Noise amplification during supercontinuum generation in microstructure fiber

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Supercontinua generated by femtosecond pulses launched in microstructure fiber can exhibit significant low-frequency (<1-MHz) amplitude noise on the output pulse train. We show that this low-frequency noise is an amplified version of the amplitude noise that is already present on the input laser pulse train. Through both experimental measurements and numerical simulations, we quantify the noise amplification factor and its dependence on the supercontinuum wavelength and on the energy and duration of the input pulse. Interestingly, the dependence differs significantly from that of the broadband white-noise component, which arises from amplification of the input laser shot noise.

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An extremely broad spectrum, or supercontinuum, can be generated by launching of femtosecond pulses from a Ti:sapphire laser into highly nonlinear microstructure and tapered silica fibers.[1,2] The resulting supercontinuum has potential applications in a number of areas, including optical coherence tomography[3] and high-precision spectroscopy.[4] Its most important application to date, though, has been in the area of frequency metrology, where the supercontinuum has permitted a simple and direct link between optical and microwave frequencies.[4,5] The supercontinuum has a number of attractive properties; it is a high-power, broad-spectrum, phase-coherent source with a single spatial mode. Unfortunately, it can also suffer from significant fluctuations in amplitude, which translate to a poor signal-to-noise ratio and can limit its utility. In recent work we discussed the fundamental limit to this amplitude noise that arises from shot noise on the input laser pulse train.[6] In this Letter we discuss an additional, potentially larger, contribution to the amplitude noise on the supercontinuum that arises from technical noise on the input laser pulse train, i.e., noise in excess of the quantum shot noise. This contribution is qualitatively and quantitatively difference from that arising from input shot noise.

At best, one might hope that the noise properties of the input laser pulse train would be transferred to the supercontinuum. In that case, at low Fourier frequencies the noise on the supercontinuum pulse train would mirror the low-frequency technical noise on the input laser pulse train, and at high Fourier frequencies the noise on the supercontinuum pulse train would, like that for the laser, approach the shot-noise limit. Unfortunately, the noise of the supercontinuum is actually much worse; the same basic Kerr nonlinearity that gives rise to the supercontinuum also results in a large amplification of any input noise seed. This amplification is closely related to the results of recent Letters that reported high-contrast fine structure on the single-shot optical supercontinuum spectrum that depended strongly on the input pulse energy.[7,8] In this Letter we present quantitative experimental, numerical, and statistical studies of the amplification of the input laser technical noise during supercontinuum generation. The resulting noise on the supercontinuum will depend on the choice of pump lasers, which largely determines the noise on the femtosecond Ti:sapphire laser.[9]

The noise on the supercontinuum can be parameterized as the relative intensity noise (RIN) measured on a photodetector. To first order, the RIN of a spectral slice of the supercontinuum, R_{SC}, at a rf Fourier frequency f, supercontinuum wavelength λ_0, and bandwidth $\Delta\lambda$ for an input pulse of energy E_0 can be written as

$$R_{\mathrm{SC}}(f, \lambda_0, \Delta\lambda, E_0) = \kappa^2(\lambda_0, \Delta\lambda, E_0) R_L(f), \qquad (1)$$

where $R_L(f)$ is the laser RIN and the square root of the amplification factor or gain κ^2 is given by

$$\kappa(\lambda_0, \Delta\lambda, E_0) = \frac{\int_{\lambda_0-\Delta\lambda/2}^{\lambda_0+\Delta\lambda/2} E_0 \dfrac{\mathrm{d}S(E_0)}{\mathrm{d}E} \mathrm{d}\lambda}{\int_{\lambda_0-\Delta\lambda/2}^{\lambda_0+\Delta\lambda/2} S(E_0) \mathrm{d}\lambda}, \qquad (2)$$

where $S(E)$ is the supercontinuum spectral density per wavelength interval for an input pulse energy of E. The factor κ can be simply interpreted as the fractional change in supercontinuum intensity at a given wavelength for a fractional change in the input pulse energy. Writing the noise amplification in this form clarifies and quantifies its relation to the previous qualitative observations of the strong sensitivity of the spectrum to input pulse energy.[7,8] Equation (1) shows that the nonlinear amplification of the noise is contained within the gain factor κ^2, which is independent of both the laser RIN and the rf Fourier frequency of the noise, under the following assumptions: First, the RIN is assumed to be small enough ($\sim<1\%$) that the linear approximation of Eq. (2) is valid; otherwise higher-order terms must be included in Eq. (1). Second, we assume no interaction between adjacent pulses because of their large temporal separation, which implies that the gain is independent of Fourier frequency for low frequencies. Finally, since the derivation assumes a change in only the overall pulse energy and not in the pulse structure, the rf Fourier frequency must be much less than the reciprocal pulse duration (10 THz). As discussed below, this final assumption is violated for shot noise, with significant consequences.

To simulate the noise amplification numerically, we use the generalized nonlinear Schrödinger equation,[10] with appropriate input laser pulse parameters, fiber dispersion,[1] and a nonlinear coefficiency $\gamma = 70\ \mathrm{W}^{-1}\ \mathrm{km}^{-1}$. For the pulse parameters of interest, a set of five spectra were simulated at pulse energy increments of 0.1%. The gain was then calculated from the derivative of the spectra by use of Eq. (2). This method of calculating the gain was significantly faster than a more brute-force Monte Carlo approach, which requires ~64 or more separate simulations with random Gaussian noise for the same result.

To measure the noise amplification experimentally, we use an apparatus shown schematically in Fig. 1.[6] An Ar^+-pumped Kerr-lens mode-locked Ti:sapphire laser provides pulses with a 45-nm bandwidth centered at 810 nm at a 100-MHz repetition rate. We first direct the laser output through a pair of double-passed fused-silica compensating prisms to control the laser chirp. A portion of the input pulse is directed to an interferometric autocorrelator for measurement of the pulse duration and inference of the pulse chirp, or quadratic spectral phase, in femtoseconds,[2] assuming a sech^2 pulse envelope. The laser output is then focused into a 15-cm-long microstructure fiber.[1] A portion of the supercontinuum that exits the fiber is diverted to an optical spectrum analyzer for measurement of the full spectrum. The remainder is spectrally filtered by a monochromator with an 8-nm bandwidth and detected by an infrared or visible detector. The low-frequency technical noise appears as an approximately Gaussian pedestal with a 440-kHz FWHM about harmonics of the laser repetition rate. The shape of the pedestal is identical to that of the measured technical noise on the input laser pulse train. However, the technical noise on the input laser corresponds to ~0.35% fluctuations in the pulse energy, whereas the measured technical noise on the supercontinuum is typically much larger. The noise amplification is exactly the ratio of these amplitudes.

Figure 2 shows the supercontinuum spectrum and noise amplification at 10-nm increments for both experiment and theory. The gain ranges from 0 to 40 dB, depending strongly on the supercontinuum wavelength. The simulated spectrum has the same general features and spectral width as the measured supercontinuum. The simulated gain has the same mean value and qualitative fluctuations as the measured gain. However, the exact structure of the gain with wavelength is not fully reproduced between simulation and experiment, presumably because it depends strongly on the exact input pulse and fiber dispersion parameters.

For practical purposes, a statistical model of the expected gain versus wavelength is useful. Since a number of independent nonlinear processes contribute to the value of the normalized derivative, κ, we assume that κ follows Gaussian statistics from the central limit theorem with zero mean.[11] The gain, κ^2, therefore has a probability distribution $P(\kappa^2) = \exp(-\kappa^2/2\sigma^2)/(2\pi\kappa^2\sigma^2)^{1/2}$, for which κ^2 has a mean of σ^2 and median of 0.46 σ^2. Figure 3 compares this probability distribution with a histogram of measured gain values, showing good agreement. Preliminary measurements indicate that the correlation of the gain across wavelengths is itself wavelength dependent, with a correlation length ranging from 1 to 10 nm.

Because the gain arises from nonlinear effects, it might depend on the pulse energy, duration, or chirp. Figure 4 gives the dependence of the median RIN, calculated across the supercontinuum spectrum, on the input pulse energy. The gain increases exponentially with pulse energy, while the width of the supercontinuum increases linearly; unfortunately, broader supercontinua come at the cost of higher gain on the

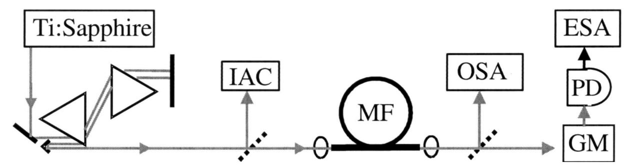

Fig. 1. Schematic of the experiment: IAC, interferometric autocorrelator; MF, microstructure fiber; OSA, optical spectrum analyzer; GM, grating-based monochromator; PD, photodiode; ESA, electrical spectrum analyzer.

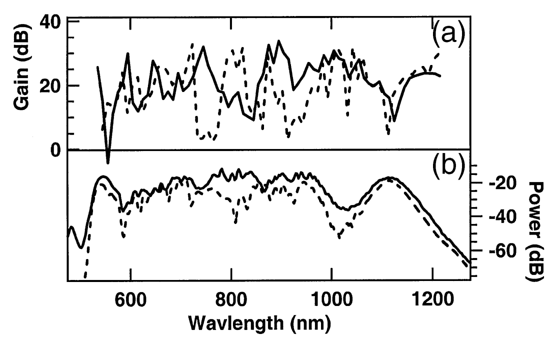

Fig. 2. (a) Noise amplification and (b) supercontinuum spectral power for experiment (solid curves) and theory (dashed curves) for a pulse energy of 0.85 nJ, spectral width of 45 nm, and pulse chirp of $-160\ \mathrm{fs}^2$, corresponding to a pulse duration of 34 fs.

