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A numerical model of CW active mode-locking of solid state lasers including intracavity second harmonic generation

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Résumé. — On développe un modèle numérique de la synchronisation active des modes de lasers à l'état solide en tenant compte du temps de cohérence du milieu actif. Les résultats sont comparés au traitement analytique de Kuizenga et Siegman pour les impulsions de très courte durée. Les effets associés à la génération de deuxième harmonique à l'intérieur de la cavité et à l'insertion d'un miroir non linéaire sont introduits dans le modèle. Le modèle est utilisé pour établir les conditions optimales pour la réduction de la durée des impulsions et la conversion de fréquence consécutives à ces effets.

Abstract. — We develop a numerical model of active mode-locking of solid state lasers taking into account the coherence time of the active medium. The results are compared to the analytical treatment of Kuizenga and Siegman for extremely short pulse durations. The effects of intracavity frequency doubling and of a nonlinear mirror are easily incorporated into the model. The model is applied to find the optimum conditions for additional pulse shortening and frequency conversion based on these two effects.

1. Introduction.

In the self-consistent theory of Kuizenga and Siegman [1] for the evolution of a short pulse in an actively mode-locked homogeneously broadened laser, the steady state pulse duration (FWHM) and the saturated gain are given by :

$$T = (\sqrt{\ln 2 / \delta / \pi f_m}) \times \left\{ 1 + \sqrt{1 + 1/4 \pi^2 \tau^2 f_m^2 g \zeta \delta} \right\}^{-1/2} \approx \sqrt{2 \ln 2 \tau / \pi f_m (g \zeta / \delta)^{1/4}} \quad (1)$$

$$g \zeta = \gamma + \ln \sqrt{1 + 8 \ln 2 g \zeta \tau^2 / T^2} \approx \gamma \quad (2)$$

where τ is the transversal relaxation time related to the laser bandwidth as $\tau = 1 / \pi \Delta f$, $g \zeta$ is the saturated gain for the electric field amplitude, γ is the amplitude loss, $f_m = \omega_m / 2 \pi$ is the drive RF applied to the modulator (half the repetition rate of the laser), and δ is the depth of modulation in the transmission function of the active modulator :

$$m(t) = \exp \left\{ -\delta \sin^2 [\omega_m (t - t_0)] \right\} \quad (3)$$

The analytical pulse shape obtained in [1] is Gaussian due to the approximate expressions used

to expand the transmission functions of the active medium and of the modulator. A similar parabolic expansion for the gain function has been used more recently [2] to extend the results of [1] for the off-resonant case. From (1) it is clear that in order to produce shorter pulses the cavity losses (which equal $2 g \zeta$) should be decreased and the modulation frequency f_m should be increased since Δf is fixed for the given laser and δ is at present limited to values < 5 for the available modulators.

Recently the use of higher frequency acoustooptical modulators ($f_m = 250$ MHz) was demonstrated in a CW mode-locked and diode-pumped laser [3]. The increased pulse rate resulted not only in reduced pulse duration but also in an improved performance and stability [3]. At $f_m = 50$ MHz, $2 g \zeta = 0.05$, $\delta = 1$ and $\Delta f = 120$ GHz (Nd : YAG laser) we obtain from (1) $T = 60.8$ ps. The low value of $2 g \zeta$ is representative of diode pumped solid state lasers. At higher modulation frequencies T would approach the phase relaxation time τ and consequently the parabolic approximation of the gain function [1] would be violated since the gain profile serves as the only bandwidth limiting factor in the absence of etalons. Earlier attempts to incorporate higher order dispersive terms in mode-locking

theory were done in reference [4] and they were extended in reference [5] to all orders of the gain coefficient.

In this paper we simulate numerically the mode-locking process (see Fig. 1) taking into account the exact lineshape function. The steady state solutions are compared to those obtained analytically in [1]. Further, we extend the model to include the effects of intracavity frequency doubling which can be successively applied to the newly developed compact diode pumped lasers [6]. Finally, we consider the effect of inclusion of a nonlinear mirror in the cavity [7] which incorporates the intracavity second harmonic generation (SHG) but instead of pulse broadening results in pulsewidth reduction due to its specific mode-locking action [8].

Fig. 1. — The model laser cavity. The numbers denote the position of the pulse which travels one round trip in the anticlockwise direction. The condition for CW steady state reproduction can be imposed in principle at each of these positions. By κ we denote either the SHG crystal (Sect. 4) or the nonlinear mirror (Sect. 5).

