Experimental Investigation and Analytical Modeling of Excess Intensity Noise in Semiconductor Class-A Lasers

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Abstract—Excess intensity noise in a low-noise single-frequency class-A VECSEL is experimentally investigated over the frequency range 10 kHz–18 GHz. An analytical model is derived, based on multimode Langevin equations, to describe the observed laser excess noise over the whole bandwidth. From 50 MHz to 18 GHz, class-A operation leads to a shot noise limited relative intensity noise (RIN), namely $-155$ dB/Hz for 1-mA detected photocurrent, except at harmonics of the cavity free spectral range (FSR). At these frequencies, the excess noise is shown to be due to the amplified spontaneous emission contained in the nonlasing side modes. The measured levels of excess noise correspond to side mode suppression ratios (SMSRs) ranging from 70 to 90 dB, in agreement with the model. At low frequencies, 10 kHz–50 MHz, the observed excess noise spectrum has the expected Lorentzian shape. Its bandwidth increases with the pumping rate to an upper limit given by the cavity photon lifetime. Below this cutoff frequency, we show that the pump RIN is the dominant source of noise, while it is filtered by the laser dynamics above. Finally, our model permits to design a semiconductor class-A laser with an intensity noise limited to the shot noise level over the whole 10 kHz–18 GHz bandwidth.

Index Terms—Laser dynamics, laser noise, pump noise, relaxation oscillations.

I. INTRODUCTION

CONTINUOUS progress in semiconductor lasers spurred numerous practical application areas such as digital optical communications [1], fiber-optic sensors [2], high-resolution spectroscopy [3], and ultrastable atomic clocks [4]. In particular, high-speed semiconductor lasers are attractive coherent light sources for broadband analog communications and more generally for microwave photonics. In this latter field, high-power and large-bandwidth components are mandatory, with bandwidths ranging typically from few tens of megahertz up to 18 GHz [5]. Unfortunately, the dynamic range of a semiconductor laser-based microwave photonic link is drastically deteriorated by the modulation nonlinearities [6] and at the low end by the optical intensity noise in which the laser noise is usually the dominant contribution [7]. Indeed, the commonly used semiconductor class-B lasers such as, e.g., DFB lasers, present damped relaxation oscillations [8], often used to achieve a broadband laser modulation [9]. However, this resonant dynamic behavior is responsible for a highly peaked relative intensity noise (RIN) spectrum around the relaxation oscillations frequency, in the 1–15 GHz frequency range [8]. This excess noise has a direct impact on the dynamic range in high-frequency systems. Thus, one would intuitively think that the use of externally modulated solid-state or fiber class-B lasers with relaxation oscillation frequencies in the 100–500 kHz range would be more beneficial for 100 MHz–18 GHz microwave photonics applications [10]. Unfortunately, this low frequency RIN is transferred around the modulation signal [11]. Furthermore, the level of this transferred noise increases proportionally with the modulation signal level, limiting the system dynamic range to a fixed value that can be exceeded only with efficient low-frequency RIN reduction techniques.

These considerations have motivated several studies which proposed clever techniques to achieve remarkably low RIN levels [12]–[17]. However the proposed methods cover limited frequency ranges. We reported in previous papers on an alternative simple approach to design semiconductor lasers which are intrinsically quantum noise limited over the whole bandwidth of interest [18], [19]. The approach consists of increasing the photon lifetime well above the carrier lifetime in order to eliminate adiabatically the carrier population effects, leading to a relaxation oscillation free class-A laser operation. Two paths have been followed to increase the photon lifetime, either by lengthening the cavity or by using a high-Q cavity. For both configurations, we demonstrated single frequency laser operation with a photodetection current RIN level limited to the shot noise (relative level at $-156$ dB/Hz) from 50 MHz to 18 GHz. This wideband, uniform, and very low RIN level achieved with class-A semiconductor lasers is undoubtedly very attractive for the development of noiseless and consequently large dynamic range microwave photonics systems. Nevertheless, some excess noise can be observed at low frequency and at multiples of the laser cavity free spectral range (FSR). Obviously, its presence is inconvenient for some stringent applications such as distribution of high purity radar signals. Consequently, the optimization of such lasers requires a quantitative prediction of their noise characteristics as function of the cavity parameters.

In this paper, we propose to investigate experimentally and then to describe quantitatively the excess intensity noise in a high-Q semiconductor class-A laser. We choose to focus on such external cavity VECSELs since they seem to be the most promising sources for applications requiring compactness and

Manuscript received June 12, 2007; revised December 4, 2007.
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Digital Object Identifier 10.1109/JLT.2008.917756

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reliability. In the first section, we study experimentally the high-frequency components of the laser intensity noise. In a second section, we propose a simple model based on Langevin rate equations for all the cavity longitudinal modes, which is used to derive analytical expressions for the laser field intensity fluctuations at all frequencies. Then, we discuss the model validity for the description of high-frequency and low-frequency laser RIN. Finally, we explore experimentally and analytically the technical noise contribution to the low-frequency excess noise.

II. EXCESS INTENSITY NOISE IN A HIGH-Q CAVITY CLASS-A SEMICONDUCTOR LASER

A. Class-A SC Laser Description

The experimental setup consists in a 1/2-VCSEL used in a short external cavity (see Fig. 1). For availability reasons, we work at 1000 nm. The 1/2-VCSEL gain chip is grown by metal organic chemical vapor deposition (MOCVD). It comprises a 99.9% reflecting multilayer mirror of 28 pairs of GaAs/AlAs with a resonant periodic gain structure consisting of five InGaAs/GaAs quantum wells. The structure is bonded onto a 150-µm Fabry–Perot étalon, output coupler: 50-mm-radius of curvature and 1% transmission at 1000 nm.

