Supermodes of High-Repetition-Rate Passively Mode-Locked Semiconductor Lasers

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Abstract—We present a steady-state analysis of high-repetition-rate passively mode-locked semiconductor lasers. The analysis includes effects of amplitude-to-phase coupling in both gain and absorber sections. A many-mode eigenvalue approach is presented to obtain supermode solutions. Using a nearest-neighbor mode coupling approximation, chirp-free pulse generation and electrically chirp-controlled operation are explained for the first time. The presence of a nonzero alpha parameter is found to change the symmetry of the supermode and significantly reduce the mode-locking range over which the lowest order supermode remains the minimum gain solution. An increase in absorber strength tends to lead to downchirped pulses. The effects of individual laser parameters are considered, and agreement with recent experimental results is discussed.

I. INTRODUCTION

PREVIOUSLY, the theory of passive modelocking has been analyzed thoroughly in the time domain [1]. Haus' analysis has provided a clear picture of the evolution of pulses through gain, absorptive, and bandwidth-limiting elements within a cavity. A steady-state solution was found when these effects are included. Certain approximations were deemed necessary in order to present an analytic solution. For example, in the steady-state solution, a symmetric and unchirped pulse envelope is assumed as limited by the approximation of all time-domain effects only up to the quadratic term. The model has been extended to include chirped pulses due to self-phase modulation (SPM) yet only for a fast absorber [2], [3], and still restricts the analysis to exponents quadratic in time and achieves symmetric pulses. No recovery is assumed to occur during pulses. Additionally, both models include an approximation of the discrete-mode spectrum by a continuous spectrum. Although the latter approximation works well for mode-locked lasers having many closely-spaced modes, and a slightly-varying gain with frequency, it, along with the assumption of no material recovery during the pulse, is not adequate for the case of high-repetition-rate passively mode-locked lasers (≥50 GHz). In this case, the difference in gain between neighboring modes can be significant, and typically only a small number of modes (around 3–10) dominate.

Active modellocking, on the other hand, has been analyzed thoroughly in both the time domain and the frequency domain [4]–[6]. It has been suggested that passive mode-locking should be analyzed in the time domain since simple products in the time-domain analysis result in cumbersome convolutions in the frequency domain analysis [7], however, in the case of high-repetition-rate passive modelocking, where few modes are involved and the induced carrier modulation is much closer to a sinusoid [8], the frequency domain approach becomes more appropriate. In this paper, we present a steady-state analysis of passive modelocking directed toward high-repetition-rate semiconductor lasers. The analysis is done in the frequency domain extending that presented in [8]. For the first time, passive mode-locking supermodes are found while amplitude-to-phase coupling from slow saturation is permitted. Section II describes the model and arrives at an equation for each mode in the supermode. It incorporates dispersive effects through the common semiconductor laser parameters and unlike previous frequency domain calculations, does not force all modes beyond (the minimum) three modes to contribute zero coupling. Section III describes the eigenvalue formulation used to arrive at a self-consistent solution of the coupled nonlinear equations. Section IV presents an approximate analytical expression based on (the minimum) three modes in order to reduce the complexity and allow one to build physical intuition about the gain requirements and amplitudes and phases of the supermode spectrum. Section V presents results for the full calculation. Section VI compares the results with experiments for high-repetition-rate passively mode-locked lasers. Finally, Section VII includes conclusions.

II. THE MODEL

High-repetition-rate modelocking (≥50 GHz) was first demonstrated by Vasil'ev [9] and by Sanders et al. [10]. To date, semiconductor lasers are the only mode-locked lasers that have been able to generate repetition rates of hundreds of GHz. Due to their large material gain coefficients, fast recovery times, and the ability to be made into short monolithic cavities, high-repetition-rate pulse trains can be generated easily. Typically, high-repetition-rate lasers involve a monolithic semiconductor laser structure, meaning no external cavity is used. The model presented is intended to analyze the monolithic multisection laser, and no intention of including an external cavity is made here although one could easily modify
were disallowed [see Fig. 2(b)]. One can write an equation for
the net gain of each mode including the coupling effects due
to each of its neighboring modes. Also there are phase effects,
and for stable mode-locking one requires that all the modes
will be equally spaced in frequency. The rest of this section
will be devoted to deriving an equation for each of these
coupled modes which will subsequently be solved to find the
supermode for the high-repetition-rate passively mode-locked
laser.

The net optical field inside the laser can be written as a
sum over individual modes,

\[ \vec{E}(\vec{r}, t) = \sum_n \vec{E}_n(t) \delta_n(\vec{r}) \]

where \( \vec{E}_n(t) \) represents the time dependence of mode \( n \),
and \( \delta_n(\vec{r}) \) represents the \( n \)th spatial eigenmode of the cold
cavity and satisfies \( \nabla^2 \delta_n(\vec{r}) + \mu_0 \epsilon_0 \Omega_n^2 \delta_n(\vec{r}) = 0 \). Here,
\( \mu_0 \) is the magnetic permeability of free space, \( \epsilon_0 \) is the
electric permittivity, and \( \Omega_n \) is the resonant frequency of
the \( n \)th mode of the cold cavity. Assuming we have some
uniform guiding (through index or gain-guiding) structure
longitudinally throughout the laser, we can write

\[ \vec{u}_n(\vec{r}) = \sqrt{2} \delta(x, y) \cos(\beta_n z). \]

These modes of the cold cavity may be delta-function nor-
malized,

\[ \int \delta_n(\vec{r}) \cdot \delta_m(\vec{r}) \, dV = V \delta_{nm}. \]

Similar to (i), the net electronic polarization can be written as
a sum over individual modes,

\[ \vec{P}_n(t) = \frac{1}{V_c} \mu_0 \frac{d}{dt} \int \vec{E}(\vec{r}, t) \cdot \delta_n(\vec{r}) \, dV \]

where \( \vec{P}_n(t) = (1/V_c) \int \vec{P}(\vec{r}, t) \cdot \delta_n(\vec{r}) \, dV \) is the projection of the polarization
on mode \( n \). Here \( \tau_{pn} \) represents the photon lifetime for the \( n \)th mode.

