Intrinsic Pulsation in Stripe-Geometry DH Semiconductor Lasers

CHANG-ZHI GUO AND KAI-GE WANG

Abstract—**The effects of lateral optical and carrier distributions and their interaction on dynamic processes in stripe-geometry DH semiconductor lasers are analyzed theoretically. The rate equations and field equations are rigorously solved simultaneously for numerical transient solutions. The calculated results show that, when the index waveguiding of the carrier distribution and the carrier diffusion become pronounced, in a certain range of injected current, the "intrinsic" pulsation, whose distinct feature is that the lateral optical beam width oscillates continuously and significantly during pulsation, may occur in the laser output.**

I. Introduction

THE self-pulsation phenomenon in a semiconductor laser is the main transient behavior limiting the device in high-speed applications, and is also one of the important problems still unsolved in the realm of semiconductor laser physics. Since the discovery of this phenomenon, a vast amount of experimental and theoretical investigations have been carried out [1]-[14], but most of them concern certain imperfections in device technology. It seems certain that, along with the improvement in device technology, such an extrinsic self-pulsation will eventually disappear, but can self-pulsation be completely overcome by that time? In addition, in order to emphasize the effects of various imperfections, the lateral distributions of electrons and photons along the stripe have often been ignored, and an oversimplified assumption of uniform distribution has been taken in most works. In conventional stripe-geometry lasers, however, nonuniform carrier distribution formed by the lateral carrier diffusion and lateral current spreading provides the gain and index guidings, which confine the optical field nonuniformly into certain tranverse mode patterns, thus there will be relatively complex interaction between them. Evidently, such physical processes related closely to the device structure will intrinsically affect the transient behavior.

In the following sections, the roles played by these fundamental processes in the transient behavior of semiconductor lasers, and the "intrinsic" self-pulsation possibly induced by them are examined theoretically in a rigorous manner. The effect of lateral optical field distribution (in this section its relative distribution, and thus, its beam width, are supposed to be unchanged) on the hole-burning in the lateral carrier distribution are treated in Section II, and the effects when the lateral carrier distribution plays an important role in waveguiding are treated in Section III.

Manuscript received April 16, 1982.

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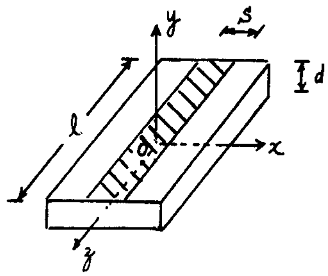


Fig. 1. The scheme of the active region in a stripe-geometry DH semiconductor laser.

II. Theoretical Evaluation

A. Rate Equations with Lateral Distributions

The scheme of the active region in stripe-geometry DH semiconductor lasers is shown in Fig. 1. The longitudinal direction (z direction) is terminated by two end facets; the y direction, perpendicular to the junction plane, is limited by the heterojunctions and the x direction, parallel to the junction plane, extends to infinity.

Assume that the photon density at the point (x, y, z) at time t is

$$N(x,y,z,t) = A\,|\psi(x,y,z,t)|^2 \tag{1}$$

where A is a constant, and ψ is the field distribution function. The total number of photons inside the cavity is

$$N = A\int_{-l/2}^{l/2}\int_{-d/2}^{d/2}\int_{-\infty}^{\infty} |\psi|^2\,dx\,dy\,dz = A\iiint_v |\psi|^2\,dv \tag{2}$$

where l is the distance between the two facets and d is the thickness of the active region. Then the rate equation of the total number of photons in the cavity is given by

$$\frac{dN}{dt} = \left(Gv - \frac{1}{\tau_p}\right)N + \gamma\frac{\mathfrak{N}}{\tau} \tag{3}$$

where the mode gain G is given by

$$G = \iiint_v g(n)|\psi|^2\,dv \bigg/ \iiint_v |\psi|^2\,dv, \tag{4}$$

the total number of electrons is

$$\mathfrak{N} = \iiint_v n\,dv, \tag{5}$$

and the electron density at the point (x,y,z) at time t is

$$n = n(x,y,z,t) \tag{6}$$

where $g(n)$ is the gain coefficient, τ_p and τ are the photon and electron lifetimes, respectively, and γ is the spontaneous emission factor.

The rate equation of the electron density is given by

$$\frac{\partial n}{\partial t}(x,y,z,t) = j(x,y,z,t)/qd - n(x,y,z,t)/\tau + D_n \nabla^2 n(x,y,z,t) - gv(n(x,y,z,t)) \cdot N(x,y,z,t) \tag{7}$$

where $j(x, y, z, t)$ is the injected current density, q is the electron charge, and D_n is the electron diffusion coefficient. Assume that

$$\psi(x,y,z,t) = \psi_x(x)\psi_y(y)\psi_z(z)e^{-i\omega t}; \tag{8a}$$

then from (1) and (2)

$$N = \mathbf{N}|\psi|^2 \Big/ \iiint_v |\psi|^2 \, dv = \mathbf{N}|\psi_x|^2 \, |\psi_y|^2 \, |\psi_z|^2 \Big/ \iiint_v |\psi|^2 \, dv. \tag{8b}$$

As the carrier and current density in the y and z directions can be considered as uniform [i.e., $n = n(x, t), j = j(x, t)$], and the relative field distribution in the y direction remains unchanged, we can integrate (7) with respect to y and z and denote the average photon and electron densities by

$$\overline{N} = \mathbf{N}/sld \tag{9a}$$

$$\bar{n} = \mathfrak{N}/sld = \frac{1}{s}\int_{-\infty}^{\infty} n \, dx, \tag{9b}$$

respectively. Then we obtain the rate equations

$$\frac{\partial}{\partial t} n(x,t) = j(x,t)/qd - n(x,t)/\tau + D_n \frac{\partial^2 n(x,t)}{\partial x^2} - gv(n)s|\psi_x|^2 N^{(t)} \Big/ \int_{-\infty}^{\infty} |\psi_x|^2 \, dx \tag{10a}$$

$$\frac{d}{dt}\overline{N}(t) = \left[G(t)v - \frac{1}{\tau_p}\right]\overline{N}(t) + \gamma\bar{n}(t)/\tau \tag{10b}$$

where s is the stripe width defined by the injected current and the mode gain may be written as

$$G(t) = \int_{-\infty}^{\infty} g(n)|\psi_x|^2 \, dx \Big/ \int_{-\infty}^{\infty} |\psi_x|^2 \, dx. \tag{10c}$$