![Fig. 1. (a) Gain chip description. (b) Experimental set up. Filter: uncoated 150-µm Fabry–Perot étalon, output coupler: 50-mm-radius of curvature and 1% transmission at 1000 nm.](image)

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In these conditions, when pumped with 1 W, the laser provides a 140-mW output power with a stable linear polarization. The analysis of the output light with a 7.5-GHz FSR Fabry–Perot interferometer shows that without any intracavity étalon, the laser spectrum is unstable and consists of several longitudinal modes separated by the 3.4-GHz laser cavity FSR. To force single-mode operation, we insert into the cavity a 150-µm-thick uncoated fused silica étalon. The laser output power then decreases to 50 mW for a 1-W pump power corresponding to an operation 2 times above threshold, but stable single-frequency operation is obtained as shown in Fig. 2. One can notice that the side mode suppression ratio (SMSR) is so large that the amplified spontaneous emission (ASE) at the side mode frequencies cannot be seen from the Fabry–Perot spectrum of Fig. 2. Notice also that with the parameters of our laser, we expect its Schawlow–Townes linewidth to be in the mHz range. We thus expect the laser phase noise to play a negligible role in our measurements.

We stress the fact that the cavity length and the output coupler reflectivity are adjusted to get a photon lifetime well above the carrier lifetime. Indeed, a cavity length of 45 mm and 1.5% net losses per round trip lead to a photon lifetime equal to 20 ns, larger than the carrier lifetime (a few ns). Thus, the laser is expected to exhibit class-A relaxation oscillation free operation and a low-intensity noise limited to the shot noise level over a broad frequency bandwidth.

B. Broad Bandwidth Measurement of the Relative Intensity Noise

To confirm the class-A white noise laser behavior, we observe the laser intensity noise from 100 MHz to 18 GHz. The RIN is measured with a bench developed and optimized at Thales R&T [18]. In this bench, we use a wideband (dc-22 GHz), high-power, and low-noise InGaAs PIN photodiode optimized for high power and wide bandwidth operation at 1550 nm. At 1000 nm, it exhibits a rather low efficiency (0.4 A/W) and is significantly nonlinear above 2 mA of detected photocurrent. Consequently, we deliberately attenuate the 50-mW laser output power to limit the detected photocurrent to 1 mA, namely in the linear response region of the photodetector. A commercial low-noise stabilized nonplanar Nd:YAG ring laser [16], [22] is used for calibration. Such a calibration spectrum obtained with that Innolight laser is reproduced in the inset of Fig. 3(a). A typical result of an RIN measurement for our VECSEL is displayed in Fig. 3(a) together with the expected shot noise level lying at −155 dB/Hz for our 1-mA photocurrent (dashed line). It is worth mentioning that this microwave spectrum throws light on details that cannot be seen on the optical spectrum of Fig. 2. We notice, first, that no intensity noise peak related to relaxation oscillations is observed over this bandwidth. Second, it clearly proves that the laser RIN is white and coincides with the shot noise relative level, i.e., −155 dB/Hz for 1-mA detected photocurrent, except for the tiny peaks at 3.4, 6.8, and 10.2 GHz.
This excess noise corresponds to beat notes between the laser line and the residual ASE at the side longitudinal mode frequencies. Consequently, the beat notes between the laser signal and the ASE at the nonlasing mode frequencies consist in peaks lying at harmonics of the laser cavity FSR \( \Delta \nu_{\text{FSR}} \). Notice that this phenomenon is also present in the case of the calibration commercial low-noise laser [see the insert of Fig. 3(a)]. The amplitudes of the peaks seen in Fig. 3(a) decrease for increasing side mode orders. The ASE level, i.e., the corresponding SMSR of the \( q \)th nonlasing mode, can be derived from the beatnote level at the \( q \)th harmonics of the laser cavity FSR, as done in the following subsection.

C. SMSR Measurement From Microwave Intensity Spectrum

The SMSR of the \( q \)th nonlasing mode is defined as the ratio between the ASE optical power at the \( q \)th nonlasing mode frequency and the optical power of the lasing mode. In order to evaluate the SMSR of the mode lying \( q \Delta \nu_{\text{FSR}} \) away from the laser line frequency, we measure the power \( P_q \) of the electrical beat note signal at frequency \( q \Delta \nu_{\text{FSR}} \) and then compare it to the noise floor corresponding to the shot noise power \( P_{\text{shot}} \) and the detected photocurrent \( i_{\text{ph}} \). The SMSR for the \( q \)th side mode can then be extracted from the experimental data through

\[
\text{SMSR}_q = \frac{P_q e B}{2P_{\text{shot}}i_{\text{ph}}} \tag{1}
\]

where \( e \) is the elementary electrical charge, and \( B \) is the measurement bandwidth. The factor 2 in the denominator of (1) comes from the fact that two side-modes located symmetrically with respect to the lasing mode contribute to the beat note. From (1) and the spectrum of Fig. 3(a), the SMSR measured at 3.4 GHz is equal to 73 dB. This proves the high extinction of the ASE at 3.4 GHz compared to the field intensity at the laser line. Moreover, the 50-kHz full width at half maximum (FWHM) of the beating noise at 3.4 GHz displayed in Fig. 3(b) proves again the drastic filtering of the ASE inside the laser cavity. At 6.8 and 10.2 GHz, the measured SMSRs are, respectively, as large as 77 and 87 dB. At higher modes frequencies, the excess intensity noise lies below the shot noise level.

This drastic filtering of the side modes evidences the remarkable filtering efficiency of the laser cavity. Indeed, one could naively state that the SMSR should be given by the product of the transmissions of the high-Q Fabry–Perot laser cavity (99.9% Bragg mirror and 99% output coupler, FSR = 3.4 GHz) and the low finesse intracavity uncoated 150-\( \mu \)m Fabry–Perot étalon (4% reflectivity per facet, FSR = 606 GHz). This would lead to SMSR values much lower than the measured ones. Therefore, to understand our results and further optimize the laser noise, it is necessary to take into account the effects of the gain medium and the filtering elements on the lasing and the nonlasing components of the laser field. This will be done in the model described in the next section.