With the optical frequency much greater than the repetition
rate, we may write \( \vec{E}_n(t) \) as the slowly varying complex
envelope of \( \vec{E}_n(t) \) such that

\[ \vec{E}_n(t) = \frac{1}{2} \vec{E}_n(t) e^{i \omega_n t} + c.c., \]

where \( \omega_n \) is the optical angular frequency of the \( n \)th lasing
mode (\( \omega_n \neq \Omega_n \) for nonzero detuning), and correspondingly
\( \vec{P}_n(t) \) may be written as the slowly varying complex envelope
of the polarization. Thus,

\[ \frac{d\vec{E}_n(t)}{dt} - i(\Omega_n - \omega_n) \vec{E}_n(t) + \frac{1}{2\tau_{pn}} \vec{P}_n(t) = \]

where \( \tau_{pn} = \frac{\mu_0}{\epsilon_0} \).
where \( \tilde{P}_n(t) \) will contain coupling terms to electric fields spaced at harmonics of the repetition rate, \( \Delta = \omega_n - \omega_{n-1} \), since the net polarization is given by

\[
\bar{P}(r, t) = \varepsilon_0 \chi(r, t) \bar{E}(r, t),
\]

where

\[
\chi(r, t) = f(\omega_n) \chi_0(r) + \sum_{k \neq 0} f(\omega_{n+k}) \chi_k(r) \cdot \cos[k \Delta t + \psi_k(r)]
\]

may possess optical-pulse-induced oscillations in the carrier density [8], and \( f(\omega_n) \) takes into account the frequency dependent gain or loss of the material. Although in general the material's loss spectrum has somewhat different center and shape than that of the material's gain spectrum, we shall not attempt to model that in this paper.

Because lasers tend to operate at their gain peak and semiconductor lasers have a significant contribution of gain-dependent phase shift at their gain peak, \( \chi_k(r) = \chi_k^G(r) + i \chi_k^L(r) \) presents not only a gain, but a change in refractive index as well. The mode-locked laser is in fact no better in this respect. It tends to operate at an even longer wavelength than a continuous wave (CW) laser (due to the presence of the absorber) [8] and is expected to produce even a slightly larger amplitude-to-phase coupling factor \( \alpha \) in its gain section [15], where \( \alpha = -\chi_k^G(r)/\chi_k^L(r) \).

Since \( \tilde{P}_n(t) \) is computed from a projection of \( \bar{P}(r, t) \) onto \( \tilde{u}_n(\bar{r}) \) over the whole length of the laser, there is a contribution from both the gain and absorbing sections

\[
\tilde{P}_n(t) = f(\omega_n) \frac{\varepsilon_0}{V_c} \xi_0 \tilde{E}_n(t)
+ \frac{\varepsilon_0}{V_c} \sum_{k \neq 0} \left[ f(\omega_{n-k})(\xi_{kg} e^{-i \omega_{kg} t} + \xi_{k0} e^{-i \omega_{k0} t}) \cdot \tilde{E}_{n-k}(t) + f(\omega_{n+k})(\xi_{k0} e^{-i \omega_{k0} t} + \xi_{kg} e^{-i \omega_{kg} t}) \cdot \tilde{E}_{n+k}(t) \right],
\]

where

\[
\xi_0 = \int \chi_0(\bar{r}) |\tilde{u}_n(\bar{r})|^2 dV
\]

and

\[
\xi_{(g/a), \pm} e^{i \psi_{(g/a), \pm}} = \int \chi_{(g/a)}(\bar{r}) \tilde{u}_{n \pm k}(\bar{r}) \cdot \tilde{u}_n(\bar{r}) dV.
\]

(The notation \((g/a)\) indicates quantities pertaining to the gain or absorber region, respectively.) We will use \( \chi' = \tilde{g}_H c/\omega_0 \) to relate the material gain coefficient, \( \tilde{g}_H \), to the imaginary part of the susceptibility, \( \mu \), being the cold cavity refractive index.

The imaginary part of (10) yields

\[
f(\omega_n) \frac{\varepsilon_0}{V_c} = \frac{f(\omega_n) \Gamma \mu_\nu c}{\omega_0} [\tilde{g}_H + \tilde{a} h_a],
\]

and this term is proportional to the average single pass gain where \( \Gamma \) is the confinement factor, \( c \) is the speed of light, \( \tilde{g}_H \) is the material gain coefficient of the absorber section \((\tilde{a} < 0)\), \( L \) is the total laser length, and \( h_g \) and \( h_a \) are the ratios of the gain section and absorber section lengths, respectively, to the full laser length. The term \( \xi_{kg} e^{-i \omega_{kg} t} \) will be determined from the carrier dynamics by using a linear approximation for the change in optical gain (loss) versus carrier density for the gain (absorber) section with \( \tilde{g} = G(n_a(t) - n_0) \). Here, \( G \) is the differential gain, \( n_a(t) \) is the time dependent carrier density, and \( n_0 \) is the carrier density at transparency. Correspondingly, \( \tilde{a} = A[n_a(t) - n_0] \) for the absorber.

