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Theory for Passive Mode-Locking in Semiconductor Laser Structures Including the Effects of Self-Phase Modulation, Dispersion, and Pulse Collisions

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Abstract—We present a theory for passive mode-locking in semiconductor laser structures using a semiconductor laser amplifier and absorber. The mode-locking system is described in terms of the different elements in the semiconductor laser structure. We derive mode-locking conditions and show how other mode-locking parameters, like pulse width and pulse energy, are determined by the mode-locking system. System parameters, like bandwidth, dispersion, and self-phase modulation are shown to play an important role in mode-locking conditions and results. We also discuss the effects of pulse collisions and positions of the mode-locking elements inside the cavity on mode-locking stability and show that these effects can be easily included in the presented model. Finally, we give a number of design rules and recommendations for fabricating passively mode-locked lasers.

I. INTRODUCTION

SEMICONDUCTOR lasers have become essential components of many opto-electronic and photonic systems. In some applications, such as fiber optic telecommunication or fiber optic data processing systems, they have formed the foundation upon which these domains have developed. The generation of short optical pulses with semiconductor laser structures is crucial for high bit-rate time-division multiplexed optical systems [1], [2] and ultra long distance soliton fiber transmission systems [3]. In order to realize these optical systems, reliable optical pulse sources must be available.

There are various methods of generating short optical pulses. Gain switching and mode-locking are the two most commonly used. Gain switching is achieved by switching a laser diode on and off. An advantage of gain switching is the flexibility to change the repetition rate of the generated pulses without modifying the cavity length and the ability to directly modulate a sequence of optical pulses. However, pulse width and pulse repetition time are restricted by the electrical characteristics of the laser diode and driving electronics.

Mode-locking is another way of generating short optical pulses [4]. In mode-locking, an intracavity gain, loss, or phase element is used to lock the longitudinal modes in the semiconductor laser structure together in order to produce short optical pulses. In order to modulate these optical pulses, an external modulator is required, in contrast to gain switching. There are three different ways in which a laser structure can be mode-locked.

In active mode-locking, the optical amplifier of the semiconductor laser structure is modulated with electrical pulses that have a repetition period equal to the round-trip time of an optical pulse in the laser cavity. Only during the peak of the electrical pulse, the optical gain of the amplifier is high enough to overcome the losses in the cavity. During this short period of positive net gain, an optical pulse is generated. Short optical pulses have been generated at various repetition rates by actively mode-locking semiconductor laser structures [5]–[11].

Passive mode-locking provides an alternative approach to generating ultra-short pulses that does not require any electrical modulation. Stable and reliable monolithic passively mode-locked pulse sources can be designed by studying the properties of the components of the mode-locking system. Passive mode-locking techniques are used in many laser systems to generate short optical pulses. The key element necessary for passive mode-locking is a saturable absorber, which locks the longitudinal cavity modes in phase, leading to a short optical pulse [9], [12]–[16].

If we combine active and passive mode-locking in the same laser structure, we find the third method of mode-locking, called hybrid mode-locking. In a hybridly mode-locked laser, the optical pulses are generated in the same way as in a passively mode-locked laser, while the pulses are synchronized by an electrical signal like in an actively mode-locked one [9], [13], [14], [17].

In this paper, we give a theoretical description of the principle of passive mode-locking for semiconductor lasers. Effects of self-phase modulation in optical amplifiers have been studied earlier [18], [19]. In this paper, we describe how self-phase modulation in an amplifier can affect a semiconductor mode-locked laser. Effects of intracavity dispersion and pulse collisions have been studied in separate papers [20]–[24]. The model presented in this paper includes the effects of self-phase modulation, dispersion, and pulse collisions simultaneously and shows that these effects do not necessarily have a negative impact on the mode-locking system. The model is also applicable to any configuration of ring or Fabry–Pérot laser structures.