III. SIMPLE MODEL FOR THE EXCESS INTENSITY NOISE

Many authors have proposed models that permit to predict the intensity noise of class-A lasers [23]. However, as we have just seen from the experimental results of Fig. 3, even if our laser is monomode, we have to deal with the noise present at all the cavity modes. We thus have to use a model describing the field and its fluctuations in several cavity modes. In the following subsection, we thus implement a semiclassical description of the laser dynamics involving the field amplitude of the different modes and the population inversion. Spontaneous emission is taken into account using Langevin forces.

A. Langevin Rate Equations

The laser dynamics is controlled by the coupled instantaneous variations of the laser electric field inside the cavity and the carrier density in the gain medium. Here, we can expand the laser electric field \( \vec{E} \) in terms of the longitudinal modes of the cavity:

\[
\vec{E}(x, y, z, t) = \sum_q \vec{U}_q(x, y) A_q(t) e^{i(2\pi n k_q z)} + c.c. \quad (2)
\]

where \( A_q \) is the slowly varying amplitude of the \( q \)th longitudinal cavity mode at the optical frequency \( \nu_q \), \( k_q \) is the optical wave number, and \( c.c. \) denotes complex conjugation. \( \vec{U}_q(x, y) \) contains the transverse spatial dependence of the mode amplitude.

The Langevin differential equations for the amplitudes \( A_q \) of the cavity modes and the number of carriers \( N \) in the 1/2-VCSEL are

\[
\frac{dA_q}{dt} = -\gamma_q A_q + \frac{K}{2} N A_q + \xi_q(t) \quad (3)
\]

\[
\frac{dN}{dt} = \Gamma(N_{\text{premp}} - N) - \frac{2\sigma}{hv} \sum_q |A_q|^2 + F_N(t) \quad (4)
\]

where \( \gamma_q \) is the cavity decay rate for the intensity of the mode of order \( q \). \( K \) is the stimulated emission rate defined as \( K = \sigma c / \nu_{\text{cav}} \), where \( \sigma \) is the laser cross section and \( \nu_{\text{cav}} \) is the laser field volume in the cavity. We notice that the gain bandwidth of our laser (40 nm) is large compared with the bandwidth we considered above (±20 GHz, i.e., 0.13 nm), allowing us to take \( K \)
independent of $q$. We also suppose that the gain is purely homogeneously broadened. To keep our model as simple as possible, we treat the active quantum wells as a four-level system where $\Gamma$ is the inversion decay rate and $N_{\text{pump}}$ holds for the pumping. $h$ is Planck’s constant. $\xi_q(t)$ is the Langevin force representing the spontaneous emission falling in the $q$th cavity mode, and $F_N$ is the Langevin force driving the carrier number fluctuations. These forces have vanishing average values, and their powers will be discussed later.

Contrary to the gain, the losses are different for the successive modes because the laser contains an étalon with a frequency-dependent transmission $T$ given by

$$T(\nu) = \left(1 + \frac{4R}{(1-R)^2} \sin \left(\frac{2\pi \nu L_{\text{opt}}}{c}\right)\right)^{-1}$$  \hspace{1cm} (5)

where $\nu$ is the optical frequency, and $c_{\text{opt}}$ and $R$ are, respectively, the étalon optical thickness and facet reflectivity. The transmission of this filter is maximum and equal to 1. We suppose that the central mode frequency that we call $\nu_0$ coincides with this transmission maximum. Then, the étalon transmission takes lower values at side mode frequencies. Hence, the cavity decay rate $\gamma_q$ for the mode of order $q$ is larger than the cavity loss decay rate $\gamma_0$ for the mode which has the lowest losses. It reads

$$\gamma_q = \gamma_0 + \delta \gamma_q = \gamma_0 + \frac{c}{L_{\text{cav}}} \left[\frac{2}{1 - T(\nu_q)}\right]$$  \hspace{1cm} (6)

where $L_{\text{cav}}$ is the optical length of one cavity round-trip (i.e., twice the cavity optical length), and where we have used the fact that the losses are very weak. Consequently, the side mode at frequency $\nu_q$ is killed by the étalon additional losses and its steady-state mean amplitude is zero:

$$\langle A_q \rangle = 0$$  \hspace{1cm} (7)

where $\langle \rangle$ means ensemble averaging. Above threshold, the average laser solutions are

$$\langle N \rangle = \frac{\gamma_0}{\Gamma} = N_0, \quad \langle A_0 \rangle^2 = \frac{\Gamma}{2\gamma} h\nu/(\eta - 1)$$  \hspace{1cm} (8) \hspace{1cm} (9)

where $\eta = N_{\text{pump}}/N_0$ and where we have defined $N_0$, the population inversion at threshold. Within the single mode operation discussed above ($\langle A_q \rangle = 0$ for $q \neq 0$), a rigorous description of the laser field amplitude should include separate equations on the lasing and the nonlasing modes. We introduce the small fluctuations $a_0, a_q$, and $n$ around these average values in the following manner:

$$A_0(t) = \langle A_0 \rangle + a_0(t)e^{i\phi_0(t)}$$  \hspace{1cm} (10)

$$A_q(t) = a_q(t)e^{i\phi_q(t)}$$  \hspace{1cm} (11)

$$N(t) = N_0 + n(t)$$  \hspace{1cm} (12)

where we have separated the phases $\phi_0$ and $\phi_q$ to allow $a_0$ and $a_q$ to be real. By introducing (10)–(12) into (3) and (4) and keeping only first-order terms, one obtains

