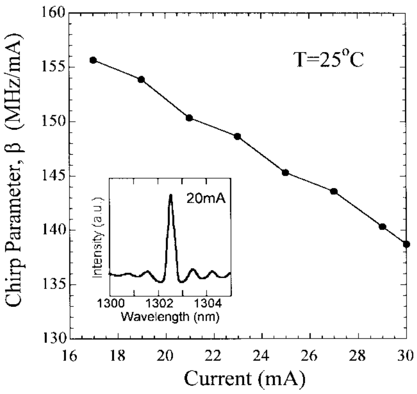

Fig. 5. The chirp parameter β as a function of current. The inset shows the radiation spectra of the same laser above threshold.

in a negative value for α_μ. For the conditions in Fig. 4, we obtained the value $\alpha_\mu \approx -1.4 \pm 0.3$. The fact that α_T and α_μ have different signs indicate that both carrier heating and the increase of carrier concentration contribute to increase of wavelength chirp. Authors in [15] obtained positive value for both α parameters. The source of the difference may arise from the fact that in pump-probe experiments none of the parameters (out of n, μ, and T) is held constant under experimental conditions.

To measure the chirp, we selected a laser from the same wafer, which had coated mirrors and a mid-band longitudinal mode, so that the right side of (4) is zero as was discussed above, and we can use (14). The threshold current was 16.6 mA. The radiation spectrum at 20 mA is shown in the inset of Fig. 5. We measured the frequency chirp using the carrier suppression method in a gated delayed self-homodyne technique [16]. In this technique, the small-signal modulation of the pumping current is gated with a duty cycle of 50% and the period is equal to twice the delay time of the fiber-optic interferometer. We used a standard interferometer (HP11980A) with a delay time of 3.5 μs. The radiation from the laser is coupled into the interferometer input and the optical signal from the interferometer output is analyzed with an RF spectrum analyzer (HP71400). The amplitude of this signal at the modulation frequency reaches its minimum when the first-order Bessel function $J_1(\beta I_{\rm mod}/f_{\rm mod})$ is zero ($\beta I_{\rm mod}/f_{\rm mod} = 3.84$). Therefore, by varying the current modulation depth $I_{\rm mod}$ to minimize the RF signal at $f_{\rm mod}$ we found the chirp parameter β. The chirp parameter as a function of dc current is shown in Fig. 5. At currents close to threshold where the effect of spatial hole burning is minimal, the chirp parameter is 156 MHz/mA. In this paper, we measured wavelength chirp at a modulation frequency $f_{\rm mod} = 100$ MHz, which is high enough to eliminate chirping due to joule heating. The modulation frequency and amplitude are also small enough to make transient effects negligible, so the measured chirp is adiabatic [10]. Using (14), we then get the rate of the change of carrier temperature with a current of approximately 0.13 K/mA. The accuracy of this estimation is about 25%.

V. DISCUSSION

The carrier heating is considered to be a cause of the wavelength chirp. This assumption is based on the results of experimental studies of fast gain dynamics in laser structures [15]. Our measurements of spontaneous emission above threshold also do not show the signature of spectral hole burning. Sources of carrier heating above threshold are injection of energetic carriers through heterobarriers and free-carrier optical absorption. The first effect depends on the injection current, and the second on the optical field. Therefore, in a first-order approximation, both effects are proportional to the difference between operating and threshold currents.

We found the proportionality coefficient between the change of the carrier temperature above threshold and the wavelength chirp. This coefficient is expressed in terms of the temperature derivative of modal optical gain and two parameters α_μ and α_T. These parameters relate the variation of the real part of the refractive index and the variation of the optical gain due to change of carrier temperature or change of the quasi-Fermi level separation.

We have measured the temperature dependence of the optical gain, the parameters α_μ and α_T, and chirp at threshold for MQW InGaAsP lasers. Using this data, we obtained the rate of the temperature increase with current above threshold, which is 0.13 K/mA. The measurements are carried out at a current close to threshold. The observed wavelength chirp above threshold is orders of magnitude smaller than the wavelength shift with current below threshold. Therefore, the variation of the carrier concentration as well as estimated change of carrier temperature are very small that justifies the use of a linear extrapolation. This allows us to estimate the carrier temperature difference between a threshold of 15 mA and a typical operating current of 50 mA to be 4.7 K, which is consistent with the results of Henry [4].

Wavelength chirp is a very important parameter for semiconductor lasers. Different applications impose different requirements on the chirp. For example, lasers for CATV applications should have high chirp [27] to reduce the interferometric noise, while digital applications require small chirp to increase the transmission distance. We believe that the new characterization technique demonstrated in this paper can be useful for understanding of wavelength chirp and its dependence on the laser design.

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