Spontaneous Emission in Semiconductor Laser Amplifiers

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Abstract—In a mode matched configuration, spontaneous emission in semiconductor laser amplifiers is enhanced by a factor which is larger than unity but which is significantly smaller than the K-factor calculated by Petermann. Using thin-slab model, we find that in typical situations, the factor is about $K/2$.

I. Introduction

IN optical communication systems, it may be advisable to amplify weak optical signals before detection. If the laser gain is large enough, the noise-to-signal ratio at the detector output may indeed be as low as

$$N/S = 4hf/P_0 \tag{1}$$

independently of the detector noise (e.g., thermal noise). In (1), S denotes the square of the signal detected current, N is the average square of the detected current fluctuations per unit bandwidth, hf represents the photon energy, and P_0 is the received optical power assumed to be independent of time. The simple result in (1) applies only when the lower level of the laser medium is unpopulated (ground state population density << upper state population density) and the signal frequency is at the line center. Among the various noise terms that can be found at the detector output, we have retained in (1) only the beat term between the signal field and the field spontaneously emitted by the amplifying medium. This is in general a permissible approximation when the laser gain is large. For a more complete discussion, see [1]-[3].

Before we can apply (1), it is necessary to examine in detail the mechanism of guiding and amplification of the optical wave. Indeed, Petermann found a few years ago that the power emitted by spontaneous emission in gain-guided semiconductor lasers is enhanced by a factor K that may be as large as 50 in typical situations [4]. Most authors concluded that gain-guided lasers are intrinsically more noisy than index-guided lasers and that, consequently, they should not be used in optical amplification.

Let us briefly review the Petermann argument. In gain-guided lasers, the real part of the refractive index is smaller in the central region than outside that region, so that the system would guide only leaky modes (that is, modes whose field increases exponentially as a function of the transverse coordinates instead of being confined, as is the case of ordinary guided modes) if it were not for the medium gain. As we increase the medium gain, there is a threshold beyond which the fundamental mode of propagation becomes normally guided with an exponentially decaying field, and the propagating wave experiences a net gain [5]. The wavefront of the mode is curved, however, or in other words, the modal field is a complex function of the transverse coordinates.

On the other hand, it is well known that the medium gain is a stimulated emission process which is always accompanied by a spontaneous emission process. The later can be modeled by an assembly of dipoles radiating independently of each other. Each dipole can be represented by a Dirac δ-function, and expanded in a series of the modes of the guiding structure, which, in the present case, are complex functions of the transverse coordinates. Petermann [4] has then shown that the spontaneous emission power in the fundamental mode, interpreted as the integral of the square modulus of the fundamental modal field spontaneously emitted, is enhanced by a large factor K, compared to the case where the modal fields are real functions of the transverse coordinates. The latter case is often improperly refered to as that of "index-guidance." This is improper because wavefront curvature appears as soon as the gain is larger in the central region than outside, whether the real part of the refractive index increases (gain-guidance) or decreases (index-guidance) as a function of the transverse coordinates. The effect of wavefront curvature is usually stronger in the case of inverted real index profiles, but strong wavefront curvature effects can also be found for normal guidance if the guidance is weak (thin slabs).

Arnaud [6] pointed out that the spontaneous emission power in the mode evaluated by Petermann is not the physically relevant quantity. At the output of the laser amplifier, the spontaneous emission field should be evaluated in a power-orthogonal system of modes. In particular, the quantity that enters into (1) as N is the integral over the detector area of the product of the signal field and the spontaneously emitted field, one of the two being complex conjugated. In that way, one gets a noise enhancement factor K' which is smaller than Petermann's K-factor. If the laser is short, that is, if mode filtering has virtually not taken place, then there is no noise enhancement at all, that is, $K' = 1$ [6]. If, on the contrary, the laser length is large, so that mode filtering is strong and only the fundamental mode of propagation survives at the laser output, and furthermore, if the input field coincides with the fundamental mode, then Arnaud's K'-factor coincides with Petermann's K-factor. The purpose of this paper is to give a precise value for