$$\frac{da_0}{dt} = \frac{K}{2} \langle A_0 \rangle n + \xi_0(t)$$  \hspace{1cm} (13)

$$\frac{da_q}{dt} = -\frac{\delta \gamma_q}{2} a_q + \xi_q(t)$$  \hspace{1cm} (14)

$$\frac{dn}{dt} = -\eta \Gamma n - \gamma_0 \frac{2\Gamma}{K} (\eta - 1) \langle A_0 \rangle + F_N(t)$$  \hspace{1cm} (15)

with $\eta = N_{\text{pump}}/N_0$ is the pumping rate and the Langevin noise terms $\xi_q(t)$ and $\xi_q(t)$ are the optical thickness and facet reflectivity. The étalon optical thickness and facet reflectivity. The transmission of this filter is maximum and equal to 1. We suppose that the central mode frequency that we call $\nu_0$ coincides with this transmission maximum. Then, the étalon transmission takes lower values at side mode frequencies. Hence, the cavity decay rate $\gamma_q$ for the mode of order $q$ is larger than the cavity loss decay rate $\gamma_0$ for the mode which has the lowest losses. It reads

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with $\eta = N_{\text{pump}}/N_0$ is the pumping rate and the Langevin noise terms $\xi_q(t)$ and $\xi_q(t)$ are independent of the other variables and can be treated independently. Hence, in the following, we derive first the intensity noise of the lasing mode using (13) and (15). Second, we will derive from (14) the intensity noises at the side mode frequencies.

**B. Relative Intensity Noise at the Central Mode Frequency**

In order to use more usual notations, the laser field component at the central frequency will be described with the number of photons $\Phi$ in the mode. It is related to the lasing mode amplitude by the following equation:

$$\Phi = \frac{V_{\text{cav}}}{c} \frac{2\langle A_0 \rangle^2}{h\nu_0}.$$

With this notation, the Langevin rate equations for the lasing mode photon number and the population inversion become

$$\frac{d\Phi}{dt} = -\gamma_0 \Phi + K N \Phi + F_\Phi(t),$$  \hspace{1cm} (17)

$$\frac{dN}{dt} = \Gamma(N_{\text{pump}} - N) - K N \Phi + F_N(t).$$  \hspace{1cm} (18)

The steady-state mean solutions are $\Phi_0 \equiv \langle \Phi \rangle = (\Gamma/K)(\eta - 1)$ and $N_0 = \gamma_0/K$. $F_\Phi(t)$ is the Langevin noise source holding for the effect of the spontaneous emission on the photon number. Its power spectral density corresponds to the usual “one photon per mode” approximation [24]:

$$S_{F_\Phi}(\omega) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} \langle F_\Phi(t)F_\Phi(t + \tau) \rangle e^{i\omega\tau} d\tau = \frac{\gamma_0^2 \Phi_0}{\pi}.$$  \hspace{1cm} (19)

The Langevin force for the population inversion takes into account both the spontaneous emission and the probabilistic nature of the upper level decay

$$S_{F_N}(\omega) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} \langle F_N(t)F_N(t + \tau) \rangle e^{i\omega\tau} d\tau = \frac{1}{\pi} (\gamma_0 \Phi_0 + \Gamma N_0).$$  \hspace{1cm} (20)
and

\[
S_{F,NF}(\omega) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} \langle F_N(t)F_\Phi(t+\tau) \rangle e^{i\omega\tau} d\tau = -\frac{\gamma_0\Phi_0}{\pi}.
\]  

(21)

Similarly to the previous subsection, we derive the small signal Langevin equations for the time-fluctuating components of the photon and the carrier number where \(F(t) = \Phi(t) + \varphi(t)\) and \(N(t) = N_0 + n(t)\). Then, we take the Fourier transform of the resulting equations and obtain a set of linear coupled equations for the Fourier transforms \(F\) and \(N\) of \(\varphi(t)\) and \(n(t)\) respectively,

\[
\dot{\eta} = (\eta - 1)\gamma N - \Gamma F
\]

(22)

\[
\gamma_0\Phi + (\eta\gamma - \dot{\eta})N = \dot{F}_N.
\]

(23)

By eliminating \(\dot{N}\) in (22) and (23) and using (19)-(21), we obtain the power spectral density of \(\Phi\):

\[
S_{\Phi}(\omega) = \frac{2\gamma}{[\omega^2 - (\eta - 1)\gamma_0\Gamma]^2 + \eta^2\omega^2\Gamma^2}.
\]

(24)

Finally, the relative intensity noise driven by the lasing mode fluctuations can be written as follows:

\[
RIN_0(\omega) = \frac{2S_{\Phi}(\omega)}{\Phi_0^2} = \frac{2\gamma m\gamma_0\mu_0}{\pi P_{out}} \left[\frac{(\omega^2 + \Gamma^2\eta)}{\omega^2 - (\eta - 1)\gamma_0\Gamma^2 + \eta^2\omega^2\Gamma^2}\right] \cdot
\]

(25)

where the factor of “2” is due to the fact that we define the RIN for positive frequencies only contrary to power spectral densities given by (19)-(21), \(\gamma_m\) is the part of the cavity decay rate due to the output coupler transmission, and \(P_{out}\) is the laser optical output power related to the intracavity photon number by the following equation:

\[
P_{out} = \Phi_0\gamma_m\mu_0.
\]

(26)

When considering a class-A laser, the photon decay rate \(\gamma_0\) is much smaller than the carrier inversion decay rate \(\Gamma\). Hence, the lasing mode relative intensity noise \(RIN_0\) can be simplified as follows:

\[
RIN_0(\omega) \approx \frac{2\gamma m\gamma_0\mu_0}{\pi P_{out}} \frac{1}{\omega^2 + \gamma_0^2 \left(\frac{\omega}{\eta}\right)^2} \left[\frac{\omega^2 + \eta\Gamma^2}{\omega^2 + \eta^2\Gamma^2}\right]^2.
\]