Gain and absorber dynamics result from the photon intensity, which is proportional to

\[
S(z, t) = s_0 + \sum_k s_k(z) \cos(k \Delta t)
\]

with

\[
s_0 = \langle s_0(z) \rangle = \frac{1}{2} \sum_{n=-q}^q |\tilde{E}_n(t)|^2 n_0^2(z)
\]

and

\[
s_k(z) = s_k \cos \left( \frac{k \pi z}{L} \right)
= \left( \sum_{n=k-q}^q \tilde{E}_n(t) \tilde{E}_{n-k}(t) u_n(z) u_{n-k}(z) \right) \lambda,
\]

where \( \langle \rangle \lambda \) represents a spatial average over a wavelength. Permitting this form, one notices from the carrier rate equation that a modulation in the light intensity will induce a modulation in the carrier density at the same fundamental and harmonics of that frequency. However, the effect of both the small number of modes and the shorter in-phase overlaps of quickly beating pairs of modes causes the coupling of higher harmonics to drop off. Ignoring the terms responsible for second nearest neighbor and higher coupling terms to simplify the problem and still keep it suitable for high-repetition rate modellocking, from the carrier rate equation

\[
\frac{dn_{(g/a)}(z, t)}{dt} = R_{p(g/a)} - (\tilde{g}/\tilde{a}) \langle z, t \rangle S(z, t)
- \frac{n_{(g/a)}(z, t)}{\tau_{(g/a)}},
\]

we find a saturated material gain \( \tilde{g} \) for the gain section, dependent on the gain recombination time \( \tau_g \) and the injection pumping \( P_{p(g/a)} \) and correspondingly \( \tilde{a} \) for the absorber section dependent on the absorber recombination time \( \tau_a \),

\[
\tilde{g} = \frac{\tilde{g}'}{1 + G \tau_g s_0}
\]

\[
\tilde{a} = \frac{\tilde{a}'}{1 + A \tau_a s_0},
\]

where \( \tilde{g}' \) and \( \tilde{a}' \) represent the unsaturated gain and unsaturated loss. Additionally, the carrier density is written \( n_{(g/a)}(z, t) = \frac{\varepsilon_0}{V_c} \xi_0 n_{(g/a)}(z, t) \).
\[ n_0(g/a) + n_1(g/a)(z) \cos \left[ \Delta \omega + \psi(g/a) \right] + \cdots, \]
and terms showing modulation at the first harmonic in the rate equation lead to:

\[ \tilde{n}_{g1}(z) = \frac{-\tilde{g}_{g1}(z)}{\sqrt{\Delta^2 + \left( \frac{1}{\tau_g} + Gs_0 \right)^2}}, \]

\[ \psi_g = -\arctan \left( \frac{\Delta}{\frac{1}{\tau_g} + Gs_0} \right), \]

\[ \tilde{n}_{a1}(z) = \frac{-\tilde{a}_{a1}(z)}{\sqrt{\Delta^2 + \left( \frac{1}{\tau_a} + As_0 \right)^2}}, \]

\[ \psi_a = -\arctan \left( \frac{\Delta}{\frac{1}{\tau_a} + As_0} \right). \]

So the carrier modulation becomes small and it lags the optical pulses by nearly \( \pi/2 \) radians since the repetition rate is well beyond the recombination rate or saturation rate. Computing the spatial integrals in (11), we find that

\[ \frac{\xi_{g1} e^{i\psi_g}}{V_c} = \left( -\alpha_g + i \right) \frac{\mu_v^2}{\omega_0 \tau_p} \kappa_g s_1 e^{i\psi_g}, \]

where

\[ \kappa_g = \frac{-G\tilde{g}_0}{\sqrt{\Delta^2 + \left( \frac{1}{\tau_g} + Gs_0 \right)^2}} \cdot \left\{ \frac{1}{2} + \frac{1}{4\pi h_g} \sin \left( 2\pi h_g \right) \right\} \]

and

\[ \tilde{g}_0 = \Gamma \tilde{g}_0 \frac{c n_p}{\mu_v} \]

is the normalized gain. Likewise for the part of the integral over the absorber,

\[ \frac{\xi_{a1} e^{i\psi_a}}{V_c} = \left( -\alpha_a + i \right) \frac{\mu_v^2}{\omega_0 \tau_p} \kappa_a s_1 e^{i\psi_a}, \]

where

\[ \kappa_a = \frac{-A\tilde{a}_0}{\sqrt{\Delta^2 + \left( \frac{1}{\tau_a} + As_0 \right)^2}} \cdot \left\{ \frac{1}{2} + \frac{1}{4\pi h_a} \sin \left( 2\pi h_a \right) \right\} \]

and

\[ \tilde{a}_0 = \Gamma \tilde{a}_0 \frac{c n_p}{\mu_v}. \]

One can write the single pass gain from (12) along with its corresponding phase contribution. Also, for generality, one should allow the inclusion of a frequency dependence [16] of \( \alpha \) [accomplished through \( \alpha(g/a) \n_0 \)], giving the single pass net gain and phase effects that are not due to coupling as

\[ \frac{\xi_{g0}}{V_c} = f(\omega_n) \frac{\mu_v^2}{\omega_0 \tau_p} \left[ \left( -\alpha_{gn} + i \right) \tilde{g}_0 + \left( -\alpha_{an} + i \right) \tilde{a}_0 \right]. \]

