Ultra-short bound states generation with a passively mode-locked high-power Yb-doped double-clad fiber laser

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Abstract

We report the generation of ultra-short bound states in a high-power ytterbium-doped fiber laser operating in the normal dispersion regime. We theoretically demonstrate that such bound states are stable solutions of the quintic complex Ginzburg–Landau equation.

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1. Introduction

The design of passively mode-locked fiber lasers is now well controlled and it is possible to generate low- or high-energy femtosecond pulses in the normal or anomalous dispersion regimes. However, depending on the experimental set-up, these lasers can operate in very different regimes that can be detrimental for some applications. Much attention has thus to be devoted to the control of the regime of emission. A commonly observed regime is multiple pulsing when increasing the pump power. Another example is the emission of bound states. This latter point has been first pointed out by Malomed et al. [1,2] who theoretically showed that the interaction of slightly overlapping solitons could lead to the formation of bound solitons with a discrete and fixed time separation in coupled nonlinear Schrödinger equations or quintic complex Ginzburg–Landau equation. The phenomenon has been observed recently in low-power passively mode-locked Er-doped fiber soliton lasers [3–6]. Experimentally, the Raman effect and the random phase variation between solitons prevent the formation of bound states in such lasers.
Multiple solitons can be observed but they are generally so far from each other that there is no interaction between them. However, by adjusting the intra-cavity polarization controllers carefully, the phase of the solitons can be locked and the time separation between solitons can be fixed on specific values by the total cavity dispersion. The spectrum is then strongly modulated. The emission is stable over several hours. The main properties of such a bound-state soliton laser have been further detailed [5]. In particular, bound solitons have been shown to be another intrinsic feature of the laser, and no single-pulse soliton can coexist with the bound solitons regime. As multiple pulsing in a single-pulse regime, multiple bound pairs (exhibiting all the same characteristics) are observed increasing the pumping rate, leading in some cases to harmonic mode-locking with bound states regularly spaced in the cavity. Note, however, that observation and control of bound state operation has never been reported with high-power passively mode-locked fiber lasers, and one can wonder whether such a regime can be observed and controlled with high energetic pulses, and at other wavelength than 1.55 μm. In this paper, we report on the experimental generation of bound states of two or three pulses from a high-power Yb-doped fiber laser operating at 1.05 μm in a stretched-pulse regime. Pulses are further compressed to 100 fs.

We demonstrate then theoretically that bound states generation is possible in the normal dispersion regime through the consideration of a cubic–quintic complex Ginzburg–Landau (CGL) equation. In particular, we demonstrate that bound states of two pulses exist as stable solutions of the cubic–quintic CGL equation.