(27)

We notice that the RIN expression comprises two frequency dependent terms. The first term corresponds to a first-order filter represented by a Lorentzian function. The second part, delimited between two square brackets, is equal to \(1/\eta\) for angular frequencies much lower than \(\Gamma\) and converges to 1 for large angular frequencies. Thus, since \(\Gamma > > \gamma_0\), we can consider that the shape of the RIN spectrum is properly described by the first term which is a first-order low-pass filter. The FWHM of this filter \(\Delta\omega_{RIN}\), depends on the photon decay rate \(\gamma_0\) and the pumping rate \(\eta\). It reads as follows:

\[
\Delta\omega_{RIN} = 2\gamma_0\frac{\eta - 1}{\eta}.
\]

(28)

The derived RIN expression describes the laser intensity noise over a wide microwave frequency bandwidth except at frequencies around harmonics of the cavity FSR \(q\Delta\nu_{cav}\). In fact, the ASE enhanced around the nonlasing modes optical frequencies is responsible for excess intensity noise at harmonics of the cavity FSR. This excess noise will be derived in the following subsection.

C. Excess Intensity Noise Derivation at Side-Mode Frequencies

The amplitude fluctuation \(a_q(t)\) of the nonlasing mode \(q\) is driven by the Langevin noise term \(\xi_q(t)\) and damped by the intracavity étalon extra losses \(2(1-T_q)\) as explained in (14). The Langevin function \(\xi_q(t)\) has the same characteristics as \(\xi_0(t)\). Its power spectral density, corresponding also to the well-known “one photon per mode” approach, reads as follows:

\[
S_{\xi_q}(\omega) = \frac{\gamma_0}{4\pi} \frac{\langle A_0^2 \rangle}{\Phi_0} = S_{\xi_0}(\omega).
\]

(29)

By using the Fourier transform of (14), we obtain a simple expression of the Fourier transform \(a_q(\omega)\) of \(a_q(t)\). Then, using (29), we deduce the power spectral density of the fluctuating amplitude \(a_q(t)\):

\[
S_{a_q}(\omega) = \frac{\gamma_0}{4\pi} \frac{\langle A_0^2 \rangle}{\Phi_0} \left[\frac{1}{(\gamma_0\eta_q/2)^2 + \omega^2}\right].
\]

(30)

The relative intensity noise of the laser at frequencies close to any beatnote frequency \(2\pi q\Delta\nu_{cav}\) will be the sum of the beatnotes of the lasing mode with the \(q\)th and \(-q\)th modes. For frequencies close to \(2\pi q\Delta\nu_{cav}\), the RIN will thus be given by

\[
RIN_q(\omega) = \frac{S_{a_q}(\delta\omega) + S_{a_{-q}}(-\delta\omega)}{\langle A_0^2 \rangle} = \frac{\gamma_m\gamma_0\mu_0}{2\pi P_{out}} \frac{1}{(\gamma_0\eta_q/2)^2 + \omega^2}.
\]

(31)

provided \(\delta\omega \ll 2\pi q\Delta\nu_{cav}\) and \(\gamma_0\eta_q \approx \\gamma_q\). We notice that this excess intensity noise exhibits a Lorentzian shape with a FWHM equal to \(\gamma_0\eta_q\). To know whether the laser noise is below or above the shot noise level, we derive the ratio of the maximum value of the RIN at the \(q\)th harmonics of the cavity FSR with the corresponding shot noise level \(RIN_{shot} = 2\rho P_{out}/(\rho P_{out})^2\):

\[
RIN_{q}(q\Delta\nu_{cav}) = \frac{\gamma_m\gamma_0\mu_0}{\delta\gamma_q} \frac{1}{\pi e}
\]

(32)

where \(\rho\) is the detector responsivity. Since \(\rho\mu_0/e \approx 1\), one can see that the RIN around the beatnote frequency can be much larger than the shot noise level.

To compare with the experimental measurements (see Fig. 3), we derive the value of the laser SMSR for the \(q\)th side mode from the variance \(\sigma_{a_q}^2\) of the nonlasing mode amplitude \(a_q\):

\[
\sigma_{a_q}^2 = \langle a_q(t)^2 \rangle = \frac{\gamma_0}{2\gamma_q} \frac{\langle A_0^2 \rangle}{\Phi_0^2}.
\]

(33)

Thus, the SMSR at \(q\Delta\nu_{cav}\) reads as follows:

\[
SMSR_q = \frac{\sigma_{a_q}^2}{\langle A_0^2 \rangle} = \frac{\gamma_0\gamma_m\mu_0}{2\gamma_q P_{out}}.
\]

(34)
D. Comparison With Experimental SMSR Results and Discussion

Let us now discuss the model validity by confronting first the experimental RIN spectra presented in Section II to the model predictions.

At high frequencies, from 100 MHz to 18 GHz, we observe excess intensity noise at the cavity FSR harmonics. It takes the shape of narrow and tiny beat notes at 3.4, 6.8, and 10.2 GHz as reported in Fig. 3(a). A complete characterization of this excess noise should include its FWHM and amplitude investigation. According to (31), the excess intensity noise peak at the $q$th cavity FSR harmonics should have a Lorentzian shape with a FWHM equal to $\frac{\delta \gamma_0}{2\pi}$. Let us focus for example on the beating noise at 3.4 GHz ($q = 1$). With our experimental parameters (150 $\mu$m étalon thickness with a facet reflectivity equal to 4%), we obtain $9 \times 10^{-5}$ for the extra losses per round-trip, leading to $\delta \gamma_1 = 3 \times 10^2$ s$^{-1}$, and thus an FWHM $\delta \gamma_1/2\pi = 47$ kHz. Such a 47-kHz-broad Lorentzian is compatible with the experimental spectrum, which is well fitted by the 50-kHz-broad Lorentzian shown in Fig. 3(b).

