Squeezing of intensity noise in nanolasers and nanoLEDs with extreme dielectric confinement

Mørk, Jesper; Yvind, Kresten

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Low-noise light sources are important for on-chip interconnects, sensors, and quantum technology. We show that, using novel cavity designs featuring deep sub-wavelength confinement, it is possible to strongly reduce quantum fluctuations over a large bandwidth. The results could enable integrated sources with extremely low amplitude noise. © 2020 Optical Society of America under the terms of the OSA Open Access Publishing Agreement

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A source of squeezed light with intensity fluctuations substantially below the Poisson limit of laser light would empower, if not revolutionize, several areas of applications. This includes continuous-variable quantum communication and information, quantum-enhanced sensing, and ultra-low power on-chip communications [1–4]. While squeezing on the order of 10 dB has been experimentally demonstrated [3], an integrated and scalable source of squeezed light does not yet exist. Here, we show that cavity-enhanced nanoLEDs employing new designs for deep sub-wavelength confinement [5–8] enable operation in a regime of strong Purcell-enhanced light–matter interaction with more than 20 dB of intensity noise squeezing within a bandwidth of several gigahertz. We extend a recently developed stochastic approach [9] to the case of squeezed-light generation in lasers and LEDs, and use it to analyze the physics of squeezing, as well as establishing the limits to the squeezing that can be achieved. In contrast to recent theoretical work [9–12], we here take into account the quantum statistics of the cavity out-coupling process, which is essential in order to describe intensity noise squeezing.

It was already shown by Yamamoto et al. that a semiconductor laser driven in constant-current mode can generate sub-Poissonian intensity-squeezed light [13,14]. This exploits the natural anti-bunching of electrons, which may be enhanced by Coulomb-blockade effects [15] or nonlinearities in the pump process [16], to realize a sub-Poissonian, i.e., “quiet,” stream of input electrons, which the laser subsequently converts to squeezed light. The degree of squeezing achieved in semiconductor lasers and LEDs so far [17–20] has been severely limited by a number of factors, including weak side modes [21], intrinsic losses in the laser cavity [22,23], and current leakage [24]. Furthermore, intensity-squeezed output has so far been observed within a relatively small bandwidth, which is typically insufficient for transmission of data. Most of these limitations may be overcome by strongly increasing the light–matter interaction. This can be accomplished using recent cavity designs featuring intensity hot spots corresponding to mode volumes that are orders of magnitude below the diffraction limit [5–8,25]. Common to these designs is a central bow-tie geometry that spatially localizes the field, while the surrounding structure provides temporal confinement (Fig. 1). Recent progress in nanolasers [12,26–28] and nanofabrication [8,25] demonstrates the feasibility of fabricating such structures. The intensity noise squeezing factor that can be obtained in nanoLEDs is approximately given by $S = 1/(1 - \eta)$. Here, $\eta$ is an effective quantum efficiency, which for nanoLEDs is limited by the spontaneous emission factor $\beta$, giving the fraction of light spontaneously emitted into the cavity mode. Microcavity lasers with $\beta = 0.97$ [27] have already been demonstrated, and achieving a near-unity $\beta$-factor is also a goal for single photon sources, further showing the relevance of our approach.

We model the dynamics of nanolasers and nanoLEDs by the following set of rate equations [9] for the number of excited emitters, $n_e(t)$, and the number of photons in the cavity, $n_p(t)$:

$$\frac{dn_e}{dt} = R_p(t) - \gamma_r(2n_e - n_0)n_p - \gamma_t n_e,$$  \hspace{1cm} (1)

$$\frac{dn_p}{dt} = \gamma_r(2n_e - n_0)n_p + \gamma_e n_e - \gamma_t n_p.$$  \hspace{1cm} (2)

Here, $n_0$ is the total number of emitters (quantum dots), $\gamma_r$ is the coupling rate between a single emitter and a photon in the cavity mode, which can be expressed in terms of fundamental material and cavity parameters [9], $\gamma_t$ is the total emitter decay rate, and $\gamma_e = \gamma_{out} + \gamma_{in}$ is the cavity decay rate, with $\gamma_{out}$ being the coupling rate into the output channel and $\gamma_{in}$ being all other cavity losses, e.g., due to residual absorption and disorder.