In Section II, we consider, as a starting point, a mode-locking system formed by a unidirectional ring laser structure. A discussion of intracavity elements and their effects on the mode-locking system will be given in Section III. In Section IV, we present a theory for passive mode-locking in semiconductor lasers. Effects of intracavity dispersion and pulse collisions are included in this theory. Section V deals with the effects of self-phase modulation in the semiconductor laser amplifier. Finally, in Section VI, we present a hybrid mode-locking scheme and some design rules and recommendations.
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in which a slow saturable amplifier and absorber are present. In our discussion, we make some suitable approximations in order to find a closed-form analytical solution to the mode-locking problem. The mode-locking system is described in terms of the bandwidth, self-phase modulation and dispersion of the system, the saturation energy of the saturable amplifier and absorber, and the ratio of the gain and absorber relaxation time to the pulse repetition time. The cavity elements that determine these parameters are described in Section II-A–II-D. In Section III, we derive the mode-locking equation and in Section IV the solution to it. In Section V, the stability of the mode-locking system is analyzed, and the conditions under which steady-state mode-locking solutions are found are derived. In this section, we also graphically show how the energy and width of the generated pulses are influenced by the different system elements. In Section VI, the effects of dispersion and self-phase modulation of the system are discussed. Finally, in Section VII, we show that the theory presented in this paper can also be applied to bidirectional ring cavities and Fabry–Perot cavities by taking into account the effects of pulse collisions that can occur in these laser structures and the position of the different system elements inside the cavity.

II. MODE-LOCKING CONFIGURATION

In this section, we derive the properties of a passively mode-locked unidirectional ring laser. The theory of passive mode-locking for (dye) lasers has been introduced and extended by Haus [20]-[22], [25]-[28]. The theory presented here follows the lines of a paper by Haus [20]. However, at several points, we have made modifications and extensions. Some of these extensions include the effects of self-phase modulation and pulse collisions on the system.

We consider a unidirectional ring configuration with a semiconductor laser amplifier (L), a saturable absorber (A), a bandwidth limiting element (B) to take into account the bandwidth and cavity loss of the system, and a dispersive element (D) representing the dispersion of the system (see Fig. 1).

Starting with our theoretical analysis, we make the following assumptions and approximations:

1) Both the amplifier and absorber are assumed to be only time-dependent. Any spatial variations of the gain and absorption coefficient inside the amplifier and absorber, respectively, are neglected. This simplification allows us to obtain analytical solutions to the mode-locking problem.

2) The dispersion of the system is determined by the wavelength and material dispersion introduced by the different system elements, like amplifier, absorber, and cavity.

3) The bandwidth of the system is determined by the amplifier/absorber. If we want to obtain analytical solutions to the mode-locking problem, we can not assign this bandwidth limitation to the gain/absorber medium. Therefore, we have introduced the bandwidth limiting element in Fig. 1, which not only takes into account the bandwidth limitation of the system, but also the cavity loss.

4) The gain relaxation time is considered long compared to the pulse width. This is a valid assumption because, in practice, the gain recovery time is about 0.2-1 ns, while the pulses generated lie normally in the low picosecond range.

5) In semiconductor lasers, the absorber relaxation time is in the order of 10–15 ps [29]; so if we consider generating pulses in the low picosecond range, the saturable absorber acts as a slow absorber.

We assume that a pulse has formed in the semiconductor laser configuration of Fig. 1, so that on the mth pass around the ring, the electric field $E_m(t)$ can be written as

$$E_m(z, t) = a_m(z, t)e^{j\omega_0 t},$$

(1)

where $a_m(z, t)$ is the envelope of the electric field at the mth pass, and $\omega_0$ is the optical carrier frequency of the electric field. We ignore the details of transverse field patterns and treat the electric field as a plane wave. Note that the Fourier transform of the envelope of the electric field $a_m(z, \omega)$ and of the electric field itself $E(z, \omega)$ are related to each other by

$$E_m(z, \omega) = a_m(z, \omega - \omega_0).$$

(2)

In the following sections, we first describe the different elements of the cavity, using the aforementioned assumptions and approximations, after which we derive the passive mode-locking equation and the solution to it.

A. Saturable Amplifier

First of all, we investigate how one single pulse is influenced by a saturable amplifier. Therefore, we consider a traveling-wave amplifier and assume that the active region dimensions of the amplifier are such that the amplifier supports a single lateral waveguide mode. We assume that the electric field inside the amplifier can be represented by (1). The evolution of the slowly varying amplitude $a(z, t)$ of the electric field along the amplifier length is then described by

$$\frac{\partial a(z, t)}{\partial z} + \frac{1}{v_g} \frac{\partial a(z, t)}{\partial t} = \frac{1}{2} [(1 - j\alpha_k)g(z, t) - \alpha_1]a(z, t),$$

(3)
where the group-velocity \( v_g = c_0/n_g \) and the group index 
\( n_g = n_{eff} + \omega(dn_{eff}/d\omega) \). The term on the right-hand side of (3) takes into account the complex internal gain 
\((1-j\alpha_L)g(z,t)\) and internal loss \( \alpha_L \) experienced by the guided mode. The linewidth enhancement factor \( \alpha_L \) represents the self-phase modulation of the amplifier.