B. Fixed Lateral Optical Field Distribution and Fixed Beam Width

We assume that, during the transcience, the relative field is independent of the change of waveguiding caused by change in carrier distribution and is taken to be a Gaussian function

$$\psi_x = \left(\sqrt{\frac{8}{\pi}}\frac{1}{w}\right)^{1/2} e^{-4x^2/w^2} \tag{11}$$

where w is the full beam width at the e^{-2} intensity points. The gain function is taken to be

$$g(n) = an - b \tag{12}$$

where a and b are constants. Suppose that the step current is applied at the time $t = 0$ and its spatial distribution is given by

$$j(x) = \begin{cases} j \text{ (constant)}, & \text{when } |x| \leqslant s/2 \\ 0, & \text{when } |x| > s/2. \end{cases} \tag{13}$$

In order to save computer time, the initial electron density distribution is taken to be the threshold distribution

$$\frac{n(x,0)}{\bar{n}_{\text{th}}} = \frac{n_{\text{th}}(x)}{\bar{n}_{\text{th}}} = \begin{cases} 1 - \frac{1}{2} e^{-s/2L_n}(e^{x/L_n} + e^{-x/L_n}), & |x| \leqslant s/2 \\ \frac{1}{2}(e^{s/2L_n} - e^{s/2L_n})e^{-|x|/L_n}, & |x| > s/2 \end{cases} \tag{14}$$

where $L_n = \sqrt{D_n \tau}$ is the electron diffusion length and

$$\bar{n}_{\text{th}} = \frac{\left(b + \frac{1}{\tau_p}\right)\int_{-\infty}^{\infty} |\psi_{x\,\text{th}}|^2 \, dx}{a\int_{-\infty}^{\infty} \frac{n_{\text{th}}(x)}{\bar{n}_{\text{th}}} |\psi_{x\,\text{th}}|^2 \, dx}. \tag{15}$$

Equations (10)-(15) are calculated numerically by the difference method as follows. The differential equations are transformed into difference equations, and the values are iterated successively from an initial time $t = 0$. The time step is 0.05 percent of τ, and the distance step is 2 percent of the stripe width s. The numerical values of relevant parameters used in the calculations are $\tau = 2.5 \times 10^{-9}$ s, $\tau_p = 2.5 \times 10^{-12}$ s, $a = 1.08 \times 10^{-16}$ cm^2, $b = 146$ cm^{-1}, $\gamma = 10^{-5}$, and the initial photon density $\overline{N}(0) = 0.0015 \times 10^{15}$ cm^{-3}.

Fig. 2. shows typical results calculated by the computer, where $j/j_{\text{th}} = 1.05$, $w/s = 1.67$, and $L_n/s = 0.1$. In Fig. 2 the optical output $\overline{N}(t)$ is a relaxation oscillation. The mode gain $G(t)$ oscillates about the threshold gain, but the average electron density $\bar{n}(t)$ in the cavity increases gradually up to a steady value, which is above the average threshold electron density $\bar{n}_{\text{th}}$. This is because the steady-state electron distribution depends on the injected current density, and for higher injection, the center of the gain distribution where the optical intensity is a maximum is depressed, and in order to achieve the threshold gain the electron density must be increased. Thus, when the injected current is changing, the clamping of the mode gain does not correspond to the clamping of the average electron density $\bar{n}_{\text{th}}$ in the cavity. This is different from the simplified uniform rate equations where clamping of the gain must correspond to clamping of the electron density.

Fig. 3 gives the calculated photon density $\overline{N}(t)$ of some beam widths at $j/j_{\text{th}} = 1.05, L_n/s = 0.5$. Fig. 3 shows that the degree of optical field spreading has a pronounced influence on the

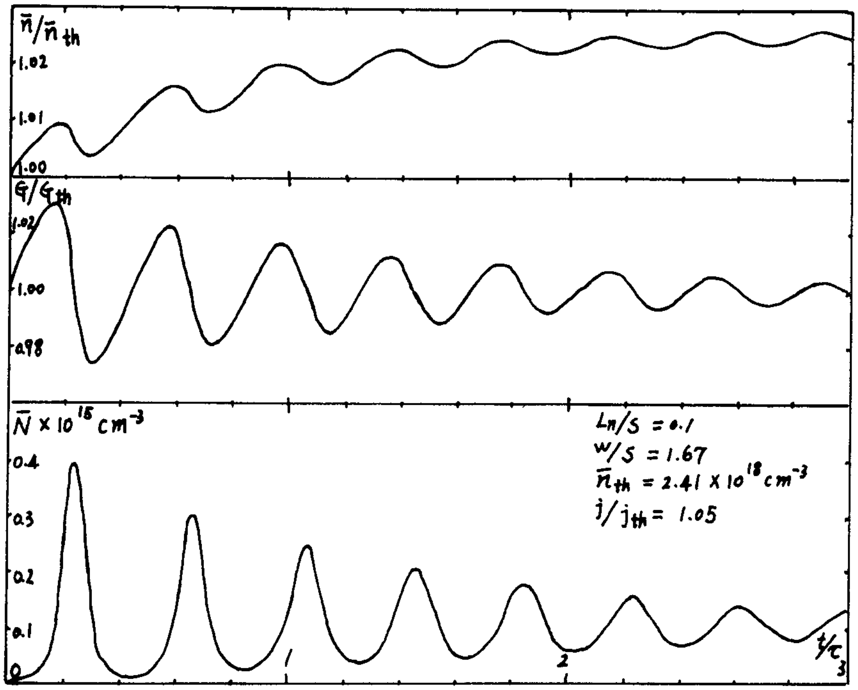


Fig. 2. Time evolutions of carrier density ($\bar{n}/\bar{n}_{th}$), mode gain (G/G_{th}), and photon density ($\bar{N}$) for $(W/S) = 1.67$, $(L_n/S) = 0.1$ at $(j/j_{th}) = 1.05$.

