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Passive switches for femtosecond lasers based on a nonlinear Fabry–Perot interferometer with a semiconductor absorber

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Abstract. An analytic model for the description of the generation of ultrashort pulses in solid-state lasers having passive switches based on semiconductor absorbers, considered as nonlinear Fabry–Perot interferometers, is proposed. The possibilities of controlling the parameters of ultrashort pulses by varying the switch parameters are demonstrated. The stability of the solutions obtained is investigated.

One of the recent major advances in the field of the methods and technology for the generation of ultrashort pulses has been the fabrication of multilayer semiconductor mirrors [saturable semiconductor mirrors (SSM) and saturable Bragg reflectors (SBR)], operating as passive switches in the time range from nanoseconds to several femtoseconds. Ultrashort pulses were generated with their aid in many solid-state laser media: Ti: sapphire [1], Cr: forsterite [2], Nd: YLF [3], neodymium glass [4], Cr: LiSAF [5], Cr: YAG [6] by using saturable semiconductor mirrors, as well as Ti: sapphire, Cr: LiSAF [7], and Cr: forsterite [8] by using saturable Bragg reflectors.

Saturable semiconductor mirrors and saturable Bragg reflectors consist of a film of a bulk or quantum-well semiconductor, placed in a Fabry–Perot cavity formed by mirrors, one of which has a 100% coefficient of reflection, of the phase shift between the mirrors. When account is taken, after an interferometer period. Suppose that the interferometer is formed by mirrors, one of which has a 100% coefficient of reflection, of the phase shift between the mirrors. When account is taken, after an interferometer period.

A theory of the soliton mechanism of the generation of ultrashort pulses in solid-state lasers has been proposed [9] to explain this phenomenon. It takes into account the operation of a semiconductor switch as a saturable absorber with a small (not more than 5%) depth of modulation, working in a strong saturation regime and introducing an additional shift of the carrier frequency of a pulse as a result of self-modulation broadening of the absorption line. The effects associated with the interaction of the field with a semiconductor switch acting as a nonlinear Fabry–Perot interferometer, and in particular the effects associated with a possible detuning of the pulse frequency from the resonance frequency of the interferometer, were not taken into account.

In the present study we shall use an approach in which a semiconductor switch is regarded as a nonlinear Fabry–Perot interferometer (NFPI). It will be shown that an interferometer exhibits both nonlinear phase and amplitude responses capable of ensuring generation of an ultrashort pulse and the stability of a pulse against energy fluctuations. Our method for the investigation of the ultrashort pulse generation regime is based on the self-consistent field theory described in the propagation of a signal $a(t)$ through a laser cavity with linear losses and consisting of an active element with the Kerr nonlinearity, a spectral filter, a dispersive element, and the NFPI. The influence of the slow absorber responsible for the saturation of the optical transitions in the semiconductor is neglected, which is permissible because typical modulation depths ensured by a saturable absorber amount to 2%–3% in practice and the absorber operates in the strong saturation regime, when the pulse energy exceeds the saturation energy by an order of magnitude. Calculations show that the same and somewhat greater modulation depths are ensured by the amplitude response of the NFPI. Our approach agrees with that of Tsuda et al. [7], who regarded a saturable Bragg switch as a Bragg reflector the operation of which is perturbed slightly by a thin film of the semiconductor switch.

We shall consider the operation of the NFPI assuming that the pulse incident on it is many times longer than the interferometer period. Suppose that the interferometer is formed by mirrors, one of which has a 100% coefficient of reflection in terms of the field, while the other has a reflection coefficient $(R)^{1/2}$ (where $R$ is the coefficient of reflection in terms of the intensity). A semiconductor film with a transmission coefficient $L$ and a nonlinear refractive index $\chi$ is placed between the mirrors. When account is taken, after $j$ passes through the interferometer, of the phase shift

$$\varphi_j = \varphi_0 + \varphi_{j-1} + \chi[a(1-R)^{1/2}(\sqrt{R})^{j-2}L^{j-1}]^2$$

(1)