2. Round trip model.

We consider plane wave pulse propagation through a resonant medium consisting of two-level systems. The integral equation which governs the solution for the slowly varying component of the electric field in the absence of single pass saturation and longitudinal relaxation during the passage of the pulse was first derived in reference [9]. The time domain solution in the case of resonant propagation and homogeneous broadening was found in reference [10]. The more general near-resonant solution for the electric field reads :

$$\bar{E}_1(\eta) = \bar{E}_0(\eta) + 2 g \zeta J(\zeta, \eta) \quad (4)$$

with

$$J(\zeta, \eta) = \int_{-\infty}^{\eta} d\eta' \times \frac{\bar{E}_0(\eta') I_1[2 \sqrt{g \zeta (\eta - \eta')}] \exp[-(\eta - \eta')/\zeta_{21}]}{2 \sqrt{g \zeta (\eta - \eta')}} \quad (5)$$

where I_1 is the modified Bessel function. The subscripts « 0 » and « 1 » of $\bar{E}$ denote the position of the pulse before and after the gain medium respectively according to figure 1 if ζ is the length of the active medium. The local time is $\eta = (t - z/v)/\tau$, $\zeta_{21} = [1 + i\tau(\omega_L - \omega_{21})]^{-1}$ is the complex lineshape factor where ω_L is the carrier frequency and ω_{21} is the frequency spacing of the two-level systems modeling the active ions. The amplification factor g has a dimension of inverse length and v is the group velocity of the pulse as determined by the dispersion of the host material.

It is clear that the use of equation (4) would lead to two basic deviations from the results obtained in [1]. First, exact consideration of the Lorentzian natural bandshape would not support a Gaussian shape of the steady state pulses, and second, complex lineshape functions result in asymmetric time domain solution even in the on-resonance case $\omega_L = \omega_{21}$.

For typical values of δ and $\Omega = \tau \omega_m$ the usual parabolic approximation may be made for the modulation function giving :

$$\begin{aligned} \bar{E}_2(\eta) &\approx \bar{E}_1(\eta) \exp[-\delta \Omega^2(\eta - \eta_0)^2] \approx \\ &\approx \bar{E}_1(\eta) [1 - \delta \Omega^2(\eta - \eta_0)^2] \end{aligned} \quad (5)$$

The reflection at the output mirror (see Fig. 1) which is assumed to contain all the cavity losses modifies the field as follows :

$$\bar{E}_3(\eta) = \bar{E}_2(\eta) \exp(-\gamma) \approx (1 - \gamma) \bar{E}_2(\eta) \quad (6)$$

The last approximation is well justified for CW lasers where the total losses are considerably reduced. For small losses of the fundamental energy in the nonlinear processes both the SHG and the nonlinear mirror introduce the following correction to the field :

$$\bar{E}_4(\eta) = [1 + \beta (|\bar{E}_3(\eta)|^2)] \bar{E}_3(\eta) \quad (7)$$

Combining (4), (5), (6) and (7) results in a self-consistency condition :

$$\begin{aligned} \bar{E}(\eta - h/\tau) &= \{ \bar{E}(\eta) + 2 g \zeta J(\zeta, \eta) - \\ &- \delta \Omega^2(\eta - \eta_0)^2 \bar{E}(\eta) \\ &- \gamma \bar{E}(\eta) + \beta (|\bar{E}(\eta)|^2) \bar{E}(\eta) \} \exp(i\chi) \end{aligned} \quad (8)$$

where we denote by h the time delay caused by the active medium and partly reduced by the modulator which should be compensated under optimum conditions by detuning of the resonator. The additional phase shift χ is insignificant and accounts for a possible deviation of the propagation constant from the assumed linear value $k_L = \omega_L n/c$.

The model can be completed now if a specific saturation scheme is assumed in order to limit the

energy. The situation is different for homogeneous and inhomogeneous broadening [1] and in general spectral decomposition of the field is needed. (Note that since single pass saturation has been neglected in (4) it remains valid in the case of inhomogeneous broadening as well provided the inhomogeneous lineshape is of Lorentzian type. The latter is a rather insignificant assumption and obviously inhomogeneously broadened systems can be treated by redefining τ as the inverse of the total linewidth.) We assume a rather simple expression for g following [11]:

$$g = g_0 / (1 + 2 f_m T_{21} \varepsilon). \quad (9)$$

In (9) g_0 is the small signal gain, T_{21} is the upper state lifetime, and $\varepsilon = \int_{-\infty}^{\infty} |\bar{E}(\eta)|^2 d\eta$ is the normalized intracavity energy density (as a normalizing constant the saturation energy density for a two-level system has been used).