For steady state, we can ignore all time derivatives and using (6), (9), (23), (25), (26), (28), and (29), the equation for mode \( n \) becomes,

\[ \{ 2i\tau_{nm} (\Omega_n - \omega_n) - 1 + f(\omega_m) \left( \left( 1 + i\alpha_{gn} \right) \tilde{g}_0 + \left( 1 + i\alpha_{an} \right) \tilde{a}_0 \right) \} \tilde{E}_n \]

\[ + \sum_{n=1}^{\infty} \left( \eta_n - \tilde{E}_{n-1} - \tilde{E}_{n+1} \right) \tilde{E}_n = 0, \]

where we have defined coupling coefficients for the nearest neighbor modes,

\[ \tilde{h}_{n+} = f(\omega_n) \left[ \kappa_g \left( 1 + i\alpha_{gn} \right) e^{-i\psi_g} + \kappa_a \left( 1 + i\alpha_{an} \right) e^{-i\psi_a} \right] \frac{\tilde{g}_0}{2}, \]

\[ \tilde{h}_{n-} = f(\omega_n) \left[ \kappa_g \left( 1 + i\alpha_{gn} \right) e^{i\psi_g} + \kappa_a \left( 1 + i\alpha_{an} \right) e^{i\psi_a} \right] \frac{\tilde{g}_0}{2}. \]

These two terms are completely determined by the structure of the laser and the average photon intensity.

Let a single detuning in the separation of modes be defined, \( \delta = \omega_n - \Omega_n - (\omega_{n-1} - \Omega_{n-1}) \), since for stable mode locking the detuning of the repetition rate, \( \delta \), must equal the detuning in the separation between all neighboring modes. The detuning, \( \delta_n \), of mode \( n \) with respect to \( \Omega_n \) is then the detuning of the zeroth mode plus \( n \) times the repetition rate detuning, \( \delta_n = \delta_0 + n\delta \). The general equation then for the \( n \)th mode with nearest neighbor coupling, for \( \alpha \) parameters incorporated for the gain and absorber, and with geometric overlap factors included is

\[ \left\{ -2i\tau_{nm} (\delta_0 + n\delta) + (1 + i\alpha_{gn}) \tilde{g}_0 + (1 + i\alpha_{an}) \tilde{a}_0 - 1 \right\} \tilde{E}_n + \sum_{n=1}^{\infty} \left( \eta_n - \tilde{E}_{n-1} - \tilde{E}_{n+1} \right) \tilde{E}_n = 0. \]

Here \( s_1 = s_1/s_0 \), and the material gain bandwidth is taken into account with \( \tilde{g}_0 = f(\omega_n) \frac{\tilde{g}_0}{\eta } \), and \( \tilde{a}_0 = f(\omega_n) \frac{\tilde{a}_0}{\eta } \).

III. THE SOLUTION

The coupled nonlinear equations (33) can be solved systematically. Also, one should solve the problem for a large enough number of equations such that the result does not depend strongly on the fact that the modes beyond those considered have been forced to have an electric field of zero. To reduce the number of parameters for the calculation, it will be helpful to transform to dimensionless parameters,

\[ s = \frac{A}{G}, \]

\[ r = \frac{\tau_a}{\tau_g}, \]

\[ \Delta = \Delta \tau_g, \]

\[ \tilde{\delta}_0 = G \tau_g \tilde{g}_0, \]

\[ \tilde{\kappa}_g = \frac{\kappa_g}{G \tau_g}, \]