2. Experimental results

A passively mode-locked Yb-doped fiber laser has been first realized [7,8]. Fig. 1 shows the experimental set-up of this laser. A diode-pumped Yb-doped double-clad fiber is used as the amplifying medium. The pumping laser diode is coupled into the fiber using the V-groove technique, which allows high pumping efficiency, while both fiber ends remain free. Passive mode-locking is achieved through nonlinear polarization rotation technique: a polarization-dependent optical isolator, and two polarization controllers before and after the doped fiber are thus introduced in the cavity. A grating pair (1200 lines/mm) is inserted in the cavity to ensure a control of the total cavity dispersion. Depending on the distance between the gratings, different regimes can be observed [8]. In our case, the distance between the gratings is 2.2 cm, which leads to a positive total cavity dispersion of +0.047 ps², and the laser operates in a stretched-pulse regime. The first grating in this scheme is used as a variable output coupler, since the part of the laser radiation reflected into the zero diffraction order (laser output) can be adjusted by rotating the phase plate just after the optical isolator. The
output signal is then analyzed with an optical analyzer (Advantest Q8384). A commercial autocorrelator, a 5-GHz-wide bandwidth sampling oscilloscope and a 8-GHz high-speed photodetector are used to study the temporal evolution of the laser output. The emission regime is strongly dependent on the orientation of the polarization controllers. Adjusting their orientation, self-starting single-pulse trains at about 18 MHz repetition rate are obtained. We measure typically 80 mW averaged output power for 1.6 W pump power. We have checked with our high-speed oscilloscope and optical autocorrelator that there is no multiple pulsing per round trip, such that the energy per pulse is 4.7 nJ. Note that the peak power of these pulses is very important (around 900 W). Duration of the pulses is measured through autocorrelation measurements. Autocorrelation traces show that the pulse duration is 5 ps at the output of the cavity in this configuration. Pulses are then compressed by another grating pair outside the cavity. After optimal compression, pulse duration is around 100 fs. Bound states generation is observed with this stretched-pulse configuration when we modify the orientation of the polarization controllers. Actually, the orientation of the phase plate before the grating pair is very important. It modifies the output coupling coefficient and thus the energy of the pulses that are reinjected in the double-clad fiber. Depending on the orientation of this element and on the pumping level, we could observe emission of bound states of two or three pulses. Let us detail one of the scenarios observed: for about 1 W pump power the laser threshold is reached and the laser delivers continuous signal. Increasing the pumping power, the mode-locking threshold is reached at 1.85 W and the laser delivers bound states of three pulses. This high-power fiber laser operates in a stretched-pulse regime, such that the wings of the pulses can slightly overlap and the pulses can be locked. Decreasing then the pumping power, this three-pulse bound state regime disappears suddenly when pumping becomes lower than 1.7 W. It is replaced by a two-pulse bound state regime. In these different bound state regimes, the laser does not present any CW component. Fig. 2 shows the autocorrelation trace of the two-pulse bound state at the output of the laser, before extra-cavity compression. Separation between the pulses is 22.4 ps, while the pulse duration is 5 ps. We measure about 75 mW output power for 1.7 W pumping power which corresponds to more than 2 nJ energy per pulse of the pair. Fig. 3 shows a typical spectrum of the bound pulses. Contrast of the interference pattern approaches unity as shown on the zoom of Fig. 3. This spectrum takes a few minutes to be recorded. The two pulses have thus to be phase-locked extremely precisely to maintain such a good contrast of the pattern, since any time jitter at the femtosecond scale would reduce significantly the fringes contrast. The precise phase relationship between the two pulses is difficult to measure in our case because of the large bandwidth emission of the laser and the large time separation between the two pulses, as previously mentioned in [6]. An exact evaluation of the period of the spectral modulation is given from the zoom of this spectrum presented in inset of Fig. 3. We obtain 0.164 nm at 1050 nm, which leads exactly to 22.4 ps time separation. The emission of bound states of two pulses was not observed when we first increased the pumping power. There is a strong hysteresis when increasing or decreasing the pumping power, leading to bistability between three-pulse and two-pulse bound state regimes. However, each of these regimes of bound states is well stable over several hours, as mentioned in the case of Er-doped bound states laser. Decreasing then the pumping power,
mode-locking disappears and CW is observed below 1.4 W pump power. Let us further point out that other scenarios can occur, depending on the orientation of the different half wave plates in the cavity. For example, with another adjustment of the phase plates, we observed similar bistability between a two-pulse bound state regime and a single-pulse operation, increasing or decreasing the pumping power. We have further compressed the pulses emitted in these different regimes through the extra-cavity grating pair. The autocorrelation trace obtained after compression is presented in Fig. 2 in the case of the two-pulse bound state regime already shown. Separation between the bound states is still 22.4 ps, while the detail of the central peak of the autocorrelation trace, presented in Fig. 4, shows that each pulse duration is now ultra-short around 100 fs assuming Gaussian pulse shape. Note that the time separation between the two bounded pulses is 220 times larger than each single-pulse duration, but pulses are still locked: the optical spectrum does not change at all after compression (bandwidth and interference pattern are the same), which confirms the precise phase-locking between the two pulses.