Steady state solutions are generated by iterating equation (8) with an arbitrary initial distribution of $\bar{E}(\eta)$. The saturated gain (9) is corrected after each round trip. Transient effects are not considered and the stationary pulse is independent of the input signal. It is characterised by high stability and deviation of $< 0.01\%$ for integral parameters like energy, duration, asymmetry etc., from round trip to round trip. 256 discrete samples have been used to cover the interval $|\eta| \leq 30$. The calculation is performed in the time domain and thus the use of apodizing filters [12] has been avoided. The parameter h is calculated from the delay experienced by the pulse and averaged over 100 round trips after the steady state has been achieved. The necessary cavity shortening in the optimum regime is equal to $hc/2$ for a linear cavity. The number of round trips necessary to reach the steady state varies between several hundreds and several thousands and is smaller for shorter steady state pulse durations.

3. Comparison with analytical results.

We use further the following laser parameters: $\gamma = 0.025$, $T_{21} = 230 \mu s$, $\Delta f = 120 \text{ GHz}$ (i.e. $\tau = 2.65 \text{ ps}$) in order to simulate the mode-locked operation of a Nd : YAG laser. The zero point of the modulator can be chosen at $\eta_0 = 0$ and in order to compare with [1] we assume $\omega_L = \omega_{21}$ and study only real solutions. The parameters which are analysed in dependence on the modulator frequency are the pulse duration T , the pulse delay h and the pulse asymmetry α defined as a ratio of the widths of the leading and the trailing edges measured at half maximum intensity. The value of $g_0 \zeta$ has no influence on these parameters and should only exceed γ in order to initiate the simulation. We note that

according to the Kuizenga-Siegman theory [1] h equals $g\zeta\tau$.

Figure 2 shows that the Kuizenga-Siegman model can adequately predict the steady state pulse duration even for spectral bandwidths comparable to the transition bandwidth. Deviations of more than 10 % can be observed for $T/\tau < 5$. The pulse duration predicted by the exact consideration of the lineshape function is always shorter than the one predicted by the Kuizenga-Siegman model. There is a physical reason for that behaviour. The Kuizenga-Siegman model assumes a Gaussian gain function which is a good approximation around line centre, but is totally wrong on the wings of the gain curve. In fact, the transmittance of the gain medium tends to unity outside the gain curve, while the Gaussian approximation predicts a null transmittance; as a result, the Kuizenga-Siegman model introduces excessive bandwidth limiting, which, by simple time-

Fig. 2. — Steady state pulse parameters *versus* modulator frequency for $\beta = 0$. Relative deviation of the Kuizenga-Siegman FWHM T_{KS} from T : $\Delta T/T = (T_{KS} - T)/T$ (a), pulse asymmetry α (b) and pulse delay h (c). T_{KS} is calculated by successive approximations of the left parts of equation (1-2), $\delta = 1$ (solid lines), $\delta = 5$ (dashed lines).

of longer pulses. The asymmetry of the steady state pulse is more pronounced for shorter pulse durations and indicates steeper leading edge of the pulse. It has the same origin: a proper account of the complex gain profile. Physically it is linked to the causal nature of the gain function or the on-resonance dispersion of the gain [13, 14]. The steady state gain $g\zeta$ is in excellent agreement with the iterative calculation based on (1)-(2) and it is not plotted in figure 2. The pulse delay in the active medium decreases with the pulse duration and exhibits opposite behaviour to that analytically derived in [1]. Similar decrease of h with the pulse duration has been obtained in [4] when reducing the cavity losses γ or increasing the modulation depth δ .

4. Addition of an intracavity frequency doubler.

The insertion of a SHG crystal inside the cavity utilizes the higher peak power available for effective frequency conversion. The previous analytical treatment of this process [15] provided an analytical extension of the Kuizenga-Siegman theory, the SHG coefficient being represented, however, by only a mean value. In our case assuming small conversion efficiencies we have $\beta = -m\kappa I(\eta)$, where 2κ is the normalized intensity conversion efficiency related to the intensity used for normalization of the field [15] and m is a telescoping factor accounting for a possible focusing in the SHG crystal. We define the normalized intensity by $I = |\bar{E}|^2$ i.e. the normalizing intensity is $\hbar\omega_L/\sigma\tau$ where $\sigma = \mu_0 \omega_L^2 \mu^2 \tau / \hbar k_L$ is the maximum interaction cross section (μ : dipole moment, k_L : wave number). From the value $\sigma = 8.8 \times 10^{-19} \text{ cm}^2$ (Nd:YAG) it can be calculated that the intensity is normalized to the value 38.2 GW/cm^2 and the dimensionless energy density — to the value 100 mJ/cm^2 .