\[ \tilde{\kappa}_a = \frac{\kappa_a}{A \tau_a}. \]
One may subtract out the detuning of mode zero from the set of equations (33). Defining a constant,
\[ R = (i\eta_0 - \tilde{E}_{-1} + \tilde{\eta}_0 + \tilde{E}_1) \frac{\delta_1}{\tilde{E}_0}, \]
(40)
\[ \text{Im} (R) \]
is the component of detuning of the center mode due to mode coupling and \( \text{Re} (R) \) is the reduction in required average gain for the center mode due to mode coupling, similar to that discussed in [17]. Taking the imaginary part of the \( n = 0 \) equation and subtracting it from the general mode \( n \) equation leads to
\[ -2i \tau_{\gamma n} n \delta + (1 + i \alpha_{\gamma n}) \tilde{\eta}_n + (1 + i \alpha_{\gamma 0}) \tilde{\eta}_0 - i \text{Im} (R) - 1) \tilde{E}_n + \delta_1 (\tilde{\eta}_n - \tilde{E}_{n-1} + \tilde{\eta}_n + \tilde{E}_{n+1}) = 0. \]
(41)
The net gain spectrum of the semiconductor material is concave downward and may be represented by the form
\[ f(uT_1) = \frac{1}{1 + (uT_1 - \omega_0)^2/(\Delta \omega)^2}. \]
Since to second order, one may write
\[ f(uT_1) = 1 - b \delta^2, \]
substituting this, and since \( b \ll 1 \) and the coupling term is of the same order, we may ignore their product which goes like \( b^2 \). Now the general equation for mode \( n \) with center mode detuning subtracted finally becomes
\[ -2i \tau_{\gamma n} n \delta + (\tilde{\eta}_n + \tilde{\eta}_0)(1 - b \delta^2) - i \text{Im} (R) - 1) \tilde{E}_n + \delta_1 (\tilde{\eta}_n - \tilde{E}_{n-1} + \tilde{\eta}_n + \tilde{E}_{n+1}) = 0. \]
(42)
Considering a set of \( 2q + 1 \) modal equations (all are complex except for the \( n = 0 \) equation), there are \( 4q + 1 \) real equations and a list of \( 4q + 1 \) unknowns including \( 4q + 1 \) unknowns to specify the fields [we may take \( \text{arg} (\tilde{E}_0) = 0 \) to define an absolute optical phase] and two other unknowns, \( \tilde{\eta}_0 \) and \( \delta \). The phase of the repetition rate is also a degree of freedom and one may specify \( \text{arg} (\delta) = 0 \). Then, the modulation response of the laser sections can be referenced relative to the phase of the optical pulses. Since physically one considers a laser operating with a specific dc pumping (or more appropriately here, a constant average output) power, one may specify a particular average cavity photon intensity for \( \tilde{\eta}_0 \). The latter two conditions, without loss of generality, reduce the number of unknowns in the field vector to \( 4q - 1 \), making the problem completely determined. Due to the nonlinear dependence of the parameters \( \tilde{\eta}_0, \tilde{\eta}_n, \delta, \text{Im} (R), \eta_0^*, \) and \( \delta_1 \) on the vector \( \tilde{E} \), the problem remains challenging. However, the solution is vastly simplified by viewing it as an eigenvalue problem. For example, one may directly write the problem in a matrix form as (43) found at the bottom of the page. Through multiplications of the rows by the appropriate complex factors one may also show that the problem can always be written, having a single complex eigenvalue, \( \lambda \), in the form
\[ [A^m(\tilde{\eta}_0, \delta, \tilde{E}) - \lambda \tilde{E}^m = 0, \]
(44)
where \( A^m(\tilde{\eta}_0, \delta, \tilde{E}) \) is a modified complex matrix and \( \tilde{E}^m \) is a modified eigenvector. The problem is more easily solved by keeping it in the form of (43), however. For a nontrivial eigensolution, we require that the real and imaginary parts of the determinant of the matrix in (43) equal zero. This gives two conditions from which one may find a best estimate for \( \tilde{\eta}_0 \) and \( \delta \), and this was done simply through Newton’s method. With this better estimate of the eigenvalue we proceed to update the relevant parameters and find a new estimate of the eigenvalue. The process is repeated as shown in Fig. 3 until convergence is reached. The computation gives the supermode solutions of the high-repetition-rate laser for the chosen average operating power \( \tilde{\eta}_0 \).

IV. RESULTS FROM AN APPROXIMATE THREE-MODE SOLUTION

The full numerical solution is complicated, involving a large number of interrelated parameters, and it does not quickly
lead to a simple intuitive picture of the effects of the device parameters. To supplement the full numerical solution, an approximate analytical description involving only three modes and an approximation of the supermode symmetry is pursued. One may show that if the $\alpha$ parameters of the gain and absorber sections are ignored, and the gain bandwidth is symmetric relative to the cavity modes, a totally symmetric (odd symmetry) supermode solution for any number of modes will result. The form of the supermode solution will be

$$\tilde{E}_n = \tilde{E}_{-n}^*,$$

and one can always find a three-mode solution having all three modes exactly in phase. However, as soon as $\alpha_g \neq 0$ or $\alpha_a \neq 0$ is chosen, the symmetry is broken and one finds that now a chirp-free supermode solution of this form will not generally exist.

Thus, no passively mode-locked supermode will exist having the form of (45) when the amplitude-to-phase coupling is taken into account. The relative phases of the modes in the supermode depend strongly on the amplitude-to-phase coupling. One finds, for numerous solutions of the full numerical analysis that once a nonzero $\alpha$ parameter is chosen, the solutions are of the even symmetric form

$$\tilde{E}_n \approx \tilde{E}_{-n},$$

since phase effects resulting from the $\alpha$ parameter greatly outweigh the effects present when the $\alpha$'s were zero.

Since a simple, analytic, and reasonably accurate result can be obtained assuming (46) when some nonzero $\alpha$ is present, we derive a solution for three mode-locked modes using this even symmetric assumption. The term $-ib(\alpha_g(\pm 1)\tilde{\eta}_0 + \alpha_a(\pm 1)\tilde{a}_0)$ is found to have little effect on the net gain, amplitudes, or phases of the supermode and will, for this reason, be ignored in this three-mode approximation.

From the $n = 1$ and $n = -1$ equations of (42), the expressions

$$\delta = \frac{\text{Im}(R)}{2\tau_p \text{Re}(\tilde{\eta}_0) + \tilde{\eta}_0},$$

and

$$\tilde{E}_1 = \frac{i\delta(\tilde{\eta}_0 + \tilde{a}_0 - \tilde{\eta}_0 - \tilde{a}_0)}{4\tau_p \delta}$$

can be obtained. Combining this with the $n = 0$ equation, we can find the reduction in required gain for the center mode,

$$\text{Re}(R) = -\tilde{\eta}_0 + \tilde{a}_0 - 1 = -\text{Im}(R) \frac{\text{Im}(\tilde{\eta}_0^2 - \tilde{\eta}_0^2)}{\text{Re}(\tilde{\eta}_0^2 - \tilde{\eta}_0^2)}.$$ (49)

From (48), we will find that a chirp-free solution will exist if

$$\tilde{\kappa}_g \alpha_g \sin \psi_g = -\tilde{\kappa}_a \alpha_a \sin \psi_a.$$

In this case, a soliton-like compensation effect occurs in the monolithic laser cavity. This condition implies that the self-phase modulation (SPM) of the absorber section may exactly oppose the SPM from the gain section [18], [19]. For a larger ratio of $\alpha_g : \alpha_a$, a net upchirp (optical frequency rising with time during the pulse) due to SPM will occur. In the frequency domain picture this corresponds to a phase term, $e^{i\Delta(\omega - \omega)^2}$, multiplying the optical spectrum, where $\alpha$ is negative. For a smaller ratio of $\alpha_g : \alpha_a$, a net downchirp due to SPM is found to occur. A plot of chirp verses the ratio of $\alpha_g : \alpha_a$ for a specific laser operating point will be shown in the next section, using the full calculation. Evidence of both these regimes has recently been demonstrated [20].

