Excess quantum noise in optical parametric chirped-pulse amplification

Cristian Manzoni,1,* Jeffrey Moses,2 Franz X. Kärtner,2,3 and Giulio Cerullo1

1IFN-CNR, Dipartimento di Fisica, Politecnico di Milano, Piazza L. da Vinci 32, 20133 Milano, Italy
2Department of Electrical Engineering and Computer Science and Research Laboratory of Electronics, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA
3DESY - Center for Free-Electron Laser Science and Department of Physics, Hamburg University, D-22607 Hamburg, Germany

*cristian.manzoni@polimi.it

Abstract: Noise evolution in an optical parametric chirped-pulse amplifier (OPCPA) differs essentially from that of an optical parametric or a conventional laser amplifier, in that an incoherent pedestal is produced by superfluorescence that can overwhelm the signal under strong saturation. Using a model for the nonlinear dynamics consistent with quantum mechanics, we numerically study the evolution of excess noise in an OPCPA. The observed dynamics explain the macroscopic characteristics seen previously in experiments in the practically important saturation regime.

© 2011 Optical Society of America

OCIS codes: (320.7110) Ultrafast nonlinear optics; (230.4480) Optical amplifiers; (190.4970) Parametric oscillators and amplifiers; (030.6600) Statistical optics.

References and links

find in this study, the amplification properties of an OPCPA uniquely affect its noise intensity fluctuations of the pump laser [15,16], affects the amplified pulse contrast. As we process, together with the amplified stimulated emission (ASE) contributions or other atoms and molecules [13].

Like any phase insensitive optical amplifier, OPCPAs are subject to amplified spontaneous emission, often called parametric superfluorescence (PSF), i.e., parametric amplification of the quantum noise due to two-photon emission from a virtual level excited by the pump field and stimulated by the signal and idler field zero-point fluctuations [14]. This process, together with the amplified stimulated emission (ASE) contributions or other intensity fluctuations of the pump laser [15,16], affects the amplified pulse contrast. As we find in this study, the amplification properties of an OPCPA uniquely affect its noise intensity fluctuations of the pump laser [15,16], affects the amplified pulse contrast. As we process, together with the amplified stimulated emission (ASE) contributions or other...
performance. In a conventional CPA, in which a laser medium amplifies a chirped signal 
pulse, a population inversion of localized emitters provides homogeneously broadened gain 
that saturates with the pulse fluence. In contrast, an OPCPA is based on instantaneous second-
order nonlinear processes, and the excited virtual levels travel with the pump pulse at a group 
velocity closely matching that of the signal; the virtual levels therefore saturate 
instantaneously on the signal’s retarded time frame. Compared to a standard OPA, an OPCPA 
adds the complexity of a map of instantaneous frequency to temporal coordinate during 
amplification due to the strong linear chirp which allows for inhomogeneous saturation of the 
available parametric gain.

Experiments have shown that amplification in the presence of PSF in an OPCPA results in 
an amplified signal field with two macroscopic components showing different energy 
localization: (i) a coherent pulse with well defined temporal chirp matching that of the 
injected signal pulse (the “seed”) and which therefore can be compressed to generate a Fourier 
transform-limited pulse; (ii) an incoherent pedestal with phase statistics similar to that of PSF, 
which remains at picosecond duration when the signal pulse is recompressed [17,18]. 
Henceforth, we refer to these phenomenological components observed at the output of an 
OPCPA as the coherent pulse and incoherent pedestal, respectively. Especially in the case of a 
broadsband signal and high gain, a severe impact of PSF on noise performance is also well 
documented [5,6,19]: it both degrades the signal stability and places an upper limit on the 
extractable signal energy, due to transfer of pump energy to the incoherent pedestal. The 
dynamics of this energy transfer during amplification have not been observed yet: their 
understanding in the highly nonlinear saturation regime is not only of fundamental interest, 
but is particularly important with regard to performance, since amplifier saturation is 
necessary for obtaining good conversion efficiency and stable output energy. Intuition, based 
on the properties of laser and electrical amplifiers, suggests that saturation should suppress 
fluctuations. However, the effects of saturation on excess quantum noise lack a satisfactory 
description: while the experiments show degradation of pulse contrast during saturation 
[6,19], a numerical analysis of output statistics of PSF in an OPCPA seeded by a distributed 
classical noise source did not isolate the effect [20].