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the K'-factor in intermediate situations. We will consider typical laser gains such as 10 dB, and typical laser lengths such as 200 μm. The K'-factor calculation is meaningful only for a multi-moded guiding system. In order to make the problem mathematically tractable and obtain closed-form expressions, we have considered as a simple model a reactive surface. A reactive surface is basically free-space with a boundary condition that relates the field and its first derivative at $x = 0$. Specifically, we set $-(dE/dx)/E = s$ and call s the susceptance. Usually, s is a real positive quantity. Then the reactive surface can support one guided mode whose field decays according to an exp $(-s|x|)$ law and a continuum of radiation modes that play the role of higher order modes. A reactive surface can be viewed as a dielectric slab whose thickness goes to 0 but whose product of permittivity and thickness remains finite (Fig. 1).

In order to get gain, we consider complex values of the susceptance s and set $s = a - ib$, $b > 0$ (an exp $(-i\omega t)$ time dependence is used throughout). Such a complex reactive surface again sustains only one guided mode, and a continuum of radiation modes. These modes can be defined from the analytic continuation of the modes in the real case.

The paper is organized as follows. In Section II, we give a general expression for the N/S ratio. In Section III, we give the modal properties of the thin slab. In Section IV, these results are applied to the general expression of Section II, and an almost closed-form expression for the K'-factor is given. In the final section, numerical results are presented and a comparison is made with purely numerical results obtained for realistic laser configurations using the beam propagation method. A brief account of these results has been reported in [7].

After submission of this paper, we were informed of an important paper by Haus and Kawakami [8]. These authors show that for very large laser lengths (were Arnaud's K' is almost equal to Petermann's K), the optimum input field is not the fundamental mode $E_0(x)$ but its complex conjugate $E_0^*(x)$. For the particular input field, the gain is enhanced by the same factor $K >> 1$ as the noise, and therefore, gain-guided laser amplifiers may not be more noisy than index-guided amplifiers. In other words, contrary to intuition, one should *not* try to mode-match the input field. While $E_0(x)$ represents a wave with diverging wavefront at the input facet, the optimum field $E_0^*(x)$ represents a wave with the same amplitude distribution but with an opposite (converging) wavefront. The optimum input field for laser amplifiers of finite length remains to be investigated, perhaps along the lines of this paper.

A word is in order concerning laser *oscillators*. Petermann's paper [4] was in fact specifically oriented toward laser oscillators; even so, his results have been widely interpreted as being also applicable to laser amplifiers. We have extended our thin-slab model to laser oscillators with plane cleaved facets and were able to show analytically that the fundamental mode intensity [from the first term in (15)] is equal to the total optical field intensity [given simply by $|u - is|^{-2}$ in the notation of (15)] well above threshold. This result validates Petermann's introduction of the K-factor, at least with respect to the evolution of the electron density well above threshold. But the situation is not so clear near threshold. The subject of laser oscillators is not addressed in detail in the present paper.

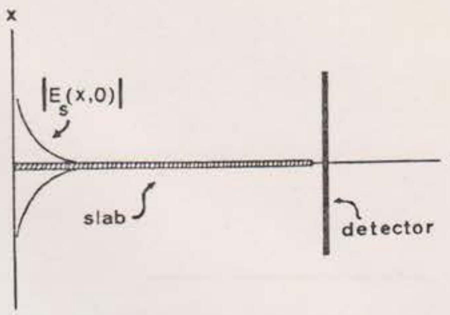

Fig. 1. Thin slab model of a laser amplifier. Input plane: $z = 0$, output or detector plane: $z = L$. The slab is normally guiding and has gain. It is characterized by a susceptance $s = a - ib$, $a > 0$, $b > 0$, $E_s(x, 0)$ represents the input field, supposed to be in the fundamental mode of propagation.