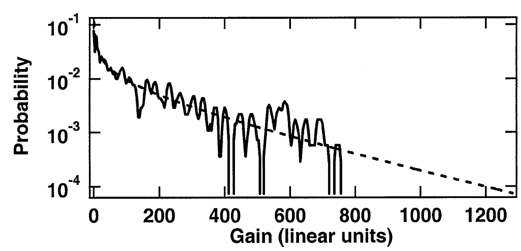

Fig. 3. Probability distribution for the gain assuming Gaussian intensity fluctuations (dashed curve) and from the measured values (solid curve) across the supercontinuum for the range of chirp values given in Fig. 5, below. The analytical distribution had a median gain equal to the measured value of 18 dB.

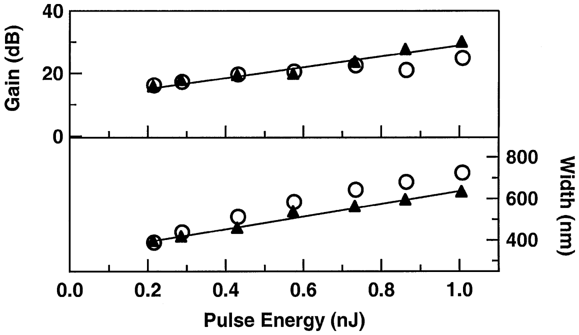

Fig. 4. Dependence of the gain and the −20-dB spectral width on the pulse energy for experiment (solid triangles) and simulation (open circles) for a pulse duration of 47 fs, a spectral width of 42 nm, and a corresponding pulse chirp of $-282\ \mathrm{fs}^2$. The solid lines are fits with a slope of 17 dB/nJ for the gain and 300 nm/nJ for the spectral width.

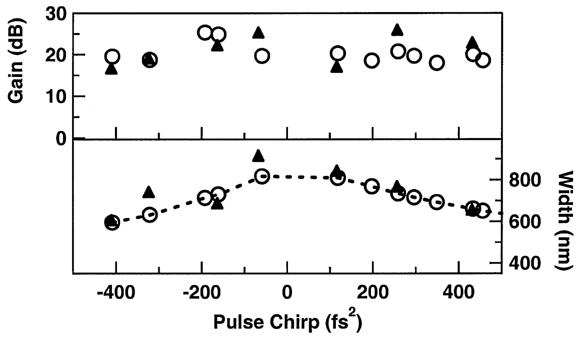

Fig. 5. Dependence of the gain and the −20-dB spectral width on the pulse chirp for experiment (solid triangles) and theory (open circles, dashed curve) for a pulse energy of 0.85 nJ and spectral width of 45 nm. Chirp variation from 0 to $\pm 650\ \mathrm{fs}^2$ corresponds to a range of pulse widths from ~20 to 90 fs FWHM.

laser technical noise. The simulation and experimental measurements are again in good agreement.

The dependence of the spectral width and gain on the pulse chirp is given in Fig. 5 for both the measurements and simulations. The average value and spread of the gain is reproduced well by the simulations. Although the spectral width is broadest for the smallest chirp, there is no dependence of the gain on the input laser pulse chirp, at least within the scatter of values.

It is interesting to compare these results with the analogous results of Ref. 6, which discusses the amplification of the input shot noise under an identical range of pulse parameters. The low-frequency technical noise discussed here corresponds to changes in the overall pulse energy from one laser pulse to the next, and its amplification arises simply from the sensitivity of the nonlinear Schrödinger equation to the input pulse parameters, as stated in Eq. (2). However, the input shot noise, unlike this bandwidth-limited technical noise, has very high-frequency components, which are greater than the reciprocal pulse duration of ~10 THz. These high-frequency components generate random high-frequency ripples across a single laser pulse, which have two important consequences with regard to the noise amplification. First, Eq. (2) is not a valid description of the noise amplification for shot noise, since its derivation assumed no change in the pulse shape. Second, and more importantly, unlike the bandwidth-limited (<1-MHz) technical noise, the high-frequency shot-noise components are subject to envelope modulation instability gain. These modulation instability effects are substantial at these high powers and frequencies and lead to the qualitatively different amplification of shot noise described in Ref. 6. Both amplifications do show a very similar dependence on supercontinuum wavelength, but the similarities end there. The amplification of the technical noise, discussed here, is typically only ~20 dB and has a pulse energy dependence of ~17 dB/nJ (slope of Fig. 4) and low dependence on pulse chirp. The amplification of the shot noise (and spontaneous Raman noise) can range from 45 to 90 dB, has a stronger pulse energy dependence of ~45 dB/nJ and is highly dependent on pulse chirp.[6]

In conclusion, we have characterized the amplification of technical input noise during supercontinuum generation in microstructure fiber and shown its dependence on supercontinuum wavelength, pulse power, and chirp. This work, combined with Ref. 6, explains the strong amplitude noise on the supercontinuum observed by a number of researchers.

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