(where $\varphi_0 = 2\pi n_d l/\lambda$ is a linear shift in the interferometer), the resultant field after reflection from the interferometer...
is given by
\[ a' = a \left\{ \sqrt{R} + (1 - R)L \exp(i\phi_0) \sum_{j=0}^{\infty} \left( \frac{(1 - R)L}{\sqrt{RL}} \right)^j \right\} \]
\[ \times \exp \left\{ i(\phi_0 + \gamma |a|^2 (1 - R)L^2 \sum_{k=0}^{\infty} (\frac{(1 - R)L}{\sqrt{RL}})^{2k}) \right\}. \] (2)

We shall next consider two limiting cases: \( \sqrt{RL} \approx 1 \), which corresponds to an interferometer with a high reflectivity (NFPI), and \( \sqrt{RL} \ll 1 \), which corresponds to a device with a low reflectivity, including a saturable Bragg reflector. In the first case, the sum \( \sum_{k=0}^{\infty} (\frac{(1 - R)L}{\sqrt{RL}})^{2k} \) can then be replaced by \( j(\sqrt{RL})^2 \), whereas in the second case it can be replaced by \( \sum_{k=0}^{\infty} (\frac{(1 - R)L}{\sqrt{RL}})^{2k} \). The corresponding expressions for the field after reflection from the interferometer assume the form
\[ a' = a \left\{ \sqrt{R} + \frac{(1 - R)L \exp(i\phi_0)}{1 - \sqrt{RL} \exp[i(\phi_0 + \gamma |a|^2 (1 - R)L^2 \sum_{k=0}^{\infty} (\frac{(1 - R)L}{\sqrt{RL}})^{2k})]} \right\}, \] and
\[ a' = a \left\{ \sqrt{R} + (1 - R)L \exp(i\phi_0) \right\} \]
\[ \times \exp \left\{ i(\frac{\gamma |a|^2 (1 - R)L^2 (\sqrt{RL})^2}{1 - \sqrt{RL} \exp(i\phi_0)}) \right\}. \] (3)

We next take into account an additional phase shift owing to a possible detuning of the carrier frequency of the field relative to the centre of the pass band of the interferometer. For this purpose, we introduced an operator of the phase shift due to the detuning of the instantaneous frequency \( \Omega \) from the resonance frequency \( \exp(i\omega t) \). We go next to the time domain, we expanded this operator as a series:
\[ e^{i\Omega t} = 1 + t_\Omega \frac{\partial}{\partial t} + \frac{1}{2} t_\Omega^2 \frac{\partial^2}{\partial t^2} + \ldots, \] (4)
where \( t_\Omega \) is the reciprocal of the width of the interferometer pass band, normalised to the reciprocal of the lasing band width \( f_\text{g} \).

To a second approximation in terms of \( t \), we then obtain for the field the expression
\[ a' = \left\{ \sqrt{R} + (D_1 + iD_2) \left[ 1 + t_\Omega \frac{\partial}{\partial t} + \frac{1}{2} t_\Omega^2 \frac{\partial^2}{\partial t^2} \right] \right\} a, \] (5)
where
\[ D_1 = (1 - R)L C \cos \phi_0 - \sqrt{RL}, \quad D_2 = (1 - R)L C \sin \phi_0, \]
\[ D_3 = \left\{ \begin{array}{ll}
(1 - R)\sqrt{\frac{L}{2}} C \gamma \sin \phi_0, \quad & \text{if } R < 0.5
\end{array} \right. \]
\[ D_4 = \left\{ \begin{array}{ll}
(1 - R)\sqrt{\frac{L}{2}} C \gamma \cos \phi_0 - \sqrt{RL}, \quad & \text{if } R > 0.5
\end{array} \right. \]
\[ C = (1 - 2\sqrt{RL} \cos \phi_0 + RL^2)^{-1}. \]

The upper (lower) row in the expressions for \( D_1 \) and \( D_4 \) corresponds to the first (second) limiting case. When account is taken of all the enumerated factors, the expression describing the evolution of lasing becomes
\[ \frac{\partial a}{\partial z} = \left\{ k - i d \frac{\partial^2}{\partial t^2} + i |a|^2 \right\} \]
\[ + (D_1 + iD_2) \left[ 1 - t_{fp} \frac{\partial}{\partial t} + \frac{1}{2} t_{fp}^2 \frac{\partial^2}{\partial t^2} \right] + [D_1 D_3 - D_2 D_4 \right] \]
\[ + i[D_2 D_3 - D_1 D_4] |a|^2 + \left( \frac{4}{3} \right) \right\} \right\} a, \] (6)
where \( z \) is the longitudinal coordinate; \( t \) is the local time; \( k = z - \gamma - \sqrt{RL} \); \( z \) is the saturated gain; \( \gamma \) is the linear loss coefficient; \( d \) is the second-order coefficient representing the dispersion of the group velocity normalised to \( t_{fp}^2 \); \( \beta \) is the coefficient of phase self-modulation in the active element; \( \delta \) and \( \theta \) are the time and phase shifts, respectively. The expression in square brackets before \( |a|^2 \) in Eqn (6) takes into account the nonlinear refractive index of the semiconductor. Its real part describes the operation of an instantaneous-saturation absorber, while its imaginary part describes the additional phase self-modulation.

Fig. 1 presents the dependences of the real and imaginary parts of Eqn (6) on \( \phi_0 \) in the case of an NFPI and a saturable absorbing layer.
Bragg reflector. Evidently, the amplitudes of variation of curves \( I \) are almost an order of magnitude smaller than the corresponding values for curves \( J \), which can be explained by the smaller fraction of the field penetrating the interferometer in the case of the NFPI. It is also seen that, depending on \( \varphi_0 \), the sign of the additional amplitude modulation introduced by the nonlinear interferometer can be both positive (which is characteristic of a bleachable switch) and negative (a ‘darkening’ switch ensuring negative feedback). The sign of the additional phase self-modulation, introduced by the interferometer, can also vary and may coincide with the sign of the phase self-modulation in the active medium or may be the reverse of the latter.

The solution of Eqn (6) is sought in the form

\[
a(t) = a_0 \text{sech}(t/\tau)e^{i\omega t},
\]

where \( a_0 \), \( \tau \), and \( \psi \) are, respectively the dimensionless amplitude, duration, and chirp of a pulse; \( \omega \) is the dimensionless frequency shift of the pulse relative to the natural frequency of the interferometer. On substituting Eqn (6) in solution (7), we obtain a system of six algebraic equations for the variables \( a_0, \tau, \psi, \omega, \delta \), and \( \theta \), which is fully solvable analytically. The solution has the following form:

\[
\begin{align*}
\omega &= -\frac{D_2\tau_0}{2 + D_1\tau_0}, \\
\delta &= -\left(D_1 + \frac{2D_2\delta + D_2^2\tau_0}{2 + D_1\tau_0}\right)\tau_0, \\
\tau^2 &= \frac{(2 + D_1\tau_0)[2\psi(2d + D_2\tau_0)^2 + (\psi^2 - 1)(2 + D_1\tau_0)]}{2k(2 + D_1\tau_0)^2 + D_2\tau_0}, \\
\alpha_0^2 &= \frac{2k(2 + D_1\tau_0^2) + D_2\tau_0^2}{2(2 + D_1\tau_0)^2(D_1\tau_0^2 + D_2D_3 - \beta)} \\
&\times \frac{4d + 3\psi(2 + D_1\tau_0^2) - \psi^2(2d + D_2\tau_0^2)}{2\psi(2d + D_2\tau_0^2) + (\psi^2 - 1)(2 + D_1\tau_0^2)}.
\end{align*}
\]

The expressions for \( \psi \) and \( \theta \) are not given because they are cumbersome. It can be seen from the set of expressions (8) that the frequency shift depends solely on the linear parameters of the interferometer, in particular on the linear phase shift \( \varphi_0 \). In both cases (for the NFPI and for the saturable Bragg reflector), this shift is negative if \( 0 < \varphi_0 < \pi \), positive if \( \pi < \varphi_0 < 2\pi \), and disappears for tuning to a resonance or antiresonance, i.e. if \( \varphi_0 = 0, \pi \) (Fig. 2). It follows from Fig. 2 that the moduli of the frequency shifts in the case of the NFPI are significantly smaller and the extrema are more strongly displaced towards \( \varphi_0 = 0, 2\pi \) than in the case of the saturable Bragg reflector.