II. General Formulation

The problem of optical amplification is defined as follows. Consider a volume V containing atoms all in the excited state. That is, we assume for the sake of simplicity a complete population inversion. (This condition is not achieved in semiconductor lasers, and the noise is increased by a factor of the order of 2 with respect to the values presently calculated.) Because of stimulated emission, this collection of atoms can be characterized by a complex index of refraction

$$n = n_r + in_i \tag{2}$$

where $-(\omega/c)n_i = g$ represents the local gain, proportional to the population density of the atoms in the upper state. The variation of g with the optical frequency is Lorentzian for free atoms and assumes a more complicated shape for semiconductors. The variation of the real part of the refractive index as a function of the optical frequency is related to that of the imaginary part, plus a constant that may be due to nonresonant states. In this paper, we will assume that the signal optical frequency is at the line center and that the line shape is symmetrical, so that both sides contribute equally to the baseband noise. It is, of course, assumed that there are many atoms per optical wavelengths, so that it is meaningful to characterize the medium by a refractive index. This is so even if very few atoms actually make a transition to the lower state.

For the sake of simplicity, we will ignore polarization effects and assume that the optical field obeys the paraxial wave quation

$$-i\frac{\partial E}{\partial z} = \frac{\omega}{c}\, n(x,y,z)E + \frac{1}{2k_0}\left(\frac{\partial^2 E}{\partial x^2} + \frac{\partial^2 E}{\partial y^2}\right) \tag{3}$$

where $k_0 = (\omega/c)n_r(0, 0, z)$ is supposed to be a real constant. Its imaginary part can indeed by neglected in the last term of the above wave equation.

The active region extends from $z = 0$ to $z = L$. Thus, the power gain can be defined as

$$G = \iint |E_s(x,y,L)|^2\, dx\, dy \Big/ \iint |E_s(x,y,0)|^2\, dx\, dy \tag{4}$$

where the integrals go from $x, y = -\infty$ to $+\infty$. The $E(x, y, z)$ field functions are so normalized that $|E|^2$ is the power density.

The s subscript refers to the input optical signal. If now $E_{sp}(x, y, L)$ denotes the optical field generated by spontaneous emission at the detector plane ($z = L$), the beat between the amplified signal field and the spontaneously emitted field results at the output of the detector in a noise spectral density N given by

$$\frac{N}{S} = 4 \iiint_V \left| \frac{\iint E_s^*(x,y,L)\, E_{sp}(x,y,L)\, dx\, dy}{\iint E_s^*(x,y,L)\, E_s(x,y,L)\, dx\, dy} \right|^2 \tag{5}$$

where $E_s(x, y, L)$ is, as before, the signal field at the detector plane and the denominator in (5) is evidently GP_0. The field at $z = z_0$ spontaneously emitted by a volume $dV = dx_0 dy_0 dz_0$ centered at x_0, y_0, z_0 is

$$(x, y, z_0; x_0, y_0, z_0) = [2hf\, g(x_0, y_0, z_0)\, dV]^{1/2} \cdot \delta(x - x_0)\, \delta(y - y_0). \tag{6}$$

Both E_{sp} and N are relative to a unit bandwidth, optical and baseband, respectively. Note that the stimulated-emission gain g enters into the expression of the spontaneously emitted field under a square root so that the spontaneously emitted power is proportional to g. The field defined at $z = z_0$ by (6) propagates along the z-axis according to the wave equation, (3). In that way, we can obtain the field $E_{sp}(x, y, L)$ that enters into the N/S expression in (5). The last integration is over the active volume V. The differential element dV appears under the square root in (6). It can be shown that (5) reduces to (1) when the medium gain is uniform, that is, when g = constant. This justifies the weighting factors introduced in (6). Note that (6) is valid only for the paraxial approximation. One should not attempt to calculate the total spontaneously emitted power from that expression. The field radiated at large angles is in fact irrelevant to our problem, since we are only concerned with the coupling to paraxial modes.