In this paper, we introduce a quantum-mechanically consistent numerical model of the 
dynamics of PSF growth in an OPCPA that captures the process of energy exchange during 
amplification between what will become coherent and incoherent components of the electric 
field after compression, and well reproduces the macroscopic characteristics observed in 
experiments. Since the purpose of this paper is to isolate the influence of PSF on the amplified 
pulse contrast, we do not consider effects due to ASE contributions or other intensity 
fluctuations of the pump laser. By virtue of the model’s adherence to quantum mechanics 
even in the highly nonlinear saturation stage, we observe the saturation dynamics of a 
quantum-noise-contaminated OPCPA, uncovering several distinguishing features. We find 
that PSF must be characterized by two observables which display different evolution 
dynamics: the shot-to-shot energy fluctuation and the ratio of coherent pulse energy and 
incoherent pedestal energy of the amplified signal field. We find that an OPCPA has well 
defined but different operating points for maximum suppression of PSF-induced fluctuations 
or pedestals. Beyond these operating points, heavy saturation leads to large excess noise that 
can be enhanced by orders of magnitude.

2. Numerical model

For the numerical description of the PSF dynamics in the amplification process, we focus on 
an OPCPA seeded by the initial quantum noise field and a chirped signal field; amplification 
occurring in a periodically-poled stoichiometric lithium tantalate (PPLT) crystal. The nonlinear 
quantum system dynamics can be described by a quasi-probability distribution, such as the 
Wigner distribution (WD) [21]. It is well known from quantum optics that for linear systems 
the evolution equation for the WD is equivalent to a classical Fokker-Planck equation, and is 
thus also equivalent to a stochastic process involving classical noise sources, resulting in a 
semiclassical picture of the quantum process. This correspondence has been exploited in
numerical studies of PSF, OPA, and optical parametric oscillation in their linear regimes [22,23]. It is less well known that for the case of weak nonlinearities, i.e., no significant nonlinear effects at the few-photon level, the Fokker-Planck approximation holds and the nonlinear quantum system dynamics can still be extracted accurately from stochastic Langevin equations [24,25], an approach used earlier to study the quantum noise in parametric amplifiers used for squeezed light generation [26,27]. These stochastic equations have a deterministic component equal to the Heisenberg equations of motion for the field operators and are complemented by relaxation terms and associated noise terms. For the case of a lossless OPA process, fluctuations stem solely from the quantum mechanical uncertainty in the input fields. Knowledge of a quasi-probability distribution allows computation of all expectation values of quantum mechanical observables, and for the case of the WD and its associated stochastic process, computed expectation values correspond to quantum mechanical expectation values of symmetrically ordered field operators [21]. Thus, this approach allows for a rigorous treatment of quantum fluctuations in weakly nonlinear quantum optical systems, such as OPCPAs with large mode cross sections. For completeness, we also mention that the quantum dynamics of a second-order nonlinear process, as is the case discussed here, can be described exactly with the help of the positive P-representation, pioneered by Peter Drummond [28]. It has been shown [29] that third- and higher-order moments of the electric field can differ significantly whether calculated by means of the P or truncated WD representation. However, for large normalized photon numbers and weak nonlinearity, differences are very small (the discrepancy was quantitatively small for pump photon numbers of 100 in [25], though some slight differences persisted at early times.). With the gigawatt peak powers typical of OPAs, we are always in the high photon number (totaling \(10^9\) in our case) and weak nonlinearity limit. This fact, together with the increased mathematical complexity of the positive P-representation, using a twice as large phase space which considerably increases computation time, led us to work within the truncated WD.