Let us now turn to the amplitude of this excess noise. Here, the noise amplitude at the microwave frequency $q\Delta \nu_{\text{cav}}$ is described by the extinction ratio of the corresponding side mode at the optical frequency $\nu_{\text{cav}}$, as given analytically in (34). It clearly states that the side mode rejection efficiency increases with larger étalon additional losses. For the considered class-A laser, the total cavity losses are estimated to be equal to 1.5%, leading to a cavity decay rate $\gamma_0$ equal to $5.1 \times 10^7$ s$^{-1}$. The 1% mirror transmission leads to $\gamma_m$ equal to $3.4 \times 10^7$ s$^{-1}$. For a 50-mW optical output power and using (34), we predict that the peak at 3.4 GHz is lying at 79 dB below the power level of the oscillating mode. Experimentally, we measure a SMSR value (for the peak at 3.4 GHz) equal to 73 dB. We remark that the experimental SMSR is independent from the detected photocurrent as long as the excess intensity noise is above the shot noise limit. Hence, it is justified to compare the theoretical SMSR calculated for the overall power (50 mW) to the experimental one measured after attenuating the overall power to 1 mA. For the second and third side modes, the étalon additional losses are respectively of about $3.5 \times 10^{-4}$ and $8 \times 10^{-4}$. We thus expect the SMSR at 6.8 and 10.2 GHz to be, respectively, equal to 85 and 89 dB while the experimental values are respectively equal to 77 and 87 dB. These experimental and analytical results are summarized in Fig. 4. The agreement between experiments and our simple model is satisfactory. The few dB difference between the theoretical and experimental SMSR values could be due to the fact that the étalon resonance is not exactly centered on the lasing mode, leading to some uncertainties on the values of the $\delta \gamma_j$’s. It could also be due to gain related nonlinear effects that are not included in the proposed model, or to more subtle effects related to the pumping statistics that could be taken into account only in the framework of a more complicated fully quantized model [26], [27]. Anyway, the predictions of our simple model are reliable enough to allow us to estimate the excess intensity noise at the laser cavity FSR harmonics. Furthermore, we can use the model to predict the optimized cold cavity response for which the excess intensity noise would be below the shot noise level. In fact, the observed noise at the nonlasing modes frequencies could be further reduced by using a higher quality factor cavity (smaller $\gamma_0$) or an intracavity filter with a larger finesse (larger values for $\delta \gamma_j$). For example, the noise at 3.4 GHz, lying 24 dB above the shot noise relative floor, could be rejected below $-155$ dB/Hz with an étalon filtering more than $2,7 \times 10^{-3}$ of the field component at 3.4 GHz. Moreover, with étalon losses equal to $1.9 \times 10^{-2}$, we predict that the peak could be rejected below the $-172$-dB/Hz level corresponding to the shot noise for a 50-mW detected power. Another alternative consists in reducing the cavity length in order to shift the intensity excess noise frequencies in the region where the étalon filtering is more efficient. This approach has been explored in [19] where we demonstrated a 8-mm cavity length class-A laser with a white RIN spectrum limited to the shot noise relative level, $-156$ dB/Hz, from 100 MHz to 18 GHz.

E. Application to the Low-Frequency RIN and Discussion

Since our simple model just allowed us to successfully reproduce the measured SMSR values, we now wonder whether it may describe conveniently the laser RIN at low frequencies. For this purpose, we experimentally investigate the laser intensity noise from 10 kHz to 500 MHz. The RIN is measured with a bench similar to the one described for high-frequency RIN measurements, except that we use this time a silicon photodiode with a lower cutoff frequency. The detected voltage is amplified with an ultralow-noise 56-dB gain electrical amplifier (Miteq AU 1291) and then analyzed with an electrical spectrum analyzer. For low-frequency measurements, the calibration of the bench can no longer be performed using the In-laser laser since its relaxation oscillations increase its noise in this spectral domain. Hence, a low-noise white source (an incandescence lamp) is used for calibration. Here again, we deliberately attenuate the 50-mW laser output power to limit the detected photocurrent to 1 mA, in order to preserve the linear response of the photodetector. A typical result of this RIN measurement is displayed in Fig. 5(a). The obtained microwave spectrum exhibits a low-pass filter shape compatible with a 40-GHz cutoff frequency and shows a RIN below $-180$ dB/Hz in this domain. The obtained RIN is compared to the theoretical predictions in Fig. 5(b) and (c). As expected, a good agreement is observed, including the flatness of the RIN for the nonlasing modes frequencies.
with our class-A laser model [see (27)]. We thus fit the laser low-frequency RIN with a Lorentzian function for several laser pumping rates and then extract the corresponding FWHM. Two typical experimental data and fit results for frequencies spanning from 10 kHz to 400 MHz are reproduced in Fig. 5(b). By performing the same analysis on several records, we obtain the evolution of the RIN bandwidth versus pumping rate \( \eta \) displayed in Fig. 6. These data fit well the prediction of (28) (full line in Fig. 6) with \( \gamma_0 = 5.1 \times 10^7 \) s\(^{-1}\), in agreement with total cavity losses equal to 1.5%.

However, even if the model is quite satisfactory to predict the bandwidth of the low-frequency noise component, it does not provide the actual RIN level. Indeed, from (23) with a 50-mW output power and inversion population decay rate \( \Gamma \) equal to 3.3 \( 10^6 \) s\(^{-1}\), we expect a RIN level as low as -173 dB/Hz at 10 kHz while the measurement of Fig. 5(a) corresponds to a RIN as large as -126 dB/Hz at 10 kHz. This huge discrepancy evidences the existence of another noise source than the intrinsic spontaneous emission. Since the bandwidth of the intensity noise created by this noise source matches the model prediction, we suspect that the corresponding noise source must be related to the gain [27]. In the following section, we thus investigate experimentally the pump intensity noise and we model and observe its transfer to the laser RIN.