III. The Thin Slab Model

In a double heterojunction semiconductor laser, guidance of the optical field is ensured by a slab of material of refractive index higher than that of the surrounding medium. The slab thickness is on the order of 0.1 μm and the optical field is strongly confined, with mode-thicknesses of the order of 1 μm. It is recalled that at a typical wavelength of 1.3 μm, the wavelength in the material is on the order of 0.4 μm. Because of the strong optical confinement, the curvature of the wavefront of the fundamental mode of propagation is negligible. Let us recall that the wave gets power only in the central active region, and that this power must flow outside into the gainless outer regions in order to sustain a growing wave. This power, however, does not propagate to infinity. Indeed, it can be shown that as soon as the wave is amplified, the field amplitude decays exponentially in the transverse direction [5].

For reasons that may be considered at the moment academic, we shall consider extremely thin active regions, with thicknesses of the order of 0.01 μm. Such very thin active regions are actually found in quantum-well lasers, but then separate confinement is added in order to increase the gain. Our motivation for considering very thin active regions (without any other optical wave confinement) is that the wave is then weakly guided, wavefront curvature may have strong effects, and large Petermann K-factors appear. In other words, while Petermann introduced the K-factor in connection with gain guidance in the junction plane, we find it more tractable from a mathematical standpoint to treat the case of an active region which is very thin in the direction perpendicular to the junction plane.

As we discussed in the introduction, for small slab thicknesses, a dielectric slab can be replaced by a reactive surface. In the limit where its thickness $2d$ goes to 0 and n^2 goes to infinity in such a way that

$$s \equiv k_0^2 d(n^2 - 1) \equiv V^2/d; \quad k_0 = \omega/c \tag{7}$$

is finite, we simply need to impose upon the field the following boundary condition at $x = 0$:

$$\frac{dE}{dx} + sE = 0. \tag{8}$$

More precisely, $-2s$ is the discontinuity of $(dE/dx)/E$ above and below the reactive surface. But here, only symmetrical solutions need to be considered. It is then well known that the unique guided modal solution is

$$E(x, z) = \exp(-s|x|)\exp(i\beta z)$$
$$\beta = k_0 + \tfrac{1}{2}s^2/k_0 \tag{9}$$

if s is real and positive. The solution in (9) describes a wave propagating along the z axis with propagation constant $\beta > \omega/c$ and exponential decay along the x axis (perpendicular to the junction plane). Now let the slab permittivity n^2 have an imaginary part: $n^2 = \epsilon_r + i\epsilon_i$. The imaginary part $\epsilon_i < 0$ expresses the fact that there is gain in the slab. As a consequence we may set from (7)

$$s = a - ib; \quad a > 0, b > 0. \tag{10}$$

The solution in (9) remains formally applicable. We have

$$E(x, z) = \exp(-a|x|)\exp\{i[b|x| + \beta_r z]\}\exp(\gamma z) \tag{11}$$

where we have set

$$\beta \equiv \beta_r - i\gamma \tag{12}$$

where γ is the gain of the wave. The power gain is

$$G = \exp(2\gamma L), \quad \gamma = ab/k_0 \tag{13}$$

where L denotes the laser amplifier length. Note incidently that the thin slab currently considered can sustain a guided wave only if it is normally guiding ($a > 0$). In order to get guided waves with inverted real-index profiles, we need two reactive surfaces separated by some nonzero spacing, say D. In the present paper, only single normally guiding reactive surfaces will be considered.

It is very easy to calculate Petermann's K-factor with our model. We get

$$K = \left(\int |E^2|\, dx \Big/ \left|\int E^2 dx\right|\right)^2 = 1 + (b/a)^2. \tag{14}$$

Thus, if $b/a = 7$, for example, we get from (14) $K = 50$. However, as we shall see, the effective noise-enhancement factor K' under mode-matched conditions is not given by K but by a more complicated expression involving the radiation modes (or, more generally the higher order modes). It is significantly smaller. The K-factor given by Petermann would be applicable directly to the expression of the signal-to-noise ratio only if there were no complex conjugation of E_s in (5), or if the spontaneous emission power were all in the fundamental mode at the laser output (and the signal field were mode-matched at the input). As a matter of fact, higher order modes (or radiation modes) play a significant role.