In order to find the spatial- and time-dependent behavior of the gain of the saturable amplifier, we start from the well-known rate-equation for the carrier density inside the amplifier

\[
\frac{\partial N(z,t)}{\partial t} = \frac{I}{eV} - \frac{N(z,t)}{\tau_L} - \frac{g(z,t)I(z,t)}{\hbar \omega}, \tag{4}
\]

where \( I \) is the current injected into the active region, \( e \) the electron charge, the volume of the active region \( V = w d L \) with \( w, d, \) and \( L \) the width, thickness, and length of the active region, respectively, \( N \) the carrier density, \( \tau_L \) the carrier life-time, and \( I(z,t) \) the intensity of the optical field given by

\[
I(z,t) = C|a(z,t)|^2, \tag{5}
\]

where \( C \) is a constant with dimension \( \Omega^{-1} \). Practical values for the carrier life-time are \( \tau_L = 0.2-1 \) ns. The carrier life-time is often referred to as the gain relaxation time.

Above transparency, the gain can be linearly approximated by

\[
g(z,t) = \Gamma A_L[N(z,t) - N_{t_{\text{L}}}], \tag{6}
\]

where \( \Gamma \) is the confinement factor of the active region, \( A_L \) is the differential gain, and \( N_{t_{\text{L}}} \) is the carrier density needed to obtain transparency in the amplifier. Typical values are \( \Gamma = 0.3, A_L = 2-3 \cdot 10^{-16} \) cm\(^2\), \( N_{t_{\text{L}}} = 1-2 \cdot 10^{18} \) cm\(^{-3}\), and \( \alpha_L = 2-6 \). Using (5) and (6), we can rewrite (4) and obtain the spatial- and time-dependent behavior of the gain

\[
\frac{\partial g(z,t)}{\partial t} = \frac{g(z,t) - g_0 - g(z,t) P(z,t)}{\tau_L}, \tag{7}
\]

with the small signal gain \( g_0 = \Gamma A_L[I\tau_L/eV - N_{t_{\text{L}}}] \), the saturation energy of the amplifier \( E_{\text{sat}_L} = \hbar \omega_0 \sigma_L / A_L \) with \( \hbar \) Dirac’s constant. The mode cross-section \( \sigma_L \) is typically \( 0.3-1 \) \( \mu \)m\(^2\), leading to \( E_{\text{sat}_L} = 10-20 \) pJ.

The power \( P(z,t) \) of the electric field inside the active region is given by

\[
P(z,t) = \sigma_L I(z,t). \tag{8}
\]

In order to simplify the former described spatial- and time-dependent behavior of the amplifier gain and electric field, we make the transformation to a reference frame moving with the pulse [30]

\[
t = t - \frac{z}{v_g}, \tag{9}
\]

We then find the following set of equations that describe the amplifier:

\[
\frac{\partial a(z,\tau)}{\partial z} = \frac{1}{2}[(1-j\alpha_L)g(z,\tau) - \alpha_L a(z,\tau) \tag{10}\]

\[
\frac{\partial g(z,\tau)}{\partial \tau} = - \frac{g(z,\tau) - g_0}{\tau_L} - \frac{g(z,\tau) P(z,\tau)}{E_{\text{sat}_L}}. \tag{11}
\]

The amplification for a pulse passing through the amplifier can then be represented by [18]

\[
a'_m(\tau) = a_m(\tau)e^{(1/2)(1-j\alpha_L)h(\tau)}, \tag{12}
\]

where \( a_m(\tau) \) and \( a'_m(\tau) \) are the input and output pulse, respectively, and the integrated gain \( h(\tau) \) is given by

\[
h(\tau) = \int_0^\tau g(z,\tau) dz. \tag{13}
\]

We use a simplified model for the saturable amplifier, in which the spatial variation of the gain is omitted. Recalling (7), we find for the time-dependent gain the following differential equation:

\[
\frac{dg(\tau)}{d\tau} = - \frac{g(\tau) - g_0 - g(\tau) P_m(\tau)}{\tau_L}, \tag{14}
\]

with solution

\[
g(\tau) = g_0 \exp \left[ - \frac{E_m(\tau)}{E_{\text{sat}_L}} \right], \tag{17}
\]

where \( g_0 \) is the gain of the amplifier before arrival of the pulse, and \( E_m(\tau) \) is defined by

\[
E_m(\tau) = \sigma_L C \int_{-\infty}^{\tau} |a_m(\tau')|^2 d\tau'. \tag{18}
\]