We simulate the evolution of noise in an OPCPA by numerically solving the coupled nonlinear equations of parametric amplification in the spectral domain, accounting for linear dispersion to all orders [30,31]; the equations describe the interaction among signal, idler and pump, which in the following will be respectively labelled \(i = 1, 2, 3\). These waves propagate along the \(z\) coordinate with carrier frequency \(\omega_i\) and wavenumber \(k_i\). To exploit the large \(d_{33}\) nonlinear coefficient of the crystal for all fields polarized along the extraordinary axis, we operate in the quasi phase-matching regime, obtained by periodically poling the nonlinear crystal: poling is accounted for by changing the sign of \(d_{33}\) along \(z\). The carrier fields therefore experience at any crystal coordinate a real phase mismatch \(\Delta k = k_3 - k_2 - k_1\). We describe the electric field of each wave as:

\[
E_i(z,t) = \frac{1}{2} \left\{ A_i(z,t) \exp \left[ j(\omega_i t - k_i z) \right] + \text{c.c.} \right\} = \Re \left\{ A_i(z,t) \exp \left[ j(\omega_i t - k_i z) \right] \right\}
\]

where \(A_i(z,t)\) denotes the field complex amplitude. The coupled equations describing the second order interaction of the fields are derived from the nonlinear propagation equation:

\[
\frac{\partial^2 E}{\partial z^2} - \mu_0 \frac{\partial^2 D}{\partial t^2} = \mu_0 \frac{\partial^2 P_{NL}}{\partial t^2}
\]

applied on the total field \(E(z,t) = E_1(z,t) + E_2(z,t) + E_3(z,t)\).

Here \(D_i(z,t) = \int e_i(r) E(z,t-r) d\tau\) is the linear electric induction field accounting for the linear dispersion of the medium [32], and \(P_{NL} = 2\epsilon_0 d_{33} E(z,t)^2\) is the nonlinear polarization, where \(d_{33}\) is the effective second-order non-linear coefficient. Since our model also accounts for a broadband noise field, and our purpose is to calculate the evolution of the noise with the highest accuracy, we avoid the slowly-varying-envelope approximation [33] typically adopted to simplify the calculations; in addition we consider the linear dispersion of the material to all orders. For this purpose, we develop Eq. (2) in the frequency domain by taking its Fourier transform and obtaining:
\[
\frac{\partial^2 \tilde{E}}{\partial z^2} + \frac{\Omega n(\Omega)}{c_0^2} \tilde{E} = -\mu_0 \Omega^2 \tilde{P}_{NL}.
\]  

(3)

where \( \tilde{E}(z, \Omega) = \mathcal{F}\{E(z, t)\} \) and \( \tilde{P}_{NL}(z, \Omega) = \mathcal{F}\{P_{NL}(z, t)\} \) are the non-unitary Fourier transforms of the electric field and of the nonlinear polarization respectively. \( \Omega \) is the angular frequency and \( n(\Omega) \) is the frequency-dependent refractive index. \( P_{NL} \) is developed rejecting components at frequencies different from \( \omega_1, \omega_2 \) and \( \omega_3, \) when the fields \( E_i \) are not overlapped in frequency, wave vector, and polarization, it is possible to split Eq. (3) into three coupled equations which separately describe the evolution of the field envelopes:

\[
\begin{align*}
\frac{\partial^2 \tilde{A}_1}{\partial z^2} - j2k_i \frac{\partial \tilde{A}_1}{\partial z} + b_i^2 \tilde{A}_1 &= -c_i \cdot e^{-j\Delta k z}, \\
\frac{\partial^2 \tilde{A}_2}{\partial z^2} - j2k_i \frac{\partial \tilde{A}_2}{\partial z} + b_i^2 \tilde{A}_2 &= -c_i \cdot e^{-j\Delta k z}, \\
\frac{\partial^2 \tilde{A}_3}{\partial z^2} - j2k_i \frac{\partial \tilde{A}_3}{\partial z} + b_i^2 \tilde{A}_3 &= -c_i \cdot e^{+j\Delta k z}.
\end{align*}
\]

(4)

In this case \( \tilde{A}_i(z, \omega) = \mathcal{F}\{A_i(z, t)\} \) is the Fourier transform of the envelope amplitude of each field and \( \omega = \Omega - \omega_i \) is the detuning from the carrier frequency \( \omega_i \); coefficients \( b_i^2 \) and \( c_i \) are defined as:

\[
b_i^2 = k_i^2 + \tilde{k}_i^2, \quad \text{with} \quad \tilde{k}_i = \frac{(\omega + \omega_i) \cdot n_i(\omega + \omega_i)}{c_0}.
\]

(5a)

and

\[
\begin{align*}
c_1 &= \left( \frac{\omega + \omega_1}{c_0} \right)^2 d_{ei} \mathcal{F}\{A_1 A_2^*\}, \\
c_2 &= \left( \frac{\omega + \omega_2}{c_0} \right)^2 d_{ei} \mathcal{F}\{A_1 A_3^*\}, \\
c_3 &= \left( \frac{\omega + \omega_3}{c_0} \right)^2 d_{ei} \mathcal{F}\{A_2 A_3\}.
\end{align*}
\]

(5b)

Here \( n_i(\omega + \omega_i) \) are the refractive index functions deduced from the Sellmeier equations, and allow to take into account the whole linear dispersion of the material. The system can be solved as follows: if a suitably small step \( \Delta z \) is chosen, the products \( A_i A_i^* \), \( A_i A_j^* \) and \( A_i A_j \) are nearly constant, and Eqs. (4) can be analytically solved. Given the fields \( \tilde{A}_i(z, \omega) \) at the beginning of a step, the fields at \( z + \Delta z \) are:

\[
\begin{align*}
\tilde{A}_1(z + \Delta z, \omega) &\approx \tilde{A}_1(z, \omega) + \frac{c_1}{\gamma_1} \cdot \exp[j(k_i - \tilde{k}_i)\Delta z] - \frac{c_1}{\gamma_1} \cdot \exp[- j\Delta k \Delta z], \\
\tilde{A}_2(z + \Delta z, \omega) &\approx \tilde{A}_2(z, \omega) + \frac{c_2}{\gamma_2} \cdot \exp[j(k_i - \tilde{k}_i)\Delta z] - \frac{c_2}{\gamma_2} \cdot \exp[- j\Delta k \Delta z], \\
\tilde{A}_3(z + \Delta z, \omega) &\approx \tilde{A}_3(z, \omega) + \frac{c_3}{\gamma_3} \cdot \exp[j(k_i - \tilde{k}_i)\Delta z] - \frac{c_3}{\gamma_3} \cdot \exp[+ j\Delta k \Delta z],
\end{align*}
\]

(6)

where \( \gamma_{1,2} = b_{1,2}^2 - \Delta k^2 - 2k_{1,2}\Delta k \) and \( \gamma_3 = b_3^2 - \Delta k^2 + 2k_3\Delta k \). It is important to remark that the representation of fields given in Eqs. (4) assumes that pump, signal and idler are three...
separate fields. Justifying this treatment, the amplifier we model employs a small non-collinear angle between signal and idler, used both to allow their separation after amplification (since they have opposite temporal chirp) and to avoid signal-idler interference for preservation of carrier-envelope phase of the signal. The results we obtain from our model do not hold for degenerate collinear OPAs, for which fields 1 and 2 of Eqs. (4)–(6) collapse into one equation, and for strongly non-collinear geometries, which would require at least one more spatial coordinate.

Our 1-D plane wave model includes all longitudinal modes, \( m \), and their associated noise. In the frequency domain, at any mode frequency, \( \omega_m \), the corresponding component of the initial signal, idler, or pump electric field is represented by a complex stochastic phasor, \( A_m(0) = B_m + n_m \). \( B_m \) is the deterministic component of the field, and is set to 0 in the case of the idler; \( n_m \) is a zero-mean, stochastic phasor representing the independent fluctuations of the field.

\[
A_m(0) = B_m + n_m
\]