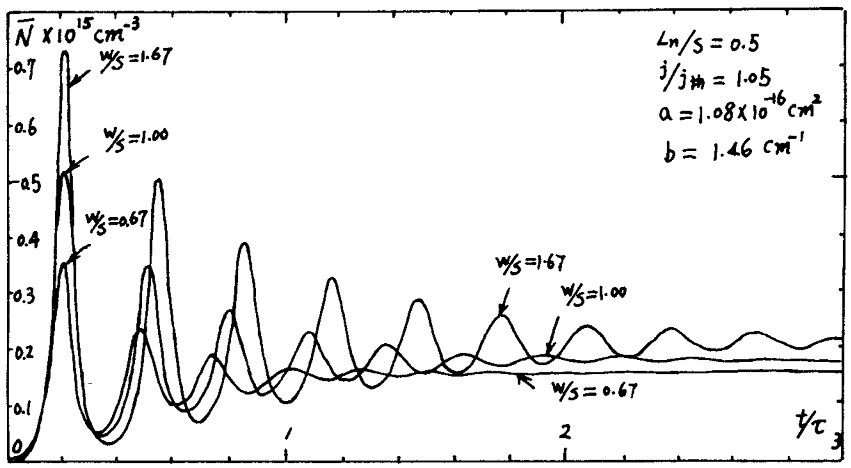


Fig. 3. Time evolutions of photon density $\bar{N}(t)$ for various beam widths at $(j/j_{th}) = 1.05$, $(L_n/S) = 0.5$.

transient behavior; when the width is wide, the optical spikes are high and attenuation is slow. It may be explained by the enhancement of the effect of the absorbing region. Fig. 4 gives the distributions of the relative optical intensity, gains and absorbing regions with the above parameters. When the beam width is narrow, the optical intensity is concentrated in the gain region and the absorbing region is not effective; when the optical field spreads out of the stripe, the fraction of its wings in the absorbing regions increases and the saturable absorber is formed, thus the spikes are enhanced. Therefore, we conclude that the "intrinsic" saturable absorber, which is formed by the stripe-geometry of the laser, cannot give rise to self-pulsation, although it may be of some help to its occurrence. (Our calculation shows that this remains true even if the slope of the negative gain is doubled, for the field in the absorbing region is still weak.)

C. The Optical Field Distribution Depends on the Carrier Distribution

The lateral optical field of the proton bombarded stripe-geometry laser is confined mainly by the gain guiding and index guiding of the lateral carrier distributions formed by lateral carrier diffusion [15]-[17]. During the transient process, variation of the field distribution is bound to be accompanied by variation of the carrier distribution. To reflect this process, we have to solve the simultaneous equations for the rate equations with the field equation. The lateral field distribution is determined by the field equation

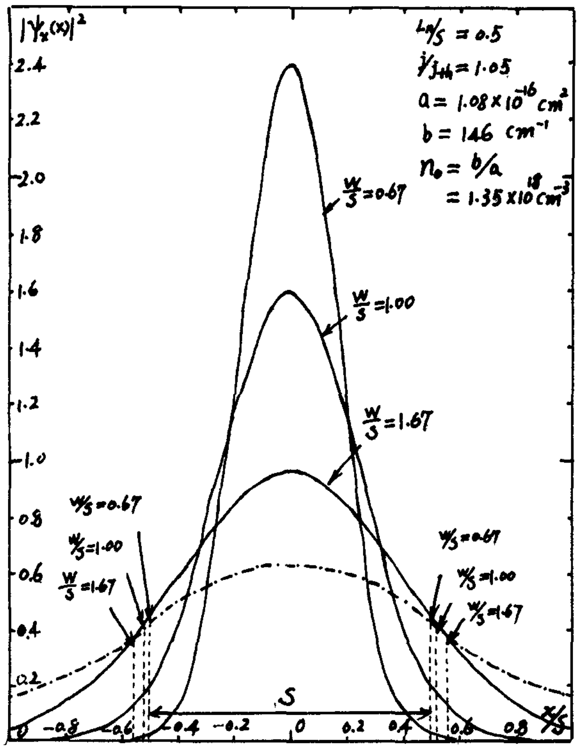


Fig. 4. Lateral relative intensity distributions ($|\psi(x)|^2$) of Gaussian modes with various beam widths (——) and lateral distribution of carrier density ($n/\bar{n}$) at $(j/j_{th}) = 1.05$, $(L_n/S) = 0.5$ (–·–·–). The dashed lines (– – –) denote the bounderies between the gain region (middle) and the absorbing regions (two wings).

$$\frac{d^2}{dx^2}\psi_x(x) + [k_0^2\tilde{\epsilon}(x) - \beta_x^2]\,\psi_x(x) = 0 \tag{16}$$

where $k_0 = 2\pi/\lambda_0$, λ_0 is the optical wavelength in vacuum, β_x^2 is the eigenvalue, and $\tilde{\epsilon}$ is the complex dielectric constant. Suppose

$$\tilde{\epsilon} = \tilde{\epsilon}_0 - \tilde{A}n \tag{17}$$

where $\tilde{\epsilon}_0$ is the complex dielectric contant when the injected carrier density equals zero and $\tilde{A}$ is a complex parameter. Let

$$\tilde{\epsilon} = \epsilon_r + i\epsilon_i \tag{18a}$$

$$\tilde{A} = A_r + iA_i \tag{18b}$$

where $i = \sqrt{-1}$. Here A_r and A_i are given by

$$A_r = -2\eta_r\frac{\partial\eta_r}{\partial n} \tag{19a}$$

$$A_i = \frac{\eta_r}{k_0}\frac{\partial g}{\partial n} \tag{19b}$$

where η_r is the real refractive index. A_r and A_i show the variation rate of the real refractive index and the gain with the carrier density, respectively. Define R as

$$R = A_r/A_i. \tag{20}$$

The magnitude of R shows the relative role played by the index guiding and gain guiding, both mainly formed by the carrier distribution. When $R \to 0$ (i.e., $A_r \to 0$), the guiding approaches pure gain guiding; when R increases, the role played by index guiding becomes more important.