The profile of the gain is assumed to be Lorentzian. If the bandwidth of the amplifier is much wider than the pulse bandwidth, we can write the gain \( G(\omega) \) in the frequency domain as

\[
G(\omega) = \frac{G(\omega_p)}{1 + j \frac{\omega - \omega_p}{\omega_L}}, \tag{20}
\]

where the amplifier has its peak gain at the frequency \( \omega_p \) and the peak gain \( G(\omega_p) \) is time-independent. The bandwidth of the amplifier \( \omega_L \) is typically \( 10-50 \cdot 10^{12} \) rad/s. Expanding (20) to second order in frequency and introducing the change
in the carrier frequency from the frequency at peak gain as
\[ \Delta \omega = \omega_0 - \omega_p, \]
we have
\[
G(\omega) = G(\omega_p) \left[ 1 - \left( \frac{\Delta \omega}{\omega_L} \right)^2 - j \frac{\Delta \omega}{\omega_L} \right] \left[ \frac{1}{\omega_L^2} \frac{d^2}{d\tau^2} a_m(\tau) \right].
\]

Using (2) and Fourier analysis, we can make the following transformations from the frequency domain to the time domain for the field envelope:
\[
\left( \frac{j \omega - \omega_0 - \omega_p}{\omega_L} \right)^2 a_m(\omega) \rightarrow \frac{1}{\omega_L^2} \frac{d^2}{d\tau^2} a_m(\tau).
\]

Taking the length of the amplifier \( l_L \), we find for the transfer of the field envelope through the amplifier
\[
a_m'(\tau) = e^{(1/2)(1-j\alpha_L)} h_L(\tau) a_m(\tau),
\]
where the transfer function \( h_L(\tau) \) of the amplifier is defined by
\[
h_L(\tau) = G(\omega_p) l_L \left[ 1 - \left( \frac{\Delta \omega}{\omega_L} \right)^2 - j \frac{\Delta \omega}{\omega_L} \right] \left[ \frac{1}{\omega_L^2} \frac{d^2}{d\tau^2} a_m(\tau) \right].
\]

Examining (24), we recognize the bandwidth limitation of the amplifier and see that the action of the amplifier upon the pulse causes a shift in time \( (d/\tau^2) \) and a diffusion in time \( (d^2/\tau^2) \).

In order to obtain analytical solutions to the mode-locking problem, we can not assign the bandwidth limitation to the gain element. Therefore, we have to introduce a bandwidth limiting element in the mode-locking configuration (see Fig. 1) that is described in Section II-C. Omitting now the bandwidth limitation and using (17), we find for the transfer of the field envelope through the amplifier (23), now with \( h_L(\tau) \) of the amplifier defined by
\[
h_L(\tau) = g(\tau) l_L,
\]
where \( g(\tau) \) is defined according to (17).

Next, we investigate what happens if succeeding pulses enter the saturable absorber. Therefore, we derive a relation between time-dependent gain \( g(\tau) \) and the small signal gain \( g_0 \). We assume that the gain at the arrival of a pulse equals \( g_0 \) and at the end of a pulse \( g_f \). Between two pulse arrivals, the gain can only relax up to \( g_0 \). Then the next pulse comes along and saturates the gain to \( g_f = g_0 \exp (-E_0/E_{sat}) \).

After the pulse passage, the gain relaxes back to \( g_0 \) according to
\[
g(\tau) = g_0 - (g_0 - g_f) \exp \left( -\frac{\tau}{\tau_A} \right).
\]

Eliminating \( g_f \) and solving (26) for \( g_0 \) gives at arrival of the next pulse at \( \tau = T_P \)
\[
g_0 = g_0 \left\{ \exp \left( \frac{T_P}{\tau_L} \right) - \exp \left( \frac{-E_0}{E_{sat}} \right) \right\}. \tag{27}
\]

Note that the time between two pulse arrivals \( T_P \) for the unidirectional ring cavity displayed in Fig. 1 equals the cavity round-trip time \( T_R \).

**B. Saturable Absorber**

We can do the same calculations for the saturable absorber, as for the saturable amplifier. In analogy, we then find for the absorption
\[
q(\tau) = q_0 \exp \left[ \frac{E_m(\tau)}{E_{sat}} \right] \tag{28}
\]
with \( E_{sat} = h\omega_0 \sigma_A / \Gamma A_A \) where \( \sigma_A \) