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Self-mode locking in a cw solid-state laser by means of a Kerr nonlinear polarization modulator

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An extremely promising method of generating stable ultrashort pulses with durations in the femtosecond range for cw solid-state lasers makes use of an ultrafast Kerr nonlinearity in a single-resonator layout. There are several effects of this kind that can be used in a passive modulator. One possible choice is the Kerr nonlinear optical effect, in which an elliptically polarized field propagates in a nonlinear medium with an intensity-dependent angle of self-precession. Such a medium behaves as a discriminator when placed between a polarizer and an analyzer crossed at a certain angle, and can be used either to increase the efficiency of mode locking or, independently, as a wide-band passive modulator for solid-state lasers. Such modulators are practically instantaneous, have a wide transmission band, are easy to tune by simply changing the angle of orientation of the modulator elements, and are capable of being either bleached or darkened. The latter feature can be used to stabilize the oscillations. In contrast to certain mode-locking methods (widely used with picosecond cw lasers) that require an additional resonator, lasers with Kerr nonlinear polarization modulators do not require long-term, interferometrically precise control of the resonator length.

However, up to now there has been no detailed study of the oscillation dynamics of a cw laser with a Kerr nonlinear polarization modulator. The issue of whether it is possible to increase the efficiency of self-mode locking in such systems and thus generate stable cw trains of ultrashort pulses has not received the attention it deserves. In light of these facts, we have investigated ways to increase the efficiency of mode locking of cw solid-state lasers by controlling the parameters of a nonlinear Kerr polarization shutter.

Our numerical simulation in Ref. 14 included an investigation of changes in the temporal profile of an oscillation field with intensity \( I \) after many passes through an active medium (with absorption cross section \( \sigma_{14} \), emission cross section \( \sigma_{32} \), maximum amplification \( A_m \), and excited-state relaxation time \( T_{31} \); Ref. 15), pumped by a flux of photons with constant density \( I_p \), and through a Fabry-Perot etalon with a characteristic time \( T_f \), which determines the group delay of the field in the etalon. The model includes linear losses (with logarithmic loss coefficient \( \Gamma \)) and the interaction with a nonlinear Kerr polarization modulator.

The Kerr polarization shutter is as described in Ref. 7. We assume that linearly polarized light from the active medium is incident on a quarter-wave plate, and that its polarization vector makes an angle \( \alpha \) with the optic axis of the plate. The elliptically polarized light that leaves the quarter-wave plate self-precesses in the nonlinear optical medium through an angle

\[
\varphi = p I \sin 2\alpha,
\]

where \( p \) is the Kerr polarization parameter, which is proportional to the corresponding component of the fourth-order susceptibility tensor and the length of the nonlinear medium. For the value of \( \sigma_{32} \) corresponding to Ti:AlO$_3$ the normalized coefficient is \( p/\sigma_{32} T_{\text{cw}} = 0.008 \) for a Cs$_2$ sample 10 cm long and for a resonator period \( T_{\text{cw}} = 10 \) ns. At the exit from the shutter, the light is converted back to linear polarization by the analyzer, which is oriented at an angle \( \beta \) with respect to the optic axis of the quarter-wave plate. The transmission of this system is

\[
T = (\sin^2 \alpha \sin^2 \psi + \cos^2 \alpha \cos^2 \psi)^{1/2},
\]

where \( \psi = \beta - \varphi \).

In Fig. 1 we show the dependence of the transmission \( T \) on the angles \( \alpha \) and \( \psi \) in the intervals \( 0 < \alpha < \pi/4 \) and \( 0 < \psi < \pi \). We see that, in the case \( p > 0 (p < 0) \), decreasing (increasing) the angle \( \psi \) by increasing \( I \) causes an increase in the transmission on the angular intervals \( 0 < \alpha < \pi/4, 0 < \psi < \pi/2 \) and \( \pi/4 < \alpha < 3\pi/4, \pi/2 < \psi < \pi (\pi/4 < \alpha < \pi/2, 0 < \psi < \pi/2 \) and \( 3\pi/4 < \alpha < 2\pi, \pi/2 < \psi < \pi) \). The modulator thus operates as a discriminator in the bleaching regime. For \( p < 0 (p > 0) \) the transmission decreases on these intervals, so the modulator acts as an intensity limiter in the darkening regime. The intervals with \( \alpha = \pi/4 \), which correspond to circular polarization of the radiation leaving the quarter-wave plate, are insensitive to the field intensity. Thus, by regulating the orientations of the analyzer and the modulator, i.e., the angles \( \alpha \) and \( \beta \), one can match the
FIG. 1. Transmission of a Kerr nonlinear polarization modulator.

FIG. 2. Evolution of the profile of the logarithm of the oscillation intensity $A = \sigma_{14} T_{cav} f_p$, shown at intervals of 2 passes during the linear stage of oscillation, up to the 20th pass, and then at intervals of 20 passes during the nonlinear stage, up to the 740th pass through the resonator. The time window was equal to the period of the resonator, $T_{cav} = 1$ ns. Here $\sigma_{14} T_{cav} f_p = 0.05$, $p/\sigma_{14} T_{cav} = 0.05$, $\beta = 0.4$, $\alpha = \pi/12$, $\Delta t = 1$ ps, $t_{coh} = 1$ ps, $t_j = 1$ ps, $A_n = 1.5$, $T_{coh} = 3.5$ ps, $S_0 = 10^{-7}$, and $T = 0.2$.

modulator can be put in a desired operating regime, with a given initial transmission and a given degree of bleaching. The latter is determined by the difference between the maximum and initial transmission levels of the modulator; regions of both positive and negative passive feedback may be specified as well. This is a considerable advantage over traditional electrooptic shutters based on a bleachable absorber.

We are primarily interested in investigating the oscillation dynamics of cw solid-state lasers with Kerr nonlinear polarization modulators. The main problem here is to choose an approximation for the oscillation field in one period of the resonator. We used two models. In the first, the oscillation field on the $k$th pass was specified on a uniform grid of local times $t_j$ spaced at $\Delta t$ within a time window, which completely filled a period of the resonator. The initial field was noise $S(t_j)$ with a Lorentz spectral profile determined by the coherence time $t_{coh}$ (Ref. 17):

$$S(t_j) = S(t_{j-1}) \exp(-\Delta t/t_{coh}) + S_0 \exp[i\xi(t_j)],$$

where $\xi(t_j)$ is a random phase and $S_0$ is the noise power level. In this formulation, the noise is then added to the field on each pass through the resonator, in order to identify the steady-state regime corresponding to generation of a train of ultrashort pulses. This model has several favorable features: it exploits the fluctuational character of the radiation from a solid-state laser to the maximum possible degree, and it can take into account fully the interaction of the field with the laser elements in the course of the self-mode locking. However, attempts to simulate such an oscillation field for real resonator periods led to a catastrophic increase in the volume of calculations. Therefore, we chose $T_{cav} = 1$ ns for this model.

In order to extend our analysis of the dynamics of mode locking to real $T_{cav}$, we used a modified model that uses (3) to describe the oscillation field in detail only for a time window considerably smaller than $T_{cav}$, and which treats the field outside the window as a steady-state background with intensity $N$. This background, which can be described by Eq. (3) if we set $t_{coh} \to \infty$ and $\xi(t_j) = 0$, interacts with the active medium and the Kerr modulator. This model, although less precise than the first, nonetheless allows us to take into account the noisy character of the oscillations and to investigate the mode locking over a wide interval of ratios of $t_{coh}$, $t_j / \Delta t$, and $T_{cav}$.

Our calculations revealed that the system generates a stable train of high-intensity ultrashort pulses only for certain choices of the parameters $\alpha$, $\beta$, $p/\sigma_{14} T_{cav}$, and $B = \sigma_{14} T_{cav} f_p$. This was true both for $T_{cav} = 1$ ns and for $T_{cav} = 10$ ns. In Fig. 2 the laser pulse train is shown on a logarithmic scale for $T_{cav} = 1$ ns. In the initial stage the oscillations are characterized by a rapid exponential growth of the intensity and a mode selection, due to the finite width of the transmission band of the frequency filter. This effect generates a large-scale heterodyning of the intensity from the initial noise. In the next stage the increase in intensity becomes smaller as the population inversion disappears, although the mode selection continues. This stage is characterized by the beginning of discrimination among the fluctuation-induced spikes within the time window and by a "suppression" of the steady-state noise outside it (for the case of the second model), due to bleaching of the modulator. This results in an ultrashort laser pulse that removes the population inversion inside it; outside the largest pulse the field is suppressed. The final stage corresponds to steady-state generation of ultra-short pulses, for which the effects of bleaching of the shutter, pumping, mode selection, and removal of the population inversion are in balance. Outside the laser pulse we observe only the low-intensity noise $S_0$.

Figure 2, which shows the linear, nonlinear, and steady-state stages of the oscillation, reveals certain features of passive mode locking that are characteristic not only of systems with Kerr polarization modulators but also of systems with additional resonators and bleachable absorbers. However, our calculations show that a system with a Kerr polarization modulator differs from systems of the latter two types in that destabilizing growth of the
ultrashort pulse intensity is limited by an automatic tendency for the modulator to go into the darkening regime when the oscillation field reaches a certain intensity.

In our study of how the nature of the mode locking depends on the system parameters, we found that effective mode locking is possible only in bounded ranges of the parameters of the laser system. This is clear from Fig. 3, where we show the region of self-mode locking in the plane of the parameters $B = \sigma_{14} T_{\text{cav}} I_p$ and $\beta$. For this case, we used a time window that coincided with $T_{\text{cav}}$, for $T_{\text{cav}} = 1$ ns and was 200 ps for $T_{\text{cav}} = 10$ ns. We adhered to the following criteria for self-mode locking: no less than 99% of the field energy over a period of the resonator must be in the largest pulse, and this concentration of energy must prevail over no less than 10% of the entire oscillation time, which we choose to be 2000 passes through the resonator.

Outside the self-mode locking regions we observe various oscillation regimes. There is no mode locking below the lower boundary; i.e., no oscillation pulse forms against the noise background. Above the upper boundary the locking becomes inefficient; symptomatic of this is the appearance of several pulses within the time window and an increased fraction of energy in the noise.

The calculations show that going from the exact model with a time window equal to $T_{\text{cav}}$ to the modified model results in a certain broadening in the region of self-mode locking, due to an expansion of its upper boundary as the temporal window decreases in comparison to $T_{\text{cav}}$. As the time interval occupied by the noise increases, it becomes more difficult to achieve mode locking, and the matching conditions imposed on the factors that influence the formation of the ultrashort pulses for efficient self-mode locking become more severe. When the noise outside the time window is taken into account within the approximate model, its contribution to the process of self-mode locking consists only of removing the population inversion of the background intensity. However, we may infer from a comparison of the regions shown in Figs. 3a and 3b that the qualitative behavior of these regions as the laser parameters change is insensitive to which noise model we use, and it also is preserved we go to realistic resonator periods.

It is clear from Fig. 3 that increasing $\beta$ causes an increase in the threshold for self-mode locking, which is associated with a decrease in the initial transmission of the electrooptic shutter (see Fig. 1). However, in this case the upper boundary of the self-mode locking region expands, because of the increase in the degree of bleaching of the modulator and a concomitant increase in its discrimination properties. Decreasing the normalized nonlinearity parameter leads to an increase in the threshold for self-mode locking and a compression of the area of the corresponding region, so the adjustments of the laser system needed for generation of steady-state ultrashort pulse trains become difficult to achieve (see regions 1 and 3 in Fig. 3).

It is possible to expand the upper boundary of the region of self-mode locking, and consequently to broaden this region, by decreasing $\alpha$. To do this we must increase the degree of bleaching of the shutter and enhance the discrimination against spikes in the oscillator field (see regions 1 and 2 in Fig. 3). However, this change causes the angle of self-precession induced by the nonlinear element to decrease, and this decrease is equivalent to a decrease in the effective nonlinearity of the latter [see (1)]. This, in turn, increases the threshold for self-mode locking. Further decreases in $\alpha$ lead to appreciable narrowing of the locking region; hence, there exists an optimal $\alpha$ in the sense that efficient mode locking is possible over the widest range of the parameters $\beta$ and $\sigma_{14} T_{\text{cav}} I_p$ when $\alpha$ has this value. For the parameters we have chosen, this optimal $\alpha$ lies between $\pi/6$ and $\pi/12$.

As the parameters of the system vary within the region of self-mode locking, the oscillation dynamics exhibit a number of peculiarities. From Fig. 4 (solid curves) it is clear that as we move upward from the lower boundary of region 1 in Fig. 3b the duration of the ultrashort laser pulses decreases markedly, reaching a minimum within the region. With further increases in pumping, $I_0$ begins to increase due to the decreased efficiency of mode locking. When $\beta$ is increased, so the efficiency of mode locking increases as a result of an enhancement of the discrimina-

FIG. 3. Regions of efficient self-mode locking for a laser with a Kerr polarization modulator. a: For a time window $T_{\text{cav}} = 1$ ns, for $p/\sigma_{12} T_{\text{cav}} = 0.05$ (1,2), $\alpha = \pi/6$ (1,3), $\pi/12$ (2), $\Delta t = 1$ ns. b: For a time window equal to 200 ps, for $T_{\text{cav}} = 10$ ns, $p/\sigma_{12} T_{\text{cav}} = 0.005$ (1,2), 0.0025 (3), $\alpha = \pi/6$ (1,3), $\pi/12$ (2), $\Delta t = 4$ ns. The parameter values are otherwise the same as in Fig. 2.
tion properties of the modulator, the minimum duration $t_0$ decreases by roughly a factor of 3. However, because the self-mode locking threshold increases, the position of this minimum duration shifts to slightly higher $I_p$.

Figure 4 shows the dependence on normalized pump intensity of the time to self-mode locking, $T_0$ (dashed curves), where $T_0$ is defined as the time interval from the onset of the oscillation to the time when an ultrashort pulse is selected that contains 99% of the field energy within a period of the resonator. It is clear that $T_0$ decreases abruptly along the lower boundary of the region near values $I_p$ corresponding to the minima of $t_0$; thereafter, it may increase slightly. The fact that $T_0$ leaves the mildly sloping portion at the lower boundary of the region when its value is less than 2000 passes through the resonator indicates that the criterion we have chosen for efficient self-mode locking (see above) is not arbitrary. In other words, the region of efficient self-mode locking is independent of the number of passes chosen to identify the onset of self-mode locking, as long as this number is not smaller than the value of $T_0$ for the asymptotic portion near the lower boundary of the region. We note that increasing the threshold for self-mode locking by increasing $\beta$ causes the duration of the transient period prior to self-mode locking to increase as well; however, the amount of increase is insignificant (see curves 1–3 in Fig. 4).

These features of the dynamics of mode locking indicate that there are zones within the region of self-mode locking in which the laser parameters of the system are optimally matched, and for which maximally stable ultrashort pulse trains with minimum duration appear rapidly.

A feature of the Kerr nonlinear polarization modulator which turns out to have an important effect on the formation of ultrashort pulses is the tendency of these devices to cause phase self-modulation of the oscillation field. In Fig. 5 (region 1) we show the region of efficient self-mode locking corresponding to Fig. 3b (region 1) in the presence of phase self-modulation with the normalized coefficient $f/\sigma_{ij} T_{\text{cav}} = 0.001$, where the phase self-modulation coefficient $f$ is proportional to the nonlinear index of refraction of the nonlinear medium and to its length. The phase self-modulation was taken into account in the same way as in Ref. 14. It is clear that in the presence of phase self-modulation mode locking is suppressed everywhere except within the region of maximally efficient locking, which corresponds to parameters lying near the minimum $t_0$ (Fig. 4).

However, the negative effect of phase self-modulation on the mode locking can be countered both by adjusting the system for the zone of maximally efficient self-mode locking and by broadening the zone itself, e.g., by narrowing the width of the transmission band for the frequency-selective element (region 2 in Fig. 5). In Fig. 5 (region 3) we also show the region of self-locking when a negative dispersion of the group velocity is introduced, which is included in the spectral representation in the form

$$I_{k+1} = I_k \exp(id\omega^2),$$  \hspace{1cm} (4)

here $I_k(I_{k+1})$ are the spectral components of the field at the $k$-th ($k+1$)-th pass; $\omega$ is the frequency; and $D$ is the dispersion parameter. The dispersion compensates for

FIG. 4. Duration of the ultrashort laser pulse $t_0$ (solid curves) and time for development of self-mode locking $T_0$ (dashed curves) versus the pump intensity $B$ for $T_{\text{cav}} = 10$ ns, $p/\sigma_{ij} T_{\text{cav}} = 0.005$, $\alpha = \pi/6$ and $\beta = 0.2\pi$ (1), $0.3\pi$ (2) and $0.4\pi$ (3). The parameter values are otherwise the same as in Fig. 3b.

FIG. 5. Region of efficient self-mode locking of a Kerr polarization modulator in the presence of phase self-modulation for $T_{\text{cav}} = 10$ ns, $p/\sigma_{ij} T_{\text{cav}} = 0.005$, $\alpha = \pi/6$, $\gamma = 1 (1, 3)$ and $2 \pi$ (2), $D/\sigma_{ij} T_{\text{cav}} = 0 (1, 2)$ and $2.5 \times 10^{-5}$ (3). The parameter values are otherwise the same as in Fig. 3b.
phase self-modulation; its effect is manifested in a broadening of the corresponding region of self-mode locking. An advantageous feature of this technique is the fact that the size of the region is quite insensitive to the ratio $D/T^2_{cav}$, whose optimum values lie in the range $5 \cdot 10^{-6}$ to $3 \cdot 10^{-5}$. These values are overestimates compared to the true ones, for the following reason. By assuming that the pump is switched on instantaneously, we overestimate the growth rate for the gain in our model at the initial stage of oscillation. This causes the nonlinear phase shift due to phase self-modulation to be large, requiring a larger dispersion to compensate for it.

In conclusion we note that our numerical simulation of the self-mode locking of cw solid-state lasers with a Kerr nonlinear polarization modulator is based on two different approximations for the fluctuating oscillator field. Both imply the existence of closed regions in the system parameter space within which high-stability ultrashort pulse trains are generated with durations close to the limiting value set by the transmission band for of the frequency filter. Analysis of the general features of self-mode locking in such systems points out ways to increase the efficiency of mode locking in the presence of phase self-modulation, by optimizing the laser parameters, narrowing the transmission band of the frequency filter, and introducing a compensating negative dispersion of the group velocity.