3. Theoretical study

From the theoretical point of view, bound states are traditionally described using a cubic–quintic Ginzburg–Landau equation or a set of two coupled nonlinear Schrödinger equations. However, it is very difficult to obtain a simple theoretical description of a stretched-pulse laser involving both anomalous and normal dispersion parts in the cavity, and including many elements like half wave plates and a polarizing optical isolator. It is known that the dynamics of the ring...
laser can be described by a master equation of CGL type [9]. Recently, we could derive a similar CGL type equation in the case of normal dispersion. It has been shown that, although the amplitude and the polarization of the pulse undergo very large variations during a roundtrip, the evolution of the electric field just after the polarizer could be described, for a large number of roundtrips, by a single scalar equation of cubic CGL type. Its coefficients depend explicitly on the orientation of the half wave plates, and good agreement with experimental results could be obtained [10]. We thus naturally carried out a first attempt to interpret bound state generation in our case of net normal dispersion regime from a sole universal CGL type equation. Mode-locking requires the existence of stable localized solutions of the equation [11–13]. In the case of normal dispersion, the cubic CGL equation does not admit such solutions, and the stability can be achieved only by taking into account higher-order nonlinear terms [14]. Therefore, we model the present situation with the cubic–quintic CGL equation

$$\begin{align*}
\frac{i}{C_0} \frac{\partial \psi}{\partial z} + \left( \frac{D}{2} - i\beta \right) \frac{\partial^2 \psi}{\partial t^2} + (1 - i\epsilon) \psi |\psi|^2 \\
+ (\nu - i\mu) |\psi|^4 = i\delta \psi,
\end{align*}$$

(1)

where $\psi$ is the normalized field amplitude. We consider here the case of normal dispersion $D = -1$. The real parameters $\delta$, $\epsilon$, $\beta$, $\mu$ and $\nu$ give account for excess of linear gain, second-order nonlinear gain, gain bandwidth, fourth-order nonlinear gain, fourth-order nonlinear index, respectively. The gain coefficients can be negative and represent absorption. It has been shown in [10] that $\epsilon$ is an effective coefficient which depends in particular on the adjustment of the optical isolator, while $\delta$ self-adjusts according to a dynamics which is related to gain saturation. The quintic terms $\mu$ and $\nu$ give account for the latter, its dependency with regard to the parameters is not explicitly known at time. Therefore, we use a dimensionless form of the CGL equation (1) with arbitrary coefficients. Notice that a wide range of values for these coefficients can be reached experimentally, by adjusting the polarization control-

On the other hand, the behaviour of the CGL equation (1) crucially depends on the sign of the group velocity dispersion $D$, taken here equal to $-1$ in accordance with the experiment. Bound states of two pulse-like solutions of this equation have been studied using either perturbation theory [2] or a energy and momentum balance approach [16, 17]. We follow the latter in the present paper. All these studies concern the anomalous dispersion regime, while we are interested here in normal dispersion, according to the experimental situation. As mentioned above, a localized stable stationary solution of Eq. (1) exists in some domain of the parameters $\delta$, $\epsilon$, $\beta$, $\mu$ and $\nu$ [14]. For sake of simplicity, we will call this solution “soliton”, according to the terminology of Akhmediev et al. It can be computed numerically. For our numerical computations, we choose a set of parameters well inside the domain where the soliton exists: $\delta = -0.1$, $\epsilon = 0.9$, $\beta = 0.8$, $\mu = -0.1$ and $\nu = -0.01$. It must be noticed that the soliton has a determined duration and amplitude, as in the case of the anomalous dispersion with gain and absorption, and contrarily to the case of propagation in conservative media. In addition to the balance between the nonlinearity and dispersion required in conservative media (which cannot be realized alone in the case of normal dispersion), a balance between gain and loss must also be realized. This results in a reduction of the number of degree of freedom of the family of localized stationary solutions from 1 for conservative systems to 0 in the nonconservative case [15].

The energy and momentum balance approach of [16, 17] reduces the problem of the interaction of two pulses to the evolution of their respective time and phase separation. Notice that this approach necessitates the existence of stable single pulses for the considered parameter values. Bound state operation will thus be predicted in domains where a single-pulse regime can exist, as observed experimentally in the domains where bistability exists between these two regimes. Further, the validity of the approach rests for a large part on the fact that no modification in the size of the soliton can occur during a two-soliton interaction. Since this property is still true in the case of normal dispersion considered here, we can use the same approach.
We recall it briefly. The pulse energy evolves with $z$ according to
\[
\frac{d}{dz} \int_{-\infty}^{+\infty} |\psi|^2 \, dt = F[\psi],
\]
where the functional $F[\psi]$ is given by
\[
F[\psi] = 2 \int_{-\infty}^{+\infty} \left( \delta |\psi|^2 + \epsilon |\psi|^4 + \mu |\psi|^6 - \beta \left| \frac{\partial \psi}{\partial t} \right|^2 \right) \, dt.
\]
Similarly, the momentum evolves according to
\[
\frac{d}{dz} \text{Im} \int_{-\infty}^{+\infty} \psi \frac{\partial \psi^*}{\partial t} \, dt = J[\psi],
\]
where the functional $J[\psi]$ is given by

Fig. 5. Zeros of the functionals $F[\psi]$ (solid line) and $J[\psi]$ (dashed line), in the plane $(\rho \cos \varphi, \rho \sin \varphi)$.