Figures 3a, b show the effect of inclusion of an intracavity SHG process in the laser. We have chosen the same parameters as in figure 2 for the case $\delta = 1$ and $f_m = 500 \text{ MHz}$. The small signal gain is $g_0 \zeta = 1$. The maximum SH intensity is $I_m^{\text{SH}} = 2m\kappa I_m^2$ and $\varepsilon^{\text{SH}} = 2m\kappa \int_{-\infty}^{\infty} I^2(\eta) d\eta$. The

SH generation introduces an intensity dependent loss which broadens the pulse and reduces its intensity. For practical purposes it is important to produce short and energetic pulses at the second harmonic. That is why the best parameter for an optimization is the maximum intensity of the second harmonic I_m^{SH} . As can be seen from figure 3 I_m^{SH} grows rapidly with $m\kappa$ then saturates and for larger values of $m\kappa$ slowly goes down. Similar behaviour of I_m^{SH} was derived in [16] for the mean intensity by a simple balance of the gain and loss mechanisms. As can be concluded from the chosen example the optimum value for $m\kappa$ lies in the region of 2 500 since I_m^{SH} has nearly reached its maximum value and T is still short enough.

Since Siegman and Heritier [12] used essentially the same equation (8) written partly in the time and partly in the frequency domain we could check our results using the following modified parameters: $g_0 \zeta = 0.03$, $\gamma = 0.007$, $\delta = 0.08$, $f_m = 250 \text{ MHz}$ and $m\kappa = 3.47 \times 10^6$ (see Sect. III of Ref. [12]). In addition the factor 2 in (8) was increased twice to account for the saturation formula of a linear resonator (Eq. (5) in [12]) and the time interval considered only for this comparison was $|\eta| \leq 150$. With these parameters we obtain somewhat longer pulse duration (426 ps compared to 405 ps in [12]). The saturated gain and the SHG efficiency are somewhat higher — 0.0254 and 2.9 % compared to

Fig. 3. — Steady state pulse parameters versus normalized SHG parameter $m\kappa$. a) Pulse duration T (solid line), delay h (dashed line) and steady state gain $g\zeta$ (dotted line). b) Maximum SH intensity I_m^{SH} (solid line), effective energy conversion $\varepsilon^{\text{SH}}/\varepsilon$ (dashed line), intracavity energy ε (dash-dotted line) and maximum intracavity intensity I_m (dotted line).

0.023 and 2 % in [12] respectively. We attribute these discrepancies to the windowing functions used in [12] to suppress the aliasing effect accompanying the discrete Fourier transforms.

5. Nonlinear mirror.

As shown in [17] in the case of the nonlinear mirror $\beta = \bar{\gamma} \tanh^2[K \sqrt{mI}(\eta)]$ where $\bar{\gamma}$ is the linear amplitude loss associated with the nonlinear mirror which is assumed to be small compared to unity and γ includes $\bar{\gamma}$ i.e. $\gamma > \bar{\gamma}$. K is related to the conversion efficiency κ corresponding to the intensity used for normalization: $\kappa = \tanh^2 K$. The model is valid for extremely short pulses since the response of the nonlinear mirror is very fast [7]. An important limitation for the achievable pulse duration becomes the group-velocity dispersion in the SHG crystal. This effect is not considered here since its time scale is less than the transversal relaxation time τ . Since

the nonlinear mirror introduces nonlinear losses we define a mean loss coefficient

$$\langle \gamma \rangle = \gamma - \bar{\gamma} \int_{-\infty}^{\infty} \beta(\eta) I(\eta) d\eta / \int_{-\infty}^{\infty} I(\eta) d\eta.$$

As can be seen from figures 4a, b the increase of the small signal gain $g_0 \zeta$ results in constant increase of the intracavity energy and intensity whereas the pulse duration is considerably reduced due to the enhanced mode-locking action of the nonlinear mirror. For small values of $g_0 \zeta$ linear dependence of the maximum conversion efficiency κ_m on I_m could be observed as discussed in [17]. Higher values of I_m lead to a saturation behaviour of κ_m (Fig. 4a). The leading edge of the pulse is always steeper than the trailing edge and this asymmetry is much more pronounced than in the cases discussed in sections 3

and 4. The pulse delay h is always smaller for shorter pulse lengths. The steady state saturated gain $g\zeta$ exceeds the mean loss and tends to follow its dependence.