The regions of existence of the solution of Eqn (6) in the form of expression (7), in terms of the coordinates \( d \) and \( \varphi_0 \), are dark in Fig. 3. The even darker regions correspond to stable solutions (an analysis of the stability is given below). We note that, when all the radiation is reflected from the interferometer, i.e. when the field does not penetrate the interferometer (\( R = 1 \)), the solutions of Eqn (6) are conditionally stable, i.e. perturbation of the pulse energy does not evolve (the decay increment is \( A = 0 \)).

By varying the linear phase shift \( \varphi_0 \) in the interferometer, we can control the nature of the dependence of the pulse parameters on dispersion. Thus, in the case of the NFPI (Fig. 3a), two different lasing regimes may be observed depending on the interferometer tuning. For \( \varphi_0 \in [2\Delta; \pi] \cup [2(\pi - \Delta); 2\pi] \), where \( \Delta = 0.314 \) is a fraction of \( \pi \) when the solution in the region \( \varphi_0 \in [2(\pi - \Delta); 2\pi] \) is unstable, the relationships \( \tau(d) \) and \( \psi(d) \) are of the usual kind for a laser system with phase self-modulation and dispersion: the minimum duration obtained for a certain negative dispersion of the group velocity and deviation from this dispersion leads to an increase in \( \tau \), which is more marked with increase in \( d \) and weaker when \( d \) diminishes; the chirp is compensated virtually down to zero in the region of negative dispersions, but it increases, being positive, with increase in \( d \) in the region of positive dispersions.

![Figure 2. Frequency shift \( \Delta \) plotted as a function of the linear phase shift \( \varphi_0 \) in the interferometer for an NFPI (1) and a saturable Bragg reflector (2).](image)

![Figure 3. Regions of existence of a solution of Eqn (6) in the form of expression (7), shown in the \( (d, \varphi_0) \) plane for an NFPI (a) and a saturable Bragg reflector (b) with \( x = -0.96 \) (a) and \( -0.51 \) (b), \( \gamma = 0.03 \).](image)
For \( \phi_0 \in [\Delta, 2\Delta] \cup [\pi; 2(\pi - \Delta)] \), a solution exists in a limited region of dispersions near zero, which expands when \( \phi_0 \) diminishes from \( 2(\pi - \Delta) \) to \( \pi \). Here, the ultrashort pulse duration is maximal for a certain value of \( d \) near zero and tends to zero at the edges of the range of existence of the solution. The intensity tends to infinity and the chirp, compensated almost to zero in the region of negative dispersions, is negative and its modulus increases with increase in \( d \) in the region of positive dispersions. On the whole, such behaviour of the pulse parameters is nonphysical (nonrealisable). In this range of parameters we should select only those parts of the solution (near the regions with zero dispersion) in which the condition \( |\phi| \leq 1 \) holds and the limits of the approximation used to obtain Eqn (6) are not exceeded. On approach to the edges of the region of existence of a solution, this condition is no longer obeyed and one may postulate that, in this range of parameters, a pulse is transformed, to assume a profile no longer described by expression (7).

A comparison of the regions of existence of a solution of this type in terms of \( \phi_0 \) with the regions of \( \phi_0 \) in which the additional amplitude self-modulation in the interferometer (Fig. 1a, curve 1) is negative shows that they coincide. It therefore follows that the existence of this solution is associated with the presence of positive feedback ensured by the interferometer.

In the case of the saturable Bragg reflector, solutions of Eqn (6) in the form of expression (7) exist for \( \phi_0 \in [0; \Delta] \cup [\pi - 2\Delta; 2\pi] \) and this solution is stable in the region \( \phi_0 \in [\pi - 2\Delta; 2\pi - 1] \) (Fig. 3b). Solutions of three types are possible. When \( \phi_0 \) varies from \( \pi \) to \( 2\pi \), the nature of the dependences of \( \tau \) and \( \psi \) on \( d \) changes from the usual type to the opposite type (going over from a solution of the first type to one of the second type): the minimum duration shifts into the region of positive dispersions and the chirp is compensated by a positive dispersion and increases, being positive, with increase in the modulus of \( d \) in the negative region. A solution of the third type, like that for the NFPI, is associated with the presence of positive feedback (darkening absorber) and is observed in the region \( \phi_0 \in [0; \Delta] \cup [\pi - 2\Delta; \pi] \).