The term $R_p(t)$ is the rate at which emitters are excited by the injected current. Its description depends on how the laser is pumped. For conventional pumping, realized by a current source with classical Poissonian distributed shot noise, we take $R_p(t) = R_{p,\text{conv}}(t) = \gamma_p (n_0 - n_e(t))$. This expression is appropriate for a finite number of emitters, where state-filling dictates $n_e \leq n_0$, and the injection efficiency decreases as more emitters are excited [9]. The total external pump rate is given by $\gamma_p n_0$, with the
For the case of quiet pumping, \( R_p(t) = R_{p,\text{quiet}}(t) \), the pump is a periodic injection of electrons at a constant rate, corresponding to sub-Poissonian electron statistics and unity injection efficiency.

We analyze the quantum noise by generalizing a recently developed stochastic approach \cite{9} to include cavity out-coupling as an additional stochastic process as well as the injection of a sub-Poisson electron average. For the latter, electrons are injected with period \( 1/R_{p,\text{quiet}} \), requiring that \( (R_{p,\text{quiet}} \Delta t)^{-1} \), where \( \Delta t \) is the time-step of the stochastic algorithm, is an integer \( M \geq 1 \). We typically use \( M \sim 10^2 \) and check that the noise measures have converged. We compare the stochastic results to a Langevin approach, where white-noise sources, \( F_\text{w}(t) \) and \( F_\text{p}(t) \), are added to the right-hand side (RHS) of Eqs. (1) and (2), with the magnitude of the noise being described through diffusion coefficients, \( 2D_{xy} = \langle F_\text{w}(t) F_\text{w}(t') \rangle \) (see, e.g., Ref. [9]).

We first consider the case of a nanolaser, using parameter values representative of the quantum dot photonic crystal laser investigated in Ref. \cite{27}: \( n_0 = 50, \beta = 0.97, \gamma_1 = 5 \times 10^9 \text{s}^{-1}, \) and \( \gamma_{\text{em}} = 0.75 \times 10^{11} \text{s}^{-1} \) (corresponding to a cavity \( Q \)-factor of 16,215 at the considered wavelength of \( \lambda = 1.55 \mu\text{m} \)). The histograms in Fig. 2 show detected photon number distributions for \( R_p = 1 \times 10^{-12} \text{s}^{-1} \) for (a) conventional and (b) quiet pumping. If the laser is used for generating low-noise pulses (“bits”) at a bit rate of \( B \), the relevant noise to consider corresponds to having a photodetector with an integration time of \( T_p = 1/B \). In Fig. 2, we consider a bit rate of 2.5 Gb/s and model the detection by a temporal (square) filter of duration 400 ps.

With noisy pumping, Fig. 2(a), the simulated distribution is well fitted by a Gaussian probability distribution, with a FWHM close to that of the corresponding Poissonian distribution with the same mean. On the other hand, quiet pumping, Fig. 2(b), results in a distribution that is clearly sub-Poissonian. The lower average photon number obtained for conventional pumping is due to reduced pump efficiency, which can be compensated for by using a larger pump, but has no effect on the conclusions.

In order to get insight into the mechanism of squeezing, Fig. 3 shows the relative intensity noisy (RIN) spectra for noisy and quiet pumping, comparing external, measurable spectra to the intracavity spectra. The analytical predictions are obtained by a small-signal analysis of Eqs. (1) and (2), including Langevin noise terms, and are given in Supplement 1.

Simulated and analytical RIN spectra in general agree very well. In the case of shot-noise-limited pumping, Fig. 3(a), the RIN of the out-coupled laser signal is reduced by 3 dB compared to the intracavity laser power at low frequencies and approaches the standard quantum limit of \( 2/\hbar \omega_\text{w} \) at frequencies exceeding the characteristic rates describing the laser dynamics. The internal RIN continues to decrease with frequency due to the filtering imposed by the laser dynamics, as expressed by the response function \( |\mathcal{H}(\Omega)|^2 \) (see Supplement 1).

Figure 3 clearly shows that the out-coupling process reduces the intensity noise at low frequencies, for both conventional and quiet pumping. This may appear surprising considering the random nature of the out-coupling process: individual photons are either transmitted through the laser output mirror or reflected back into the cavity. Yamamoto \textit{et al.} explain the noise reduction as being due to destructive interference with external vacuum fluctuations \cite{14}. Alternatively, the squeezing may be understood as the result of the lossless transfer of the sub-Poissonian stream of input electrons to a sub-Poissonian stream of output photons \cite{29}. At low frequencies, corresponding to long observation times, it is guaranteed that a given number of electrons in the input electron stream results in the same number of photons in the output. In contrast, the intracavity energy can be distributed among the photons and the emitters, and the fluctuations of each of these populations may be large.

The black solid line in Fig. 3 shows the effect of increasing the internal losses to 2 cm\(^{-1} \) (\( \gamma_{\text{int}} = 1.71 \times 10^{10} \text{s}^{-1} \) and \( Q = 71,000 \)), corresponding to an out-coupling efficiency of 0.77. Although a small loss, it is seen to strongly reduce the squeezing. Since the small gain region in nanolasers dictates a small out-coupling loss, to achieve lasing, internal losses make it difficult to realize nanolasers with strong noise suppression. Instead, we shall show that nanoLEDs allow to overcome this limit.

Figure 4 shows photon number distributions for a nanoLED in bit slots of duration 100 ps, corresponding to operation at 10 Gb/s, and Fig. 5 shows the corresponding RIN spectra. The parameters rate \( \gamma_p n_r \), representing leakage. A more accurate model could take into account an upper pump reservoir, e.g., constituted by carriers in the wetting layer of a quantum dot laser, but the present model already elucidates important physics related to the squeezing process. For the case of quiet pumping, \( R_p(t) = R_{p,\text{quiet}}(t) \), the pump is a periodic injection of electrons at a constant rate, corresponding to sub-Poissonian electron statistics and unity injection efficiency.

We analyze the quantum noise by generalizing a recently developed stochastic approach \cite{9} to include cavity out-coupling as an additional stochastic process as well as the injection of a sub-Poisson electron average. For the latter, electrons are injected with period \( 1/R_{p,\text{quiet}} \), requiring that \( (R_{p,\text{quiet}} \Delta t)^{-1} \), where \( \Delta t \) is the time-step of the stochastic algorithm, is an integer \( M \geq 1 \). We typically use \( M \sim 100 \) and check that the noise measures have converged. We compare the stochastic results to a Langevin approach, where white-noise sources, \( F_\text{w}(t) \) and \( F_\text{p}(t) \), are added to the right-hand side (RHS) of Eqs. (1) and (2), with the magnitude of the noise being described through diffusion coefficients, \( 2D_{xy} = \langle F_\text{w}(t) F_\text{w}(t') \rangle \) (see, e.g., Ref. [9]).

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With noisy pumping, Fig. 2(a), the simulated distribution is well fitted by a Gaussian probability distribution, with a FWHM close to that of the corresponding Poissonian distribution with the
used are: $N_0 = 10$, $\beta = 1$, $\gamma_c = 1 \times 10^{12} \text{s}^{-1}$, $\gamma_t = 6 \times 10^{13} \text{s}^{-1}$, and $\gamma_{int} = 4.3 \times 10^{10} \text{s}^{-1}$. The large emitter–cavity coupling rate, $\gamma_c = \beta \gamma_t$, reflects that the emitters are embedded in a sub-wavelength cavity (see Supplement 1).

The nanoLED RIN spectra show less frequency dependence than the nanolaser due to the large emitter–cavity coupling rate and large cavity decay rate. The latter also results in a small population of intracavity photons, and the transfer function has an amplitude smaller than one, meaning that the noise is not enhanced beyond the value given by the standard quantum limit, in contrast to the nanolaser, cf. Fig. 3. The reduced squeezing seen at low frequencies for the simulated results as compared to the analytical results is attributed to dynamical pump blocking and the resulting leakage. Figure 5 also shows the RIN spectrum for a reduced $\beta = 0.9$, and the black dashed line shows the standard quantum limit.

Nanolasers may also employ sub-wavelength cavity designs, and to compare with nanolEDs, Fig. 6 shows different measures of the quantum noise (defined in Supplement 1) in dependence of $\gamma_t$. As $\gamma_t$ increases, a transition occurs from lasing to LED operation. This is evidenced by the zero-delay intensity correlation for the intracavity photon distribution, $g^{(2)}(0)$, which changes from a value of one, indicating coherent light, to a value approaching two, indicating thermal statistics. The variation of the RIN with $\gamma_t$ clearly shows the advantage of operating in the nanoLED regime. Whereas the RIN quickly saturates for conventional pumping, it decreases with $\gamma_t$ for quiet pumping.

The correlation of photons detected externally, with a response time of 100 ps, $g^{(2)}(0)$, hardly changes with $\gamma_t$ and is almost independent of the statistics of the pump. From the relation

$$
\langle \Delta n_p^2 \rangle = \langle g^{(2)}_{\text{ext}}(0) - 1 \rangle (n_p^2 + 2)$$

it is thus seen that if the average photon number is large (corresponding, e.g., to a long detector response time), even a small reduction of $g^{(2)}_{\text{ext}}(0)$ below unity indicates fluctuations that are strongly sub-Poissonian. On the other hand, Mandel’s $Q$-factor [31] (see Supplement 1 for the definition) is seen to provide a sensitive measure of squeezing, with a value of $\sim 1$ indicating the maximum degree of squeezing.

The degree of squeezing can be represented by an effective Fano factor. Consider a quiet input electron stream consisting of a sequence of time (bit) slots of duration $T_B$, each containing exactly $N$ electrons (number states). If these number states are transferred to the detector with efficiency $\eta$, the resulting probability distribution is a binomial counting distribution with mean $\langle n \rangle = \eta N$, variance $\sigma^2 = \eta(1 - \eta)N$, and Fano factor $F = \sigma^2 / \langle n \rangle = 1 - \eta$ [2]. It is easily seen that in this case, the noise is reduced by a factor of $S$, with $1/S = \RIN_{\text{quiet}} / \RIN_{\text{conv}} = F$, so that the electron number states are conserved, and the output shows no intensity fluctuations for $F = 0 (\eta = 1)$.

The efficiency $\eta$ has different contributions: $\eta = \eta_c \eta_{\text{dyn}} \eta_{\text{det}}$, with $\eta_c$ being the efficiency by which electrons are converted into out-coupled photons, $\eta_{\text{dyn}}$ the dynamical transfer efficiency, depending on the bit duration $T_B$, and $\eta_{\text{det}}$ the detector efficiency including link losses (assumed to be unity in the simulations). For a laser operating well above threshold, $\eta_c = \eta_c(\gamma_c - \gamma_{\text{int}}) / \gamma_c$, while for an LED, $\eta_c = \eta_c \beta$. Here, $\eta_c$ is the pump injection efficiency, which can be close to unity [32], and $\eta_{\text{dyn}} \sim (H(\xi / T_B))^2$, where $\xi$ is a number of order unity depending on the temporal waveform and the variation of the RIN within the signal bandwidth. This implies that a modulation bandwidth considerably exceeding the signal bandwidth is needed to have significant squeezing. For lasers, a large bandwidth is normally achieved by operating the laser with a high rate of stimulated emission. Due to the small gain, however, nanolasers tend to be overdamped [10,33]. On the other hand, Purcell-enhanced nanolEDs may feature a large modulation bandwidth [34] when exploiting extreme dielectric confinement [7]. This is key to generating squeezed light within a bandwidth of several gigahertz.

The pumping rate for the nanoLED corresponds to a current density less than 10 kA/cm², which is high, but feasible [35]. Compared to plasmonic structures, the dielectric structures considered here have no intrinsic losses, and the high-efficiency nanoLED structures can feature low or even negative heat dissipation, due to the energy carried by the photons [32].

A very important application of low-noise sources is on-chip interconnects, e.g., between computer cores [36]. Figure 7 shows the number of photons required in the 1-bit of an on-off keying system versus $F = 1 - \eta$ to achieve bit-error ratios (BERs) of...
Fig. 7. Average photon number in one bit to reach given BER values versus overall efficiency. $F = 1$ corresponds to Poissonian statistics. Solid curves are analytical estimates from Ref. [2].

$10^{-12}$ and $10^{-20}$ for a decision threshold at the one-photon level (lower curves) and at the midpoint (upper curves). Notice that communication links between the cores of a computer requires much lower BER than conventional communication systems. The curves show the very significant reduction in power level that can be obtained by reducing the Fano factor.

In conclusion, we have shown that nanoLEDs are promising sources of intensity-squeezed light when pumped by a quiet current source and exploiting new designs for extreme dielectric confinement. We have established fundamental limitations to the degree of squeezing and have shown that nanoLEDs are advantageous compared to nanolasers. Furthermore, an efficient approach for stochastic simulations of squeezing was introduced, which may be used for further investigations.

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See Supplement 1 for supporting content.

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