To simplify the calculation, an approximate solution was made by a variation method. According to the variational principle [18], the field equation (16) is equal in value to the variation equation

$$\delta[\beta_x^2] = \delta\left[\frac{\int_{-\infty}^{\infty}\psi_x\left(\frac{d^2}{dx^2} + k_0^2\tilde{\epsilon}(x)\right)\psi_x\,dx}{\int_{-\infty}^{\infty}\psi_x^2\,dx}\right]. \tag{21}$$

Suppose that the trial function ψ_x in the variation equation (21) is the Hermite-Gaussian function, we consider only the fundamental mode as a first approximation

$$\psi_x = ce^{-1/2\,\tilde{\alpha}x^2} \tag{22}$$

where c is the normalized constant and $\tilde{\alpha}$ is a complex parameter. Substitute (22) into (21) and let [18]

$$\frac{\partial}{\partial\tilde{\alpha}}[\beta_x^2] = 0 \tag{23}$$

giving a transcendental equation

$$\int_{-\infty}^{\infty} n(x)(2\tilde{\alpha}x^2 - 1)e^{-\tilde{\alpha}x^2}\,dx = \frac{\sqrt{\pi\tilde{\alpha}}}{k_0^2\tilde{A}} = \frac{\sqrt{\pi\tilde{\alpha}}}{k_0^2A_i(R+i)}. \tag{24}$$

Equation (24) can be taken as an approximation for the field equation (16), and when $n(x)$ is given, the values of $\tilde{\alpha}$ and ψ_x can be obtained. Let

$$\tilde{\alpha} = \alpha_r + i\alpha_i, \tag{25}$$

the relation between beam width w and α_r is

$$w = \sqrt{\frac{8}{\alpha_r}}, \tag{26a}$$

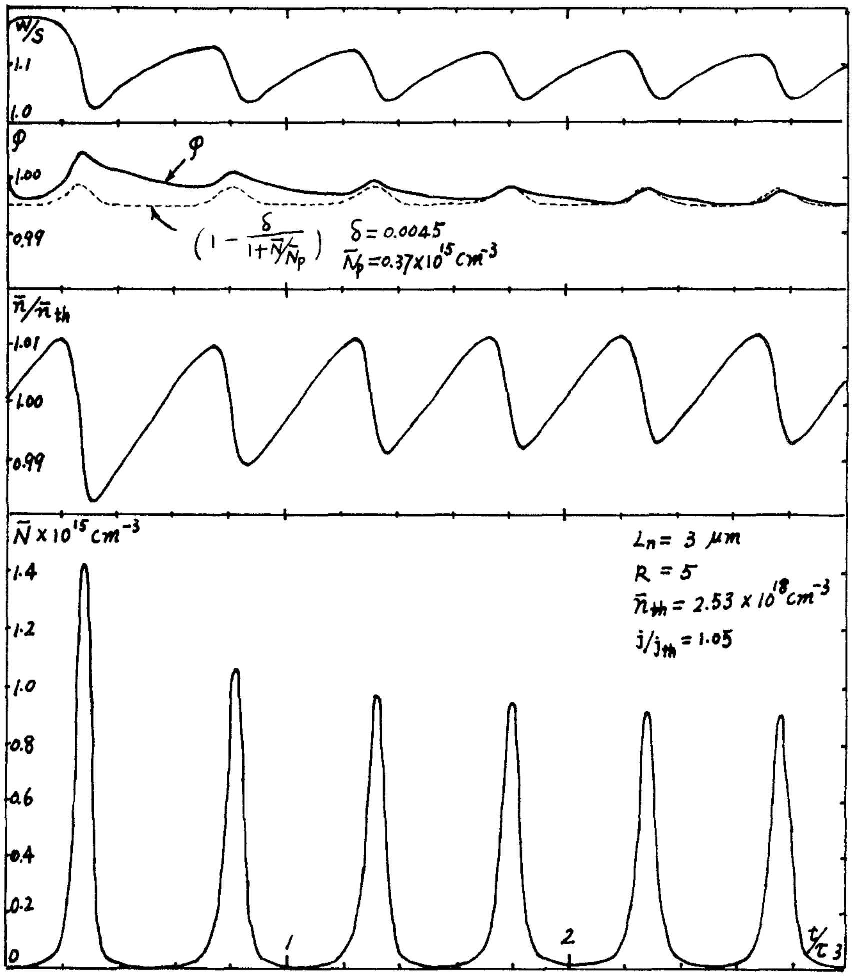
and the relation between α_i and the curvature ρ of the phase front of the beam at the center ($x = 0$) of the facet is

$$\rho \approx \frac{\alpha_i}{k_0\eta_r} \tag{26b}$$

approximately. This means that when $\rho < 0$ the waist of the Gaussian beam is inside the cavity, otherwise it is outside the cavity.

We combined the rate equations (10) with the approximate field equation (24) and made the numerical calculation with (12)-(15), (17)-(20), (22), (25), and (26). The algorithm, steps, and parameter τ, τ_p, a, b, $\bar{N}(0)$ used are the same as above, other parameters used are $s = 10\ \mu m$, $\lambda_0 = 0.89\ \mu m$, $\eta_r = \sqrt{\epsilon_{0r}} = 3.6$.

The calculated results show that different values of R and L_n give rise to different transient behaviors. When R and L_n take on large values, then in a certain range of current, the self-



(a)

Fig. 5 Time evolutions of W/S, φ, $\bar{n}/\bar{n}_{th}$, $\bar{N}$ for various values of R and L_n at $(j/j_{th}) = 1.05$, $S = 10\ \mu$m. (a) $R = 5$, $L_n = 3\ \mu$m.

pulsation occurs. The transient behaviors for different values of R and L_n at $j/j_{th} = 1.05$ are compared in Fig. 5. When $R = 5$, $L_n = 3\ \mu$m, self-pulsation occurs in laser output as shown in Fig. 5(a). When $R = 2, L_n = 3\ \mu$m [Fig. 5(b)] and $R = 5, L_n = 1\ \mu$m [Fig. 5(c)], relaxation oscillations appear in the laser outputs. Fig. 6 gives the results for $R = 5$ and $L_n = 5\ \mu$m at $j/j_{th} = 1.01$ [Fig. 6(a)] and $j/j_{th} = 1.2$ [Fig. 6(b)]; both outputs are in self-pulsation.

Fig. 7 gives the results for different R and j/j_{th} at $L_n = 5\ \mu$m. It shows that self-pulsation can only appear in a certain range of injected current and will disappear as the current gradually increases. Fig. 8 gives the results calculated at $j/j_{th} = 2.0$, with $R = 5$ and $L_n = 3\ \mu$m. Fig. 8(a) shows the variations of curvature ρ of the beam phase front, beam width w, average carrier density $\bar{n}$, and average photon density $\bar{N}$ with time. Fig. 8(b) shows the carrier spatial distributions at some typical times.

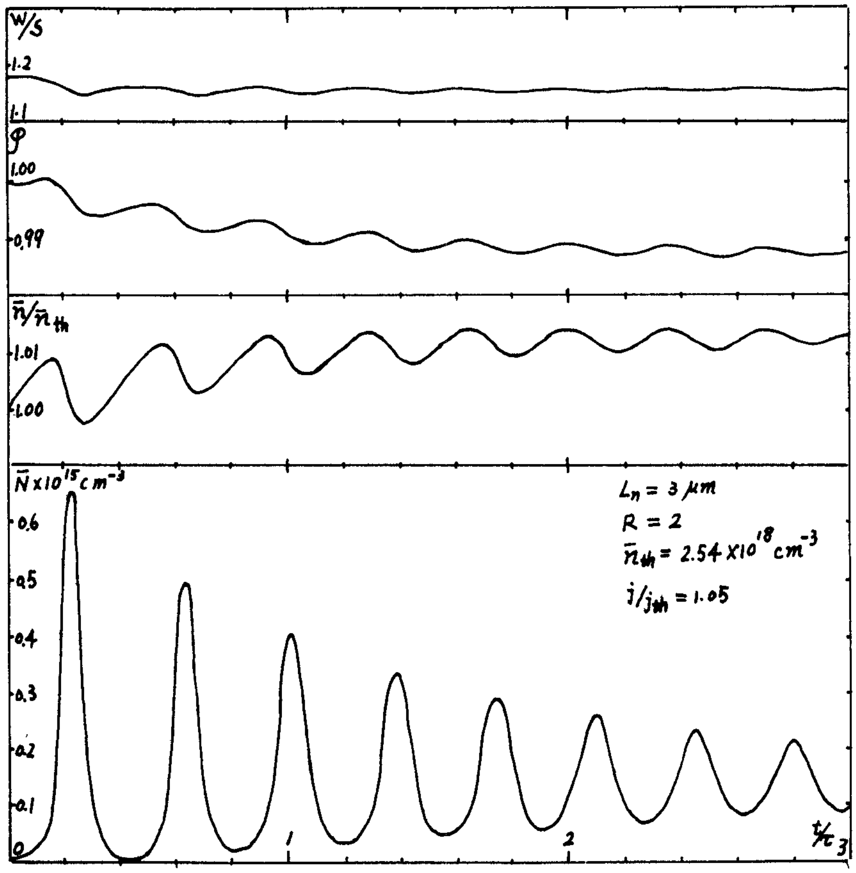
III. Mechanism and Condition for Appearance of Intrinsic Self-Pulsation

On examining the calculated results of Section II-C it is found that, whenever the self-pulsation appears, the beam width oscillates distinctly with large amplitude, as shown in Fig. 5(a) and Fig. 6; but whenever the variation of the beam width is very small, there is no pulsation, as shown in Fig. 5(b) and (c). This indicates that the appearance of this kind of self-pulsation and the instability of the optical field distribution (beam width) are closely related.

We believe that, in general, the appearance of self-pulsation in semiconductor lasers is because the cavity gain or the mode gain G, for some mechanism, depends on the photon density and, thus, becomes nonlinear and, moreover, satisfies the condition

$$\frac{\partial G}{\partial \bar{N}} > 0. \tag{27}$$

For example, if there is some saturable absorber in the cavity, then the above condition can be satisfied and self-pulsation may appear [2], [19]. The intrinsic self-pulsation mechanism considered here will also confirm this view. The instability of the beam width will affect the mode gain causing it to be dependent on the photon density and thus causing it to become nonlinear. When laser spikes occur, the hole-burning effect decreases the beam width and concentrates the field distribu-