### V. THE FULL SUPERMODE CALCULATION

As formulated in Section III, the high-repetition-rate laser supermode can be found numerically. This may be accomplished even while eliminating all assumptions on the modal phase and removing any restrictions on the number of participating Fabry–Perot modes. One finds that if a large enough number of modes is allowed such that the outermost modes have powers of $<10^{-6}$ compared to the strongest modes, there is little further change in the result if additional modes are included.

Given reasonable parameters for laser material and structure, such as those shown in Table I, one can find the supermode solution. In general, one would not expect the $\alpha$ parameter from the gain and absorber regions to be equal. Previously [15], the dependence of the interband transition component of this parameter has been calculated. One would expect a smaller $\alpha$ parameter for laser sections pumped to lower carrier densities. This, in fact, is found to be an important consideration in finding a stable supermode solution. Lau [8] has calculated supermode solutions for three modes with $\alpha = 0$ for both sections. We find reasonably good qualitative agreement with these results even as the number of modes considered is increased. The plots resulting from $\alpha = 0$,

<table>
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<tr>
<th>Variable</th>
<th>Symbol</th>
<th>Value</th>
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<td>Number of Modes Considered</td>
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<td>Center Wavelength</td>
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<td>$\mu$m</td>
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Fig. 4. Calculated mode structure of supermode assuming no amplitude-to-phase coupling, $\alpha_g = 0$, $\alpha_a = 0$, $\delta_0 = 2.5$, and using the other parameter values as given in Table I.

$\delta_0 = 2.5$, a 15-mode calculation, and the parameters in Table I are shown in Figs. 4 and 5. From here on, the frequency dependence of the cold cavity loss is neglected so $\tau_{pn} = \tau_p$.

Fig. 4 shows the calculated field strengths for the 15-mode supermode. Fig. 5 shows the corresponding modal phases, where $\phi_n$ is defined as the optical phase in $E_n e^{i(\omega nt + \phi_n)}$.

Clearly, the symmetry of (45) is present here. Fig. 6 shows the threshold gain difference, $Re (R)$, as a function of $\delta_0$, defined previously and is displayed in units of $10^{-4}$ times the cold cavity loss (from $\tau_p$). The right side scale of this plot shows the expected detuning, $\delta$, of the cavity repetition rate. Fig. 7 shows the modulation depth at the first harmonic as a function of average intensity. The threshold gain for single mode operation must be greater than the mode-locking threshold gain, meaning $Re (R) > 0$ for stable modelocking to be realized [8]. In ideal amplitude modulated (AM) passive modelocking, a minimum mode coupling is required in order to obtain simultaneous lasing from 3 or more modes of a homogeneously broadened laser. This requires a minimum nonlinearity to be present. Hence, if the average cavity intensity, $\delta_0$, is too low, an inadequate amount of mode coupling is generated, and mode-locked operation cannot be obtained. Additionally, if the cavity intensity is large such that the absorber is strongly saturated to a point far beyond the knee of the nonlinearity, the minimum mode coupling again cannot be obtained. This explains why mode-locking may only be obtained over a finite range in Fig. 6. The right scale in Fig. 7 shows in this case where $\alpha_g = 0$ and $\alpha_a = 0$, one does not expect SPM to generate any pulse chirp effects and the quadratic phase $\phi_1 + \phi_{-1} - 2\phi_0)/2 = 0$ indicates that, to first order, no linear chirp is present in this case.

As discussed in the previous section, the $\alpha$ parameter can have a large effect on the phase of each optical mode. Assuming an $\alpha$ parameter of $\alpha_g = 4$ for the gain section, only a limited range of values for $\alpha_a$, the $\alpha$ parameter for the absorber section, was found to give stable self-consistent solutions. A calculation of the approximately linear chirp (quadratic phase) at the center of the optical spectrum, $(\phi_1 + \phi_{-1} - 2\phi_0)/2$, versus $\alpha_a/\alpha_g$ is plotted in Fig. 8 for the range of stable mode-locked solutions. The range is quite narrow and corresponds to a region where the SPM effects from absorber and gain nearly cancel as discussed in [20]. The dependence of the $\alpha$'s on frequency is ignored in this and subsequent plots.

The same plots as shown in Figs. 4–7 can be shown for the case including effects of reasonable nonzero $\alpha$'s. The new calculated field strength for the 15-mode supermode with...
Fig. 8. A calculation of the linear chirp, the quadratic phase around the center of the optical spectrum, for different values of $\alpha / \alpha_g$.

Fig. 9. Calculated mode structure of supermode when allowing amplitude-to-phase coupling, $\alpha_g = 4$, $\alpha_a = 2.1$, and $\delta_0 = 2.5$. Parameters are exactly as shown in Table I.