Thus, let us consider the general expression of the field radiated by a $\delta(x)$ source located at $z = z_0$ and $x = 0$ in the presence of a reactive surface whose (possibly complex) susceptance is denoted s as before. The field at location x, L is [9]

$$G(x, L, z_0) = se^{-sx} \exp\left[i\left(k_0 + \frac{1}{2} s^2/k_0\right)(L - z_0)\right] + 2\int_0^\infty \frac{u}{\pi(u^2 + s^2)} (u \cos ux - s \sin ux) \cdot \exp\left[i\left(k_0 - \frac{1}{2} u^2/k_0\right) (L - z_0)\right] du. \quad (15)$$

In this equation, x stands for $|x|$, the solution being even in x. The first term in (15) is the field of the guided mode excited by the source and the second term is an integral over the radiation modes. It can be verified that the rhs of (15) satisfies the wave equation, (3), and the boundary condition in (8). Furthermore, the rhs of (15) reduces to a $\delta(x)$ function as $L \to z_0^+$ [the cos term in the integral is one form of a Dirac δ-function, while the sine term cancels out with the exp $(-sx)$ term). Note the orthogonality in the sense of the simple product of the modal functions in (15), even if s is complex, provided the integrals exist.

Let us now consider the field spontaneously emitted by a length dz_0 of the reactive surface. It is given by the rhs of (15) multiplied by the factor

$$[2hf(b/k_0)\, dz_0]^{1/2}. \quad (16)$$

To obtain this factor we have replaced g in (6) with $b/2k_0d$ according to (7), which gives a correspondence between the slab of nonzero thickness and the b-parameter relative to the reactive surface. We have also performed an integration over x_0 from $-d$ to $+d$ and taken the limit $k_0 d \to 0$.

The incident field, on the other hand, is taken to be in the fundamental mode, that is,

$$E_s(x, L) = \sqrt{aP_0}\; e^{-sx} \exp\left[i(k_0 + \tfrac{1}{2} s^2/k_0)L\right] \quad (17)$$

where P_0 denotes as before the incident coherent optical power to be amplified. Note that the wavefront is curved in a diverging manner. Without loss of generality, the phase at $x = z = 0$ is taken to be zero.

When we substitute the expressions in (17) for the signal field and (15) and (16) for the spontaneously emitted field into the general expression in (5), we find that some of the integrations can be performed in closed form.

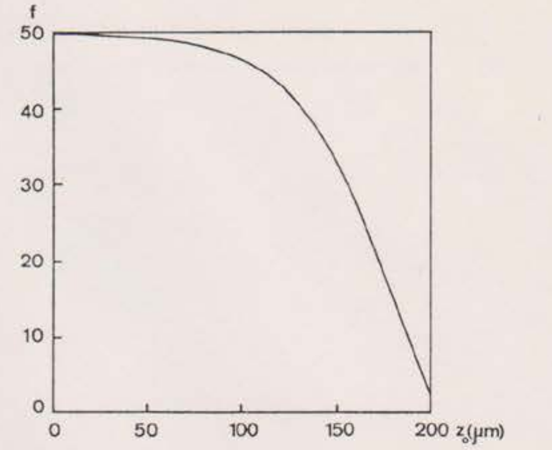

Fig. 2. The function f expresses the noise enhancement factor for spontaneous emission from a location at a distance z_0 from the input plane. $f = 1$ when $z_0 = L = 200$ μm. The slab susceptance (in μm^{-1}) $s = 0.12 - i\,0.84$ and the free-space wavelength is 1 μm.