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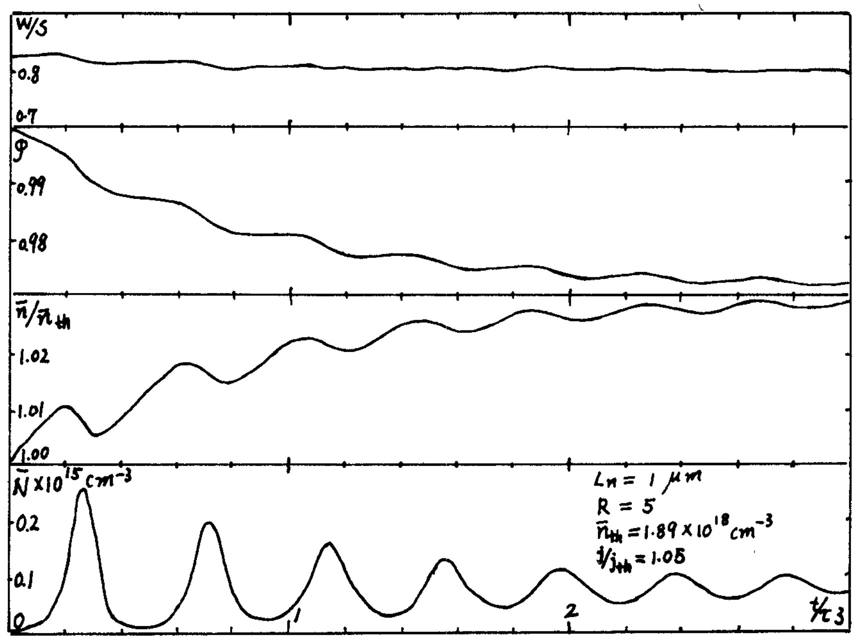


(c)

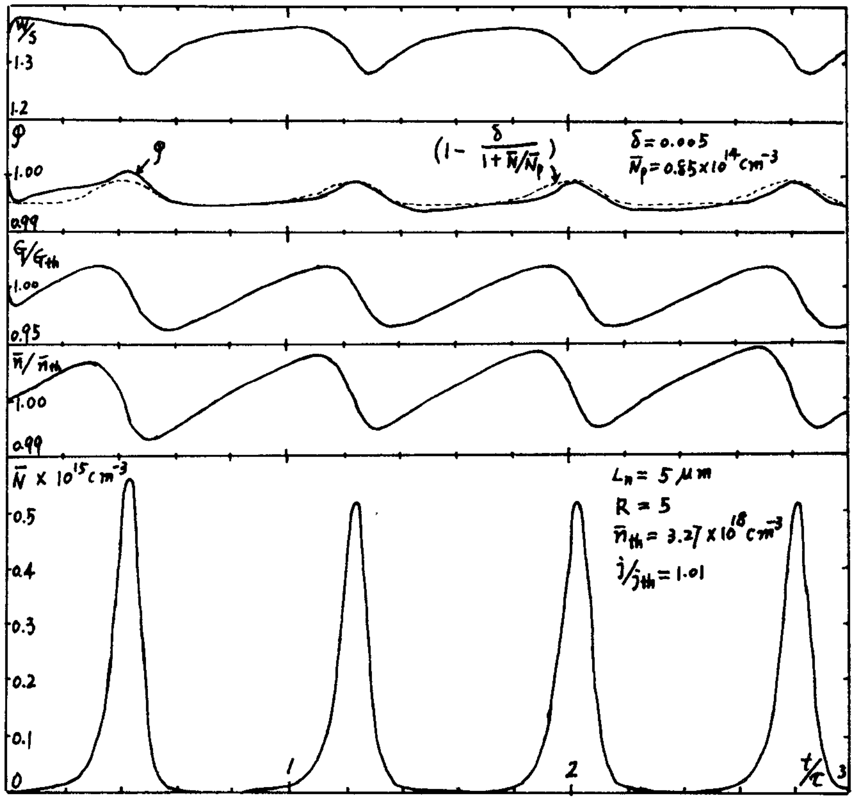
Fig. 5. (*Continued.*) (b) $R = 2, L_n = 3\ \mu\text{m}$. (c) $R = 5, L_n = 1\ \mu\text{m}$.

tion. This will make the mode gain increase and condition (27) will be satisfied.

For example, from (10c) and (12) we obtain

$$G(t) = a \int_{-\infty}^{\infty} n(x, t)|\psi_x|^2\, dx \Big/ \int_{-\infty}^{\infty} |\psi_x|^2\, dx - b. \tag{28}$$

Now let



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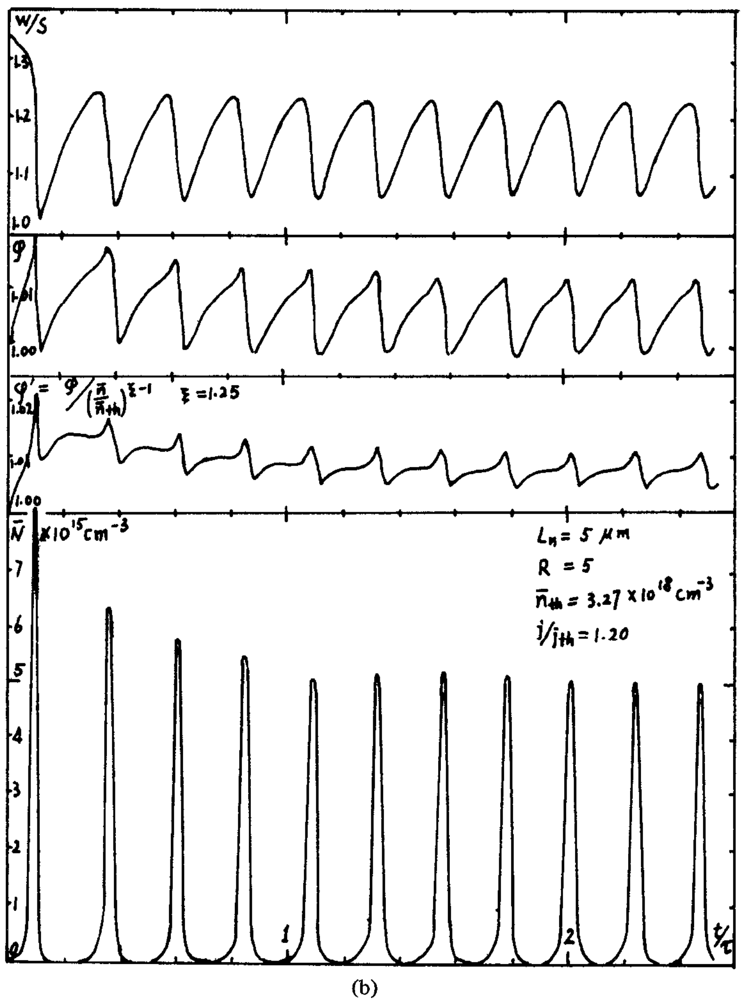


(b)

Fig. 6. Time evolutions of W/S, φ, φ', G/G_{th}, $\bar{n}/\bar{n}_{th}$, $\bar{N}$ for $R = 5$, $L_n = 5\ \mu m$, $S = 10\ \mu m$. (a) $j/j_{th} = 1.01$. (b) $j/j_{th} = 1.2$.