$\alpha_g = 4$ and $\alpha_a = 2.1$ is shown in Fig. 9. Fig. 10 shows the corresponding modal phases, $\phi_n$. The previously discussed change in supermode symmetry is mainly shown in this plot of $\phi_n$. Before discussing the other three plots, it should be mentioned that physically as the gain current in the laser is increased to raise the average intensity, $\delta_0$, one weakens the absorber section through the relation $\delta_0 = \delta_0^0/(1 + \tau s \delta_0)$, where $\delta_0^0$ is the section’s normalized unsaturated absorption. The strength of the gain is also weakened since we require $\delta_0 + \delta_0 - 1 \approx 0$. Thus, the two sections both operate closer to transparency as $\delta_0$ is increased. This implies a change in each section’s $\alpha$ parameter also occurs and their dependence on $\delta_0$ will be approximated to first order here by $\Delta \alpha_g = -\alpha_1 \Delta \delta_0$ and $\Delta \alpha_a = \alpha_1 \Delta \delta_0$, where $\alpha_1$ takes into account a linear decrease (increase) in $\alpha_g (\alpha_a)$ as the cavity intensity is increased. Here, $\alpha_1$ is taken as 0.25 around the point $\delta_0 = 2.5$.

Fig. 11 shows the plot of required gain reduction, $\text{Re}(R)$, and the expected detuning in the repetition rate as a function of $\delta_0$. Fig. 12 shows the modulation depth and an estimate of the mode-locked laser’s linear chirp $(\phi_1 + \phi_{-1} - 2\phi_0)/2$. One can see that the expected mode-locking range over which the coupled equations can be simultaneously satisfied is severely limited when the phase condition including the $\alpha$ parameter is considered (although other parameters remain identical). This is a direct result of the presence of the $\alpha$ parameters in the coupling terms and occurs consistently regardless of whether or not one includes more allowed modes in the calculation.

It is expected that the mode-locked laser’s operation will change if one modifies the structure or bias parameters. These
effects are important if one intends to understand or optimize the laser's operation. We have calculated results one would expect from modifying key laser parameters and using the nearest-neighbor mode coupling approximation for the range of supermode solutions that exist around the case considered in Fig. 9.

One finds that if $s$, the ratio of the differential absorption to differential gain is increased, a larger mode coupling is obtained. This leads to a larger value of $R_e$, the reduction in the mode-locking threshold relative to the single-mode threshold, as shown in Fig. 13, which is expected to lead to a more stable mode-locked supermode. Fig. 13 also shows that a decreased upchirp or increased downchirp is expected to occur if a larger $s$ is present and all other parameters are unchanged.

The effect of $r$, the ratio of absorber recovery time to recovery time, is expected to be nearly the opposite. Shown in Fig. 14, an increased $r$ leads to a decrease in the mode-locked gain reduction and ultimately a loss of a stable mode-locked solution altogether as the ratio is increased above $r = 0.46$ in this case. Simultaneously the increased value of $r$ will lead to an increased upchirp as shown in Fig. 14. It is known that one can reduce the value of $r$ through stronger reverse bias or ion implantation into the absorber section.

An increase in $\tilde{\alpha}_u$, the unsaturated absorption strength of the saturable absorber, is shown to lead (as shown in Fig. 15) to an increase in the mode-locked gain reduction, $R_e$. The strength of the unsaturated absorption is proportional to this section's length and absorption coefficient. An increase in either of these is expected to lead to a more strongly downchirped pulse as shown in Fig. 15. This agrees with expectations described in [20] where a stronger saturable absorber is cited as the reason for a significant downchirp being obtained over most of the experimental chirp-versus-current curve.

Unlike the time-domain analyses [1]–[3], the frequency domain analysis allows one to account for the spatial geometry of the laser. Recently, this has been noted by Martins-Filho et al. [13]. The ratio of the physical length of the absorber to the total laser length, $h_a$, is expected to change the effectiveness of mode coupling. If the same absorber strength can be incorporated into a smaller segment of the laser, one can achieve a more effective mode coupling and obtain a larger mode-locked gain reduction, $R_e$. This is consistent with results determined in [21]. Fig. 16 also shows the effects on pulse chirp when the parameter $h_a$ is varied.

The mode-locked laser's round-trip frequency is determined by the laser's cavity length. A larger cavity round-trip fre-
frequency is expected to result in reduced mode coupling due to a reduction in $k_0$ and $\kappa_0$. This will eventually lead to a point where the minimum mode coupling cannot be obtained and no stable mode-locked supermode exists. Although the point is $\approx 105$ GHz in this case (Fig. 17), using larger values of $s$ ($s \geq 5$), we have obtained stable supermode solutions slightly beyond 200 GHz. This agrees well with the theoretical results presented by Lau [8]. In this case, larger mode coupling effects resulted in a reduced downchirp.

As intuitively expected, lasers having a larger gain and absorber bandwidth will obtain a greater mode-locked gain reduction. Fig. 18 shows the expected increase in Re($\mathcal{R}$) as one solves the supermode equations allowing successively larger material bandwidths. Even larger advantages are found to occur if one assumes a gain bandwidth wider than the laser’s absorption bandwidth. An expected decreased pulse chirp for larger material bandwidth is also shown in Fig. 18.