Using the identity

$$\int_0^\infty (u \cos ux - s \sin ux)\, e^{-s^*x}\, dx = -2ibu/(u^2 + s^{*2}) \quad (18)$$

where the star denotes as before a complex conjugation, we arrive at the following expression for the effective noise enhancement factor:

$$K' = \int_0^L 2\gamma e^{-2\gamma z_0} f(z_0)\, dz_0 \quad (19)$$

where

$$f(z_0) = K \left| 1 - \frac{4iab}{\pi s} e^{-i\beta(L - z_0)} \int_0^\infty \frac{u^2 \exp\left[i\left(k_0 - \frac{1}{2}\frac{u^2}{k_0}\right)(L - z_0)\right]}{(u^2 + s^2)(u^2 + s^{*2})} du \right|^2 \quad (20)$$

is equal to 1 when $z_0 = L$, and $K = 1 + (b/a)^2$ when $L - z_0$ is large.

$$\beta = k_0 + \tfrac{1}{2} s^2/k_0 \quad (21)$$

is the propagation constant of the fundamental mode.

IV. Numerical Results

We have plotted in Fig. 2 the function $f(z_0)$ given in (20) for a set of the a and b parameters. As expected, for that part of the spontaneously emitted field that originates near the output end of the laser amplifier, there is no noise enhancement at all: $f = 1$, because the spontaneously emitted field essentially remains a δ-function. On the contrary, for that part of the spontaneously emitted field that originates from the input end of the laser amplifier, the spontaneously emitted field is essentially in the fundamental mode at the output and Petermann's factor applies: $f \cong K$. What matters in a laser amplifier is, however, the total noise emitted by the active region from $z = 0$ to $z = L$. This average quantity K' (or rather K'/K) is given in Fig. 3, for the case of a 10 dB laser gain and a laser length of 200 μm. It is plotted as a function of the modal width $2e$. This quantity is so defined that half the power of the fundamental mode is comprised between $x = -e$ and $+e$.

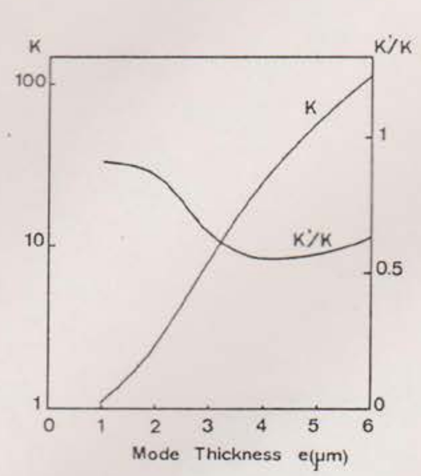

Fig. 3. Ratio of the effective noise enhancement factor K' to Petermann's K-factor for a very thin slab semiconductor laser amplifier of length $L = 200$ μm, gain $G = 10$ dB, and different mode thicknesses $2e = \log(2)/a$. Also shown is Petermann's K-factor from (14), in logarithmic scale. The free-space wavelength is 1 μm. The medium outside the slab is assumed to be free-space, for simplicity.

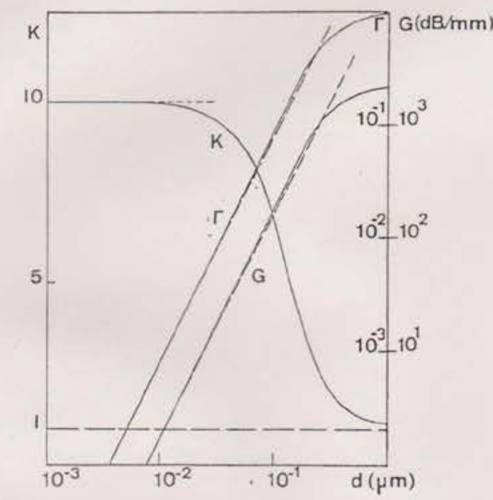

Fig. 4. This figure compares the thin-slab (or reactive-surface) approximation (dashed lines) to exact numerical results (plain lines). In the latter, the free-space wavelength is 1 μm, the slab relative permittivity is $13 - i\,0.3$ in a medium of permittivity 12.9. The slab thickness is $2d$. Γ represents the confinement factor, G the power gain in dB/mm, and K Petermann's K-factor.

As one can see, for typical laser gains and lengths, K' may be on the order of half the K-factor.