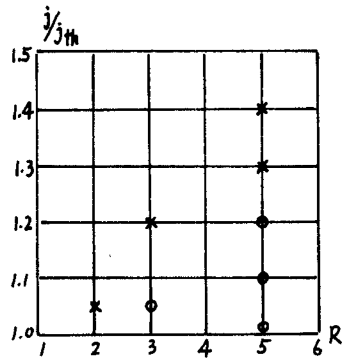


Fig. 7. Calculated injection current regions of relaxation oscillation (+) and self-pulsation (0) for various values of R and L_n at $S = 10$ μm and $L_n = 5$ μm.

$$G(t) = a'\bar{n}(t)\varphi(t) - b \tag{29}$$

and

$$a' = \frac{a\int_{-\infty}^{\infty} n_{th}(x)/\bar{n}_{th}\,|\psi_{x\,th}|^2\,dx}{\int_{-\infty}^{\infty}|\psi_{x\,th}|^2\,dx} \tag{30}$$

for the normalization of φ. From (28)-(30) we obtain

$$\varphi(t) = \frac{G(t)+b}{a'\bar{n}(t)} = \frac{\int_{-\infty}^{\infty}\frac{n(x,t)}{\bar{n}(t)}|\psi_x|^2\,dx\Big/\int_{-\infty}^{\infty}|\psi_x|^2\,dx}{\int_{-\infty}^{\infty}\frac{n_{th}(x)}{\bar{n}_{th}}|\psi_{x\,th}|^2\,dx\Big/\int_{-\infty}^{\infty}|\psi_{x\,th}|^2\,dx}. \tag{31}$$

Here, φ is a nonlinear factor, showing the dependence of the mode gain on the photon density. Using a computer to evaluate the variation of $\varphi(t)$, we found that during self-pulsation φ clearly depends on the photon density. For instance, in Figs. 5(a) and 6, the spikes appearing in the φ curve synchronize with the spikes in the laser output. In some cases, the variation of φ can be approximated to the form $(1 - \delta/1 + \bar{N}/\bar{N}_p)$ (as the dashed line in the figure). For example, we only have to let $\delta = 0.0045$, $\bar{N}_p = 0.37 \times 10^{15}$ cm^{-3} in Fig. 5(a), and $\delta = 0.005$, $\bar{N}_p = 0.85 \times 10^{14}$ cm^{-3} in Fig. 6(a). This kind of mode gain can be approximately equivalent to [9]

$$G = a'\bar{n}\left(1 - \frac{\delta}{1+\bar{N}/\bar{N}_p}\right) - b \tag{32}$$

which obviously satisfies the condition $\partial G/\partial \bar{N} > 0$.

In (32), the value of δ in the nonlinear part which reflects dependence on the photon density is quite small, but it has a decisive influence on the appearance of self-pulsation, although it can hardly be large enough to influence the steady optical power output versus driven current characteristics [20].

In general, the approximate expression (32) does not necessarily hold during self-pulsations. Especially for higher currents, the nonlinear situation may be complicated and φ can even depend on the average electron density $\bar{n}$. If we let