VI. EXPERIMENTAL RESULTS AND DISCUSSION

Experimental measurements of the spectrum, pulse chirp, and the variation of pulse chirp with injection current have previously been published [20]. The laser used was a monolithic two-section quadruple quantum well GaAs laser having a repetition rate of 73 GHz. It showed qualitatively the same characteristics as the calculation for 80 GHz in the previous section. A broader spectrum and longer pulses as found from streak camera results can typically be obtained at higher bias conditions. An optical spectrum for the laser operating at 30 mA has previously been shown in [20]. The chirp of this spectrum has been measured through cross-correlation techniques [20], and integration of these results leads to phase values, $\phi(\lambda)$, of the optical spectrum plotted in Fig. 19. The figure shows a phase of the optical spectrum corresponding to a train of pulses with a 1.7 ps/nm downchirp and a time-bandwidth product, $\Delta T \Delta \nu$, which is 18% larger than the compressed pulse time-bandwidth product achieved in the experiment. This regime of operation qualitatively corresponds to the calculated optical phase in Fig. 10. Additionally, the experimental measurement of 1.7 ps/nm downchirp for this laser is equivalent to $(\phi_1 + \phi_{-1} - 2\phi_0)/2 = 0.07$ rad.

Previously presented experimental results have demonstrated the effect of changes in the dc gain section injection current from the preceding case. The experimental results are shown in Fig. 20. While the laser is above threshold, the changes in dc injection current are nearly linearly related to the average photon intensity inside the cavity, $g_0$. Hence, we expect Fig. 20 to show agreement with the calculated pulse chirp in Fig. 12. Both show a sequence of upchirped, chirp-free, and downchirped operation as the photon intensity inside the cavity is increased.

Although good agreement between theory and experiment is obtained, we do not intend to imply that we have found the actual parameters of the mode-locked laser. However, we believe that the chosen parameters place the calculation in qualitatively the same regime of operation and that the calculated effects of the $\alpha$ parameter, the laser structure parameters, and the bias parameters will show a good correspondence with further experimental results.

Additionally, all results presented in this paper are believed to be for the lowest-order supermode—the one which possesses a minimum threshold gain. We have found some relatively small regions in the parameter space in which a second supermode solution could be found as a self-consistent solution. For a set of reasonable parameters and an arbitrary $\tilde{s}_0$, we have obtained stable supermode solutions slightly beyond 200 GHz. This agrees well with the theoretical results presented by Lau [8]. In this case, larger mode coupling effects resulted in a reduced downchirp.

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and downchirped pulses all from a single laser under different gain section bias. Results of the full supermode calculation (with nearest-neighbor-only coupling) were presented. The supermode magnitude and phase were plotted in the case where no amplitude-to-phase coupling exists in either laser section. Here, supermode solutions could be obtained over a broad range of cavity intensities. In this case, parameter values and mode-locking ranges show good agreement with previous calculations by Lau [8]. Other characteristics of the supermode solution were plotted as a function of cavity intensity also.

When reasonable amplitude-to-phase coupling factors were chosen for both laser sections, the supermode symmetry was severely changed. The phase was found to take on a predominantly quadratic shape in the region of the spectrum where the mode strengths are significant. This indicated the presence of essentially linearly chirped pulses. The presence of a nonzero $\alpha$ parameter was found to drastically limit the range (in terms the variation of cavity intensity) over which stable mode-locked solutions could be found owing to the added phase effects. Near regions where the phase effects from gain and absorber nearly compensated each other, the effect of the $\alpha$ parameters on the reduction in gain due to coupling were neither very advantageous nor very harmful. They typically led to a slight weakening in $\text{Re} \left( R \right)$, the reduction in threshold gain due to mode coupling. To facilitate understanding and optimization of high-repetition-rate passively mode-locked lasers, calculations of the reduction in gain provided through mode coupling and of the expected linear chirp were presented for variations in parameters of the laser structure and bias. Comparisons were made to expectations and to results from other models.

Next, experimental results from a high-repetition-rate passively mode-locked laser at 73 GHz were compared to the supermode calculations in this paper. A good qualitative agreement for the spectral shape, chirp, and variation in chirp with changing injection current was found. The calculated supermodes analyzed were typically not as broad as the measured supermode. The reason for choosing narrower supermodes is that in this case the higher-order coupling effects (e.g., second nearest neighbor, third nearest neighbor coupling, etc.) are expected to be smaller. Thus, in this case, the nearest-neighbor-coupling approximation is expected to be more accurate. However, by including second-nearest-neighbor coupling and higher-order coupling in the matrix for the supermode solution, one may more accurately model lower repetition rate mode-locked lasers down to lower repetition rates ($\leq$5 GHz) which are viable for data rates in communication systems which are practical today.

**REFERENCES**


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Amnon Yariv (S'56-M'59-F'70-LF'95) was born in Tel Aviv, Israel. He received the B.S. (1954), M.S. (1956), and Ph.D. (1958) degrees in electrical engineering from the University of California in Berkeley. In 1959, he went to Bell Telephone Laboratories, Murray Hill, NJ, joining the early stages of the laser effort. He came to the California Institute of Technology in 1964 as an Associate Professor of Electrical Engineering, and became a Professor in 1966. In 1980, he became the Thomas G. Myers Professor of Electrical Engineering and Applied Physics. On the technical side, he took part (with various co-workers) in the discovery of a number of early solid-state laser systems, in proposing and demonstrating the use of semiconductor integrated optics, and the invention of the semiconductor distributed feedback laser and in pioneering the field of phase conjugate optics. His present research efforts are in the areas of nonlinear optics, semiconductor lasers, and integrated optics with emphasis on communication and computation. He is a founder and chairman of the board of the ORTEL Corporation. He has published widely in the laser and optics fields, and has written a number of basic texts in quantum electronics, optics and quantum mechanics.

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