Note that fluctuations are included for each of the signal and idler fields; the quantum noise of the pump is negligible, as we also confirmed independently by simulations not reported here. Real and imaginary components of \( n_m \) are taken as uncorrelated Gaussian distributions [22] with variance \( \sigma_m \propto \omega_m \). A representation of these fluctuating fields is given in Figs. 1(a) and 1(b), where we show \( B_m \) (vectors), and \( A_m \) and \( n_m \) (scatter) for three modes \( \omega_m \) of the signal field. Our numerical method treats identically the initial electric field components whether originating from the deterministic field or the vacuum fluctuations. We apply this tool to the study of an ultra-broadband OPCPA system known to be sensitive to PSF [5,6]. We model a typical high-gain pre-amplifier, in which noise begins at the vacuum level and that establishes the noise content of later stages of a multi-stage system. The amplifier, pumped by a 9-ps FWHM Gaussian pulse at 1.047 μm and seeded by a broadband (69-THz FWHM) pulse at 2.094 μm for operation around degeneracy, uses a 3-mm long...
PPSLT crystal with poling period $\Lambda = 31.2\mu$m. These parameters are close to the experimental conditions of Ref. [6]. The pump-to-seed energy ratio is $10^6$.

Table 1. Key Parameter for Three OPCPA Configurations

<table>
<thead>
<tr>
<th>Config</th>
<th>Input ($z = 0$ mm)</th>
<th>Output ($z = 3$ mm)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Seed duration [ps]</td>
<td>Pump intensity [GW/cm$^2$]</td>
</tr>
<tr>
<td>I</td>
<td>1.04</td>
<td>6.254</td>
</tr>
<tr>
<td>II</td>
<td>7.34</td>
<td>8.356</td>
</tr>
<tr>
<td>III</td>
<td>10.48</td>
<td>8.825</td>
</tr>
</tbody>
</table>

*For efficiency $\eta$, bandwidth $\Delta\nu$ and efficiency-bandwidth product columns, maximum values are highlighted in bold.

The variances $\sigma_m^2$ are determined by the quantum fluctuations due to the longitudinal modes of the 100-ps-long simulation window. This number is further increased by a factor equal to the number of transverse modes amplified assuming a pump beam of 100-µm radius in the 3-mm long crystal, which is estimated as about 25. We note that this choice results in a calculated amplified pulse contrast that closely matches that observed in equivalent experiments [6]. This system was investigated previously in order to find the conditions of maximum efficiency-bandwidth product [34]. Following that analysis, we choose 3 values of seed chirp, summarized in Table 1, which correspond to: an under-chirped amplifier, with maximum amplified signal bandwidth but limited conversion efficiency (Configuration I); an amplifier chirped for maximum efficiency-bandwidth product (Configuration II); and an over-chirped amplifier (Configuration III), with excellent conversion efficiency but significant spectral narrowing. Table 1 also provides the pump intensity corresponding to each of the three regimes. For each configuration, we evaluated 50 independent trajectories triggered by uncorrelated noise fields. The averages taken over this ensemble of classical solutions correspond to quantum-mechanical expectation values [35]. The results of a batch of simulations are depicted in Fig. 1(c).

The incoherent nature of the amplified noise is evidenced in panel (d), representing the same fields after compression. A complete description of an amplified signal field is given in Fig. 2, showing the WD before (a) and after (b) compression. The WDs clearly reveal the presence of a strong incoherent field, arising from PSF, superimposed to the coherent chirped amplified signal; as a feature common to all configurations, this component has a spectrum corresponding to the phase-matching bandwidth of the OPCPA, and duration comparable to...
the pump pulse. After applying a linear chirp to compress the coherent pulse (see Fig. 2(b)),
the incompressible temporal pedestal degrades the pulse contrast, as observed in experiments.

In addition to this pedestal, PSF also strongly affects the shot-to-shot energy stability of
the amplified signal. The noise contaminating the signal field during the OPCPA process can
thus be characterized by two quantities: (i) the signal-to-noise ratio (SNR), which evaluates
the shot-to-shot energy fluctuation of the signal pulse; (ii) the signal-to-pedestal ratio (SPR),
which measures the ratio between the energies of the coherent pulse and the incoherent
pedestal. To calculate the SNR we must derive the expectation value and variance of a mode
intensity $I_m(z) = a_m^\dagger(z)a_m(z)$ of signal or idler, (where $a_m^\dagger(z)$, $a_m(z)$ are the photon creation
and annihilation operators), in terms of expectation values of the classical stochastic field
variables $A_m(z)$. Equating the computed expectation values of the stochastic field
variables to the expectation values of the corresponding symmetrically ordered quantum
mechanical field operators, we derive:

$$
\langle I_m(z) \rangle = \langle |A_m(z)|^2 \rangle - \langle |\hat{p}_m|^2 \rangle
$$