Figure 5 illustrates that optimization of the pulse parameters is possible by variation of the nonlinearity. The pulse duration reaches a minimum and the pulse intensity — a maximum at a given value of K whereas the energy increases with K . Further increase of the nonlinearity or the focusing factor m results in pulse broadening. This behaviour can be explained by the high conversion efficiency. If the conversion efficiency is close to unity for the most part of the pulse the shortening action of the nonlinear mirror is cancelled since the second harmonic repeats the fundamental [8]. For the parameters used in figure 5 the conversion efficiency in the maximum of the pulse κ_m is close to unity for $K \geq 50$.

The pulse delay follows the dependence of T and is not shown in figure 5. The dependence of

Fig. 4. — Steady state pulse parameters *versus* small signal gain in the case of a nonlinear mirror. $f_m = 500$ MHz, $\delta = 1$, $\gamma = 0.05$, $\bar{\gamma} = 0.025$ and $K = 50$. a) Intracavity pulse energy density ϵ (solid line), maximum intensity I_m (dash-dotted line), pulse duration T (dashed line) and maximum conversion efficiency κ_m (dotted line). b) Mean loss $\langle \gamma \rangle$ (solid line), saturated gain $g\zeta$ (dashed line), pulse delay h (dash-dotted line) and asymmetry α (dotted line).

Fig. 5. — Steady state pulse parameters *versus* nonlinearity K in the case of a nonlinear mirror. $f_m = 500$ MHz, $\delta = 1$, $\gamma = 0.05$, $\bar{\gamma} = 0.025$, $g_0 \zeta = 1$. a) Intracavity pulse energy density ϵ (solid line), maximum intensity I_m (dash-dotted line) and pulse duration T (dashed line). b) Mean loss $\langle \gamma \rangle$ (solid line), saturated gain $g\zeta$ (dashed line) and asymmetry α (dotted line).

the asymmetry α shows that the leading edge of the pulse broadens when the conversion efficiency begins to saturate. The saturated gain $g\zeta$ follows the dependence of the mean loss $\langle\gamma\rangle$ (Fig. 5b). The two curves are closer in the limits $K \rightarrow 0$ and $K \rightarrow \infty$. This could be expected since both of them correspond to pure active mode-locking with total losses equal to 0.05 and 0.025 respectively.

6. Conclusion.

The paper presents a numerical study of actively mode-locked CW operation of solid state lasers with intracavity second harmonic generation. The laser parameters used are most appropriate to describe the operation of diode-laser pumped Nd: YAG arrangements.

In the regime of pure active mode-locking the deviations from the analytical results of Kuizenga and Siegman [1] can be attributed to the exact consideration of the lineshape function. The most important of them is the reduced pulse duration as it approaches the transversal relaxation time.

The analysis of the intracavity second harmonic generation shows that there exists a maximum value for the intensity of the second harmonic when the nonlinearity parameter is varied and this dependence can be used for optimization of the operation regime.

In the analysis of the nonlinear mirror we included

not only nonlinear losses as in [17] but also additional linear losses. The latter modifies in fact only quantitatively the results of [17]. Up to now the considered passive mode-locking technique has been applied only to flashlamp pumped pulsed systems [19-21] and only in a recent experiment an active mode-locker has been added [22]. All of these systems are characterized by losses comparable to unity and consequently all resonator elements can substantially affect the pulse parameters during a single pass. That is why direct comparison with the above experiments is not possible. Estimations of the peak intracavity intensity in CW actively mode-locked diode pumped lasers [17], however, indicate that it can be high enough in order to initiate the nonlinear processes.

Since linear and nonlinear losses are included the model is adequate to describe coupled cavity configurations as that realized in [21]. Our results are, however, valid for only limited range of the parameters used and direct comparison with the pulsed system described in [21] is not possible since both γ and $\bar{\gamma}$ are comparable to unity there. Nevertheless our estimations can be useful if CW generation in this so called additive mode-locking regime could be achieved.

Finally we demonstrated that there exists an optimum value of the nonlinearity and/or the focusing parameter giving minimum achievable pulse duration.

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