$$G = a'(\bar{n}/\bar{n}_{th})^{\xi}\varphi' - b \tag{33}$$

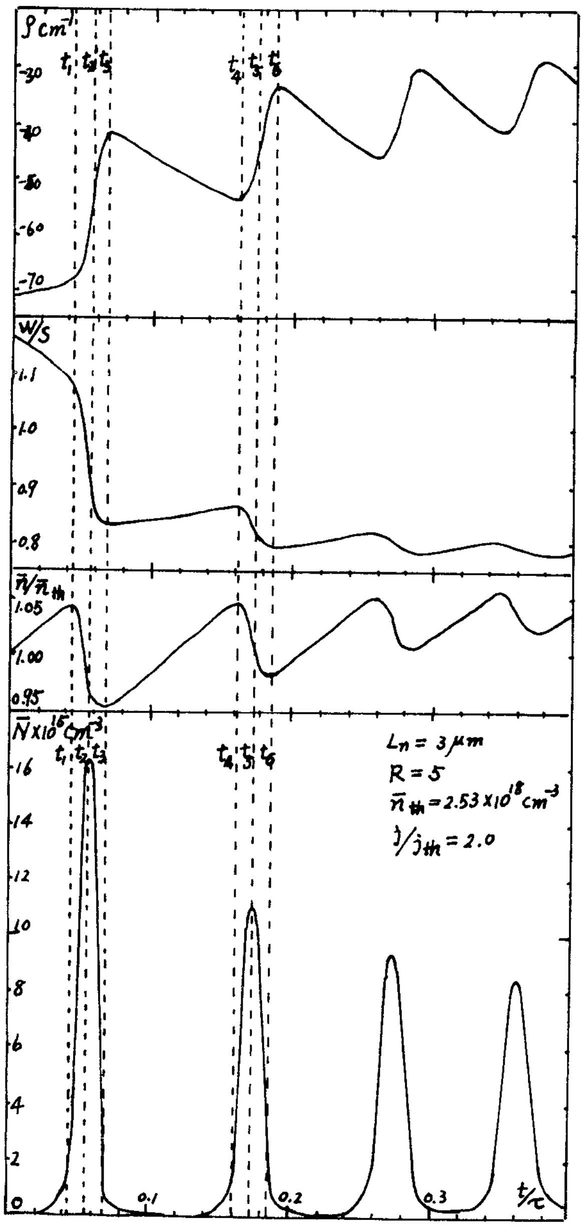
then

$$\varphi' = \frac{\varphi}{(\bar{n}/\bar{n}_{th})^{\xi-1}}. \tag{34}$$

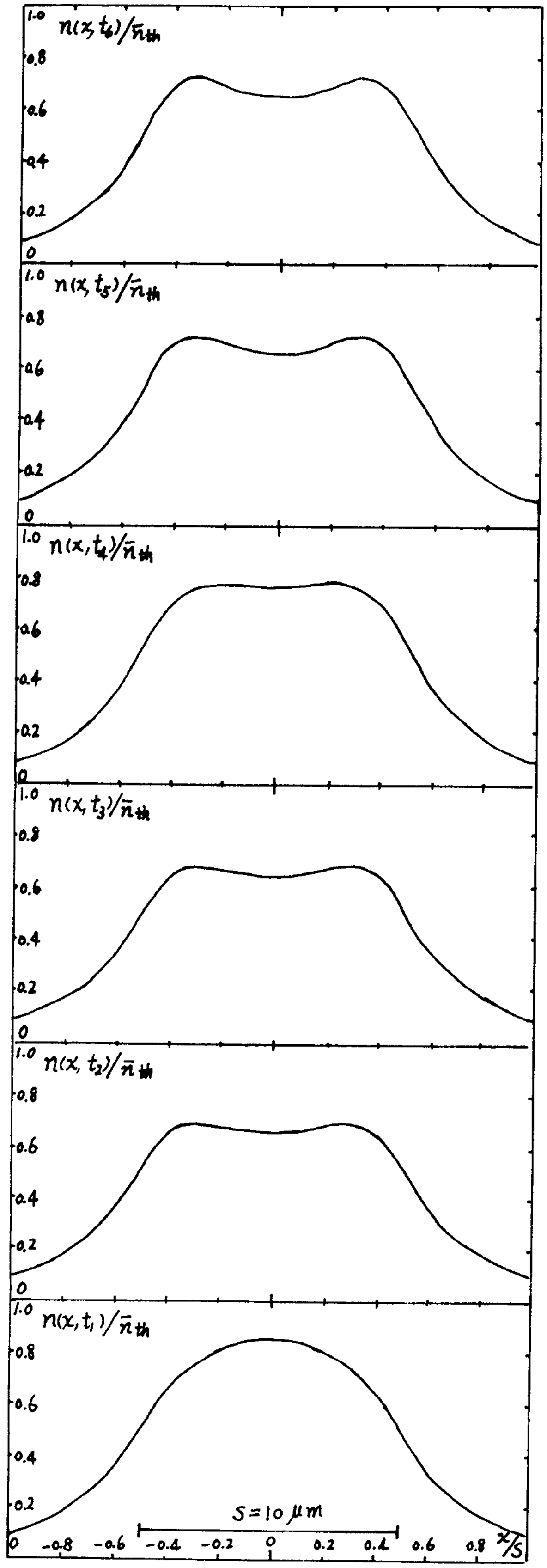
φ' is obtained from φ by deducting the factor depending on $\bar{n}$. φ' shows more clearly the dependence of the mode gain on the photon density, as can be seen from Fig. 6(b), where the synchronization between the spikes in $\varphi'(t)$ and $\bar{N}(t)$ curves is more distinct than for the $\varphi(t)$ and $\bar{N}(t)$ curves.

In the case of nonself-pulsation, the curve for φ has no spikes synchronizing with the optical spikes; its peaks merely synchronize with the peaks of $\bar{n}(t)$, as shown in Fig. 5(b) and (c). This shows that, when the variation of the beam width is small, the mode gain can hardly become nonlinear and thus the self-pulsation can hardly appear.

The appearance of the intrinsic self-pulsation depends mainly on the growth and decline of two processes: the spatial hole-burning of the carrier distribution (this refers in a general sense to any distortion of the carrier distribution by the optical field) and the antihole-burning (filling up the "hole" by carrier diffusion), they both affect the waveguiding, but in opposite ways. During an optical spike, the former process (mainly due to the change of index guiding) concentrates the field and thus increases the mode gain. But during the interval between two optical spikes, the latter process spreads the optical field. If the optical field can be completely restored in such an interval, the oscillation of the beam width, and thus, of the mode gain, will be sustained, and sustained intensity pulsation will appear. Therefore, the appearance of this intrinsic pulsation depends on the magnitude of R, L_n, and j/j_{th}. When R is large, the index guiding induced by the carrier distribution may be pronounced enough to cause significant variation of the beam width; when L_n is large, the antihole-burning is fast enough to maintain repeated oscillation of the beam width, thus self-pulsation results. When j/j_{th} is high, the time interval between every two optical spikes is too short for the antihole-burning to be completed, thus, the beam width oscillation will be impossible to maintain self-pulsation. Fig. 8 shows the whole process. It also shows that the so-called "self-focusing effect" [16], caused by a depression at the center of the carrier distribution, is not necessarily related to self-pulsation. Our calculations have shown that the center of the carrier distribution does not depress into a hole, although the self-pulsation appears in Fig. 5(a) and Fig. 6(a) and (b).



(a)



(b)

Fig. 8. Transient responses for $S = 10\ \mu$m, $R = 5$, $L_n = 3\ \mu$m at $j/j_{th} = 2.0$. (a) Time evolutions of ρ, W/S, $\bar{n}/\bar{n}_{th}$, $\bar{N}$. (b) Lateral carrier distributions at some typical instances during the transiences.

IV. Conclusion

The distinguishing feature of the intrinsic self-pulsation model discussed in this paper from other self-pulsation models is the existence of distinct oscillations of the beam width (about 10 percent). This intrinsic self-pulsation can appear in the proton bombarded stripe-geometry (or oxidized stripe-geometry) lasers. This phenomenon has just been observed in experiments [21]. For the buried stripe-geometry laser, the lateral field is much more stable, thus self-pulsation may often be related to some defects in the device.

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