Equations (7) give us a physically satisfying answer for the initial field intensity and
variance: $\langle I_m(0) \rangle = |B_m|^2$, i.e., the deterministic component of the initial field, and
$[\Delta I_m(0)]^2 = 2|B_m|^2 \langle |\hat{p}_m|^2 \rangle$, i.e., the corresponding photon number fluctuations or shot noise.

These equations also correctly give a null expectation value and null variance for a vacuum
state ($B_m = 0$): this is of key importance, since the simulation is initially dominated by modes
in the vacuum state. From these quantities we can calculate the expectation value $\langle E(z) \rangle$ and variance $\Delta E(z)$ of the pulse energy:

$$
E(z) = \sum_m I_m(z) \Delta \omega_m
$$

$$
[\Delta E(z)]^2 = \langle (E(z) - \langle E(z) \rangle)^2 \rangle
$$

and define $\text{SNR}(z) = \langle E(z) \rangle / \Delta E(z)$. Figure 3 shows the evolution with propagation length of
$\langle E(z) \rangle$ and $\Delta E(z)$ for signal and idler, while Fig. 4(a) shows $\text{SNR}(z)$ (solid lines). The
pedestal field is deduced from the comparison between the PSF-contaminated signal field
$A_m(z)$ and a reference field $A_{\text{ref}}(z)$ obtained from the simulation of a noise-free amplifier ($n_m = 0$). In Figs. 1(c) and 1(d) the dots correspond to the PSF-contaminated signal fields and the
arrows to the deterministic reference fields. We find that PSF does not significantly modify
the amplification regime, in terms of pump depletion; the pedestal field $A_{\text{ref}}(z)$ can be
therefore evaluated as $A_{\text{ref}}(z) = A_m(z) - A_{\text{ref}}(z)$. The SPR is finally evaluated by taking the ratio
between the signal and the pedestal pulse energy.

It should be noted that, while the SNR can be measured by the shot-to-shot fluctuations of
the amplified pulses, the SPR is difficult to characterize experimentally. In fact, the effect of
PSF cannot be isolated by measuring an unseeded OPCPA, since pump depletion is much
lower than in the seeded case. The experimental measurement of the evolution of these
quantities during the OPCPA process poses an even greater challenge. Our numerical
approach, on the other hand, has the unique capability of isolating the pedestal field from
the signal under realistic conditions, allowing exploration of its dependence on amplifier
parameters, as well as SNR and SPR evolution in the OPCPA crystal.
3. Results and discussion

In Fig. 3 we report the evolution of energy mean and standard deviation of signal and idler, for configurations I and II; for completeness, we have extended the calculation beyond the optimum crystal length of 3 mm. From the curves we can clearly identify three stages of amplification: growth, saturation, and over-saturation.

![Figure 3](image_url)

Fig. 3. (Color online) Evolution of energy mean and standard deviation for signal and idler; the energy growth of the pedestal mean is also given. In the case of configuration II we indicate the coordinate at which the pump peak is fully depleted (see inset). Trends calculated for configuration III (not shown) are comparable to the ones of configuration II.

At the start of amplification, the unseeded idler intensity catches up with the signal and the signal standard deviation grows relative to the mean as a result of mixing with the idler, causing a degradation of the initial SNR [Fig. 4(a)]. After this, exponential growth of both waves sets in. In this regime we observe a gradual decay of SNR.

![Figure 4](image_url)

Fig. 4. (Color online) (a) Evolution of SNR and SPR for the three configurations of Table 1. (b) Temporal profiles of the signal pulses after compression, each normalized to its peak.