The minimum possible pulse durations are attained in both cases for \( \phi_0 = \pi \). For the parameters which we selected, the corresponding values of \( \tau \) are approximately 50 and 15 fs (Fig. 4). A very important feature is that, for \( \phi_0 \approx \pi \), the pulse duration remains close to the minimum value when the negative dispersions are varied over a wide range. When a saturable Bragg reflector is employed, it is possible to generate stable ultrashort pulses with a significant nonzero chirp in the region both of positive and of negative dispersions (Fig. 4b, curve 2). This is possible in the range of values of \( \phi_0 \) ensuring going over from a solution with the dependences of \( \tau \) and \( \psi \) on \( d \) of the type usual for a laser system to a solution with the opposite dependences as regards the sign of \( d \).

Under these conditions, the sign of the chirp is the same as the sign of the dispersion, the rate of increase in the ultrashort pulse duration with increase in \( |d| \) is approximately the same for negative and positive dispersions, and the minimum duration appears near zero. The \( \tau (d) \) and \( \psi (d) \) dependences have a discontinuity at a position which depends on \( \phi_0 \); for \( \phi_0 \rightarrow \pi \) or \( \phi_0 \rightarrow 2\pi \), the discontinuity is displaced to infinity. A similar discontinuity in the solution is reported in Ref. [10]. Its presence is associated with balance between the dissipative (amplitude response of the interferometer) and the nondissipative (phase self-modulation) factors in the system.

It is interesting to consider the characteristic features of the experimental realisation of the regimes governing generation of ultrashort pulses, which follow from our model. Thus, an important control parameter in the experiments is the pump energy \( E_p \) which in our model is an implicit
parameter and may be calculated with the aid of the relationship
\[
\alpha = \alpha_{\text{max}} \left( \frac{1 - \exp\left(-E_p - E/2\right)}{1 - \exp\left(-E_p - E\right)} \right),
\]
where \(\alpha_{\text{max}}\) is the gain for total inversion of populations; \(E = 2|\alpha_0|\tau\) is the energy of a pulse. The insets in Fig. 4 give the dispersion dependences of the pump power necessary for the realisation of the lasing regime illustrated in Fig. 4. In particular, it can be seen that in the case of the NFPI, the pump powers required to achieve the regime described by curve 1 are lower than those for the regime described by curve 2; furthermore, in the latter case the maintenance of minimal pulse duration with increase in the negative dispersion requires a simultaneous increase in the pump power.

In the case of the saturable Bragg reflector (Fig. 4b), the attainment of the lasing regime with the dispersion dependence of the pulse duration opposite to the traditional dependence (curve 2) requires higher pump powers, especially in the region of negative dispersions, which in fact accounts for the difficulty of observing this regime in practice.

In order to investigate the possibility of realisation of our solutions in a real laser system, it is necessary to analyse their stability. We therefore multiply Eqn (6) by \(a^*\), add it to the corresponding complex-conjugate equation, and integrate with respect to time. The relationship obtained yields the law of conservation of the pulse energy. We next introduce a small perturbation of the energy in the form \(e = e_0 e^L\) and, linearising the above relationship with respect to it, we obtain the condition ensuring decay of the pulse energy perturbation as a function of time:
\[
\dot{e} = \frac{4}{3} i t_0^2 (D_1 D_3 - D_2 D_4) + k + D_1 - \omega t_0 D_2 - \left(1 + \frac{i}{2} \frac{t_0^2 D_1}{2} \right) \frac{1 + \phi^2 + 3\omega^2\tau^2}{3\tau^2} < 0.
\]

The regions of existence of a solution, where the above condition holds, are identified in Fig. 2 by the darker shading. A comparison with the curves describing the additional amplitude self-modulation, introduced by the interferometer (Fig. 1a), shows that the regions of stability in terms of \(\phi_0\) coincide with those regions in which such self-modulation corresponds to the operation of a saturable absorber.

A model describing generation of ultrashort pulses in cw solid-state lasers with semiconductor mirrors is thus developed in the present study. It is shown that such mirrors can in general be regarded as an NFPI and they introduce an additional amplitude–phase modulation. Two limiting situations, corresponding to lasers with an NFPI and with a saturable Bragg reflector, are investigated. The optimal parameters of such devices, corresponding to the generation of the shortest ultrashort pulses, are found. The stability of the ultrashort pulses generated in these lasers is analysed.

References