### IV. Experimental Study of Pump-to-Laser RIN Transfer and Comparison to Model Predictions

Here, we fix the laser pumping rate to 1.8 corresponding to 66 mW of output optical power. In this case, the pump diode pumping rate is equal to 3.2 which allows us to observe comfortably its intensity fluctuations. The laser low-frequency RIN half width at half maximum is equal to 3 MHz (see Fig. 6). We first measure the laser and the pump RIN for frequencies below the laser microwave bandwidth. The resulting RIN spectra are presented in Fig. 7(a). Two points should have our attention in these microwave spectra. First, the measured pump RIN is large enough, -135 dB/Hz, to play a major role in the laser excess intensity noise. In fact, the pump source is a multimode diode whose low frequency RIN is probably enhanced by mode partition noise. Second, the laser RIN lies 7 dB above the pump RIN.

To clarify the origin of this 7-dB offset, we investigate the laser response to pump intensity fluctuations. Other works studied this point [27] in which a Langevin noise approach is used to include the pump intensity fluctuations. Here, we use a simpler deterministic approach where we represent the pump intensity fluctuations by a fluctuating pumping rate \( \delta I(t) \) around a mean value \( I_0 \). The differential equation governing the lasing mode intensity \( I(t) = 2A_0(t)A_0^*(t) \) then reads as follows:

\[
\frac{dI}{dt} = \gamma_0 I(t) \left( -1 + \frac{\eta(t)}{1 + I(t)/I_{\text{sat}}} \right)
\]

where \( I_{\text{sat}} = (\Gamma/\sigma)\mu_0 \) and the steady-state mean value is \( I_0 = I_{\text{sat}}(\eta_0 - 1) \). With \( \eta(t) = \eta_0 + \delta \eta(t) \), the small-signal equation for the intensity fluctuations \( \delta I(t) = I(t) - I_0 \) is

\[
\frac{d\delta I}{dt} = -\gamma_0 - \frac{1}{\eta_0} \delta I + \gamma_0 I_0 \delta \eta,
\]

After a Fourier transform of (34), we derive a simple equation coupling \( \delta I \) and \( \delta \eta \), the Fourier transforms of \( \delta I(t) \) and \( \delta \eta(t) \), respectively, leading to

\[
\left| \frac{\delta \eta}{\delta I} \right|^2 = \frac{(\gamma_0 I_0/\eta_0)^2}{\omega^2 + \gamma_0^2 \left( \frac{\eta_0 - 1}{\eta_0} \right)^2}.
\]

![Fig. 5](image)

(a): RIN spectrum measured from 10 kHz to 500 MHz with a RBW = 10 kHz. The laser pumping rate \( \eta \) is fixed to 1.5 corresponding to an optical output power equal to 50 mW. The RIN level coincides with the shot noise relative level at -155 dB/Hz for 1-mA detected photocurrent. The dashed black line is the fit result with a Lorentzian function. The FWHM is equal to 9.5 MHz. The full gray line is for a pumping rate equal to 1.7. The continuous black line is the fit result with a Lorentzian function. The FWHM is equal to 5.25 MHz. The full gray line is for a pumping rate equal to 2.4.

![Fig. 6](image)

Squares: low-frequency RIN FWHM measured for different pumping rates. Full line: theoretical evolution using (28) for 1.5% total cavity losses.
is the shot noise relative level at 1 mA. The data in gray is the laser RIN for a pumping rate equal to 1.8. The corresponding pump RIN spectrum is represented by the black line. The dashed line is the shot noise relative level at $-155$ dB/Hz for 1-mA detected photocurrent. (b) Gray line: experimental pump-to-laser RIN transfer factor for a pumping rate equal to 1.8 [ratio of the spectra of Fig. 7(a)]. Black line: fit using (38) for total cavity losses equal to 1.5% and a pumping rate equal to 1.8.

Fig. 8. (a) Comparison between the laser and the pump RIN spectra from 3 to 50 MHz measured with a RBW $= 10$ kHz and for a detected photocurrent equal to 1 mA. The data in gray is the laser RIN for a pumping rate equal to 1.8. The corresponding pump RIN spectrum is represented by the black line. The dashed line is the shot noise relative level at $-155$ dB/Hz for 1-mA detected photocurrent. (b) Gray line: experimental pump-to-laser RIN transfer factor for a pumping rate equal to 1.8 [ratio of the spectra of Fig. 8(a)]. Black line: fit using (38) for total cavity losses equal to 1.5% and a pumping rate equal to 1.8.

The function $H_{\text{RIN}}$, which describes the transfer of the pump RIN onto the laser RIN is obtained by normalizing $\langle \delta I \rangle^2$ and $\langle \delta \eta \rangle^2$ to $I_0^2$ and $\eta_0^2$

$$H_{\text{RIN}}(\omega) = \frac{RIN_L \times N_L}{RIN_P \times N_P} = \frac{\gamma_0^2}{\omega^2 + \gamma_0^2} \left(\frac{\eta-1}{\eta_0}\right)^2$$

where $RIN_L$ is the laser RIN driven by the pump intensity fluctuations and $RIN_P$ is the pump RIN. In the same way as in (27), the formula in (38) proves the low-pass filter behavior of the class-A laser regarding intensity fluctuations. However, for frequencies well below the laser cutoff frequency $\gamma_0(1-1/\eta_0)/2\pi$, the pump RIN converted into the laser RIN is enhanced by a factor larger than unity and varying as $\gamma_0^2/(\eta_0-1)^2$. This factor is close to unity when the laser is operating at large pumping rates. We obtain experimentally the pump to laser RIN transfer factor, $H_{\text{RIN}}$, from 10 kHz to 3 MHz, by dividing the two experimental spectra of Fig. 7(a). The resulting transfer factor is reproduced by Fig. 7(b). We obtain a good agreement between the experimental result and the theoretical curve obtained from (38). The full line displayed in Fig. 7(b) is obtained for a pumping rate of 1.8 and for total cavity losses equal to 1.5%. In particular, we obtain theoretically and experimentally a 7-dB transfer factor at low frequencies (typically below 1 MHz). Thus, we confirm that the observed excess intensity noise is totally driven by the pump intensity fluctuations with a large enhancement for low pumping rates.

This enhancement is predicted to vanish for frequencies above the laser cutoff frequency. This is confirmed by our measurements of the laser and pump RIN from 3 to 50 MHz [see Fig. 8(a)]. The experimental pump-to-laser RIN transfer factor fits well the factor predicted using (38). As shown in Fig. 8(b), the pump RIN can be reduced by up to 15 dB before disappearing below the shot noise limit, $-155$–5 dB/Hz, for 1-mA detected photocurrent.

V. CONCLUSION

In summary, we have presented an experimental study of excess intensity noise in a low-noise high-Q class-A semiconductor laser. The laser under study consisted in a 1/2-VCSEL in an external cavity including an etalon filter where the cavity length and the output coupler transmission were optimized for oscillations-relaxation free class-A operation. This dynamics led to a relative intensity noise level limited to the shot noise relative level, $-155$ dB/Hz for 1-mA detected photocurrent, over a wide bandwidth from 50 MHz to 18 GHz. Excess intensity noise was observed below 50 MHz and at harmonics of the cavity FSR. It has a low-pass filter shape at the low-frequency part while it takes the form of tiny peaks at the high-frequency part. Its highest level reaches $-128$ dB/Hz at 10 kHz for a pumping rate equal to 1.8. This level is much lower than the excess intensity noise ($-100$ dB/Hz) at the relaxation oscillations frequency in the Innolight laser [16], well known for its low RIN performance. Furthermore, the excess noise level in the high-frequency range is comparable to the excess noise at the FSR harmonics of the Innolight laser (see Fig. 2). These results have been obtained without any active stabilization.
For an analytical description of excess noise in a single-frequency high-Q class-A laser, we proposed a simple model based on Langevin rate equations for the laser field components at the multiple cavity modes frequencies and the gain medium population inversion. The excess noise at the qth cavity FSR harmonics has a Lorentzian shape with a FWHM and a peak level totally determined by the étalon additional losses. With our experimental parameters (150 μm uncoated solid étalon), the excess noise at 3.4 GHz is predicted to be 47-kHz broad. This is in good agreement with the 50-kHz broad measured beat note at 3.4 GHz. Moreover, the model predicts SMSRs as high as 79, 85, and 89 dB, respectively, at 3.4, 6.8, and 10.2 GHz, in good agreement with the experimental values respectively equal to 73, 77, and 87 dB.

On the low-frequency side, the class-A operation is expected to lead to a Lorentzian RIN spectrum with a FWHM increasing for an increased laser pumping rate and reaching a stationary value proportional to the photon decay rate. The expected Lorentzian RIN shape agrees with the experimental RIN spectra. Moreover the RIN bandwidth values measured for several pumping rates are compatible with the model predictions for total cavity losses equal to 1.5%. For the 45-mm long-cavity and assuming 1.5% total losses, the RIN microwave bandwidth should converge to 7.5 MHz.

However, the ~126-dB/Hz RIN level at 10 kHz measured for a pumping rate equal to 1.5 is incompatible with the ~173-dB/Hz expected RIN level when assuming the spontaneous emission as the noise source driving the laser field fluctuations. Hence, we investigate experimentally the effects of pump noise source on the laser noise properties. Below the laser microwave cutoff frequency and for a pumping rate equal to 1.8, the pump intensity fluctuations are transferred to the laser RIN with an enhancement factor equal to 7 dB. For frequencies larger than the laser bandwidth, the laser filters out the pump intensity fluctuations. For still larger frequencies, these fluctuations disappear below the expected shot noise relative level. The derived simple model for pump-to-laser noise transfer description is in excellent agreement with the experimental observations. It predicts that the pump-to-laser transfer factor vanishes to unity for high pumping rates.

The predictions of our simple model are reliable enough to allow us to use it for the optimization, in terms of intensity noise, of this kind of laser. As an example, the observed noise at 3.4 GHz could be further rejected below the ~172 dB/Hz level corresponding to the shot noise for a 50-mW detected power with étalon additional losses larger than 1.9 × 10^{-2} at the side mode frequency. For a 150-μm-thick étalon, this corresponds to an étalon with facet reflectivities equal to 72%. At the low-frequency side, the RIN bandwidth gets narrower with a longer photon lifetime. Its peak level could be reduced to the shot noise limit with a shot-noise limited pump source allowing pumping rates larger than 10. Furthermore, we can take advantage of the simple first-order laser response to reduce the low-frequency laser noise with a feedback on the pump signal.

Finally, the high-Q class-A laser noise properties offer new opportunities for many applications including the development of ultra-low noise microwave photonics systems. Further refinements include the investigation of the laser RIN at very low frequencies (below 10 kHz) through microwave signal phase noise measurements and the use of telecom wavelengths compatible with WDM applications. The pump-to-laser RIN transfer should also be explored in the case of optical pumping with summed single-frequency diodes and in the case of electrical pumping of the gain medium. Besides, a fully quantum theoretical description of the system using Heisenberg–Langevin equations for the successive laser modes is under investigation.

REFERENCES


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