Once we enter the pump depletion regime, however, the signal energy fluctuation is dramatically reduced, thus improving the SNR. The SNR reaches a sharp maximum shortly before the exit facet of the crystal, close to the propagation distance for which the pump peak fully depletes and the maximum pump to signal/idler conversion occurs. After this point, a rapid drop in SNR is seen. This clearly shows that, while in the pump depletion regime the SNR can be considerably enhanced, in over-saturation the mean intensity sags due to back-conversion to the pump, while noise continues to grow on average across the pulse. Figure
4(a) also shows that, while the three configurations exhibit very similar trends up to saturation, configuration II is preferable due to its higher SNR value.

After saturation, the low-chirp configuration results in the most dramatic drop in SNR, a result of strong growth of PSF in the unseeded temporal wings while the amplified signal converts back to pump at the pulse center. In comparison, the SPR trend reveals dramatic differences with that of SNR, indicating that the energy of the incoherent pedestal, evaluated through the SPR, is not directly correlated to the pulse energy fluctuations measured by the SNR. In fact, maximum suppression of energy fluctuations in configuration I occurs simultaneously with strong degradation of pulse contrast, and in this case the SPR is sharply reduced at the peak of saturation, i.e., at $z = 3$ mm, where conversion efficiency and efficiency-bandwidth product are maximized. The two parameters therefore offer complementary information for evaluating the impact of excess noise in the OPCPA. The profile of the incoherent pedestal after pulse compression is given in Fig. 4(b); as expected, configuration I is strongly contaminated by an incoherent pedestal generated at unseeded temporal coordinates. Note that the different pedestal durations are due to pulse compression, which imparts a vertical shearing to the WDs (see Fig. 2). These observed dynamics explain the experimental results of Ref. [6], including the non-intuitive result that a saturated amplifier can exhibit reduced pulse contrast. We note, the observed degradation of SPR in saturation, and of both SPR and SNR in oversaturation, are characteristic phenomena singular to quantum noise growth in an OPCPA. In a conventional CPA, there is no oversaturation regime; while nonlinear wave mixing in an OPCPA allows back-conversion of signal and idler to pump after amplification, saturation of a population inversion in a laser is not reversible. Laser gain saturation is also homogeneous, and, therefore, unlike the OPCPA, unseeded frequency modes of the signal field cannot continue to amplify exponentially while the gain for seeded modes saturates. In a conventional OPA, the absence of a chirp ensures that saturation and back-conversion is uniform with respect to frequency, and no macroscopic pedestal is produced, thus allowing SNR to fully characterize the noise performance.

4. Conclusions

In conclusion, we have performed a quantum-mechanically consistent numerical investigation of the dynamics of PSF growth in a realistic high-gain OPCPA. Thanks to the model’s capability to simulate also the saturation stage of amplification, this investigation for the first time captures all dynamics of a quantum-noise-contaminated OPCPA, addressing the important practical issues of signal energy stability and pulse contrast. Both quantities are related to the incoherent pedestal superposed to the signal, but display different evolution dynamics throughout the amplification process. Three operating conditions were explored, characterized by different chirps of the input seed and maximizing, respectively, the bandwidth, the efficiency-bandwidth product, and the conversion efficiency. We find that the chirp maximizing the efficiency-bandwidth product is also characterized by the smallest contribution of the noise, both in terms of energy fluctuation and of pulse contrast. Significantly, we find that while amplifier saturation improves the signal’s shot-to-shot energy stability, it does not necessarily improve the pulse contrast. In fact, for the case of an amplifier optimized for bandwidth, strong degradation of the pulse contrast (i.e., growth of the PSF-derived incoherent pedestal energy relative to the coherent signal pulse energy) is observed in saturation. Over-saturation, sometimes employed to boost amplifier bandwidth at the expense of conversion efficiency, uniformly reduces both pulse contrast and energy stability.

Knowledge of these dynamics thus provides important insight for the optimization of OPCPA systems applied to the study of strong-field laser physics, as well as increases our fundamental understanding of quantum noise in parametric amplification.

Acknowledgments

This work was partially supported by the U.S. Air Force Office of Scientific Research (AFOSR) under grants FA8655-09-1-3101, FA9550-09-1-0212 and FA9550-10-1-0063, and the Progetto Rocca.