One may wonder whether the thin-slab approximation, in which we let the refractive index n go to infinity, is applicable to guiding systems where the relative change of index is on the order of 1 percent. Fig. 4 gives an answer to that question. For a slab of thickness $2d$, and real permittivity change of 0.1 ($\Delta n/n \cong 0.4$ percent), we have plotted the exact numerical results for the confinement factor, the gain, and the K-factor (plain lines) and compared them to the thin-slab approximation. We find that the approximation for K in (14) is very good up to slab thicknesses on the order of 0.02 μm.

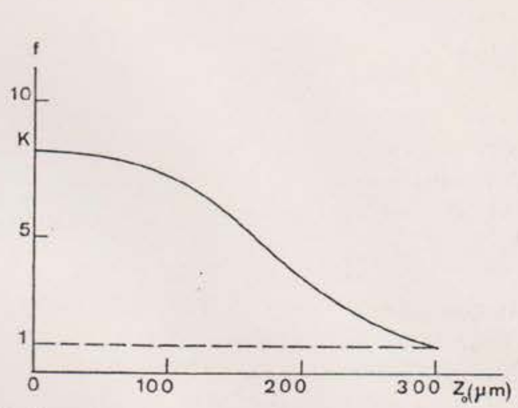

Fig. 5. This figure is analogous to Fig. 2, but applies to the junction plane of a gain-guided laser, and is obtained by purely numerical calculations and application of (5). The integration over z, however, has not been performed. Stripe width = 13 μm, current density $J = 2.7$ kA/cm^2. Net gain = 10 dB.

We have also used a purely numerical technique to evaluate the f-function of (19) for a realistic model of gain-guided laser. This time, the f or K' factors refer to the junction plane as in Petermann's work. The result is shown in Fig. 5. Here again, we notice that f is close to K when spontaneous emission takes place near the input ($z_0 = 0$) and is unity when $z_0 = L$ (here $L = 300$ μm). To obtain these results, we have first used the effective-index method in the junction plane. Then the beam propagation method has been used to proceed from some initial field distribution to the output plane. To obtain the fundamental mode of propagation, we let an arbitrary field propagate back and forth along the structure until a steady situation occurs. This gives us the signal field, since we have assumed, somewhat arbitrarily, that the signal field at the input is in the fundamental mode of propagation. For the spontaneously emitted field, we have used as the input field a narrow Gaussian field (in place of the Dirac function) located at $z = z_0$, and again used the beam propagation method. Finally, the basic equation (5) is used. The integration over z_0 from 0 to L, however, is rather trivial and is not presented here. The realistic calculations presented here for gain-guided lasers are seen to be in qualitative agreement with the results obtained before for thin slabs.

V. Conclusion

In a previous paper [6] one of us had suggested that the spontaneous-noise enhancement factor K obtained by Petermann had to be reconsidered, because the noise powers must be defined in a power-orthogonal system at the laser output. In this paper we have calculated precise values for the effective noise-enhancement factor (K') and found that under the mode-matched input field configuration, it is about half the K value (laser length = 200 μm, gain = 10 dB). Let us emphasize that these results apply to laser *amplifiers* under *mode-matched* conditions. It is not entirely clear how they apply to laser *oscillators*. In laser oscillators the useful field is essentially in mode-matched conditions, because the cleaved facets are plane and perpendicular to the propagation axis. Therefore, this condition of validity of our paper is fulfilled. The next question is whether or not spontaneous emission is strongly filtered. On that respect, it should be remembered that, because of spontaneous emission, the laser gain per pass is significantly smaller than the reciprocal of the facet power reflectivity, and that consequently, an elementary radiating dipole (simulating spontaneous emission) has its radiation emitted only after a few passes, and modal filtering need not be large. If this is the case, this paper has relevance also to laser oscillators near threshold. This will need further investigations.

Acknowledgment

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[9] V. V. Shevchenko, *Continuous Transitions in Open Waveguides*. Boulder, CO: Golem, 1971. Note a few differences in the formulations: our integration over u goes from 0 to ∞, we use a paraxial wave equation, and our "source" is a $\delta(x)$ field function instead of a current.

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