Fig. 6. Trajectories in the plane $(\rho \cos \varphi, \rho \sin \varphi)$, around the stationary points $A$ (a), $B$ (b) and $D_1$ (c).
\[ J[\psi] = 2 \text{Im} \int_{-\infty}^{+\infty} \left[ \left( \delta + \varepsilon |\psi|^2 + \mu |\psi|^4 \right) \psi \frac{\partial \psi^*}{\partial t} - \beta \frac{\partial \psi}{\partial t} \frac{\partial^2 \psi^*}{\partial t^2} \right] \mathrm{d}t. \] (5)

At this stage we can note that both functionals \( F \) and \( J \) do not depend on the dispersion. However, the soliton solution of Eq. (1) depends on the sign of the dispersion. Therefore, the stationary solutions for a bound state strongly depend on the dispersion regime. The stationary solutions satisfy \( F[\psi] = 0 \) and \( J[\psi] = 0 \). (6)

We use as a trial function
\[ \psi = \psi_0(t - \rho/2) + \psi_0(t + \rho/2)e^{iu}, \] (7)
where \( \psi_0 \) is the soliton solution, computed numerically by solving the evolution equation (1).

The two parameters \( \rho \) and \( \varphi \) represent the distance and the dephasing between the two pulses, respectively. The zeros of \( F \) and \( J \) are represented in the plane \( (X, Y) = (\rho \cos \varphi, \rho \sin \varphi) \), see Fig. 5. The figure is more complicated than in the case of anomalous dispersion \( D = +1 \) [17]. Instead of five in this latter case, we find 11 intersection points inside the domain represented on Fig. 5, and others exist outside of this domain (we did not find any stable bound state among the latter).

To determine the stability of the intersection points, we compute the velocity \( \vec{v} = \frac{d(X, Y)}{dz} \) according to the evolution equations (2 and 4), and draw the corresponding streamlines, cf. Fig. 6. This way we find that the points \( A, B \), and the symmetrical ones \( A', B' \) on Fig. 5 represent stable solutions, while \( C, C', D_1, D_2, D_3, D_4 \) and \( D_5 \) are unstable. The shape of the trajectories is very different from the case of anomalous dispersion, and the domain where attraction by the stable points \( A \) or \( B \) occurs is much smaller. (Notice, however, that the trajectories are computed here using the energy and momentum balance approximation, and not solving numerically the CGL equation (1) as in [16,17].) The stability of the bound states determined this way is checked by numerically solving the CGL equation (1). An example of computation corresponding to point \( B \) is shown in Fig. 7. We have thus shown that bound states of two solitons of the cubic–quintic CGL equation exist in the case of normal dispersion. For the parameters of our simulation, and apart from the symmetries of the problem, we have found two different bound states, instead of a single one in the case of anomalous dispersion. However, the quantitative difference between the two bound states is small, as shown in Fig. 8.

4. Conclusion

In summary, we have reported for the first time to our best knowledge the experimental
observation of bound states in high-energy passively mode-locked fiber laser operating in the normal dispersion regime at 1.05 μm (stretched-pulse operation). Stable bound states of two and three pulses have been obtained with a high-power ytterbium-doped double-clad fiber laser. These nanojoules pulses could be compressed extra-cavity to 100 fs. The laser has been modeled with a cubic–quintic CGL equation. In the region where localized solutions are stable, we have demonstrated that bound states of two pulses are stable. Experimentally, the selection of either single- and double-pulse operation is related to the pumping power. It is taken into account in the theoretical model through the total energy of the considered solution. A more complex model needs now to be elaborated to explicit the influence of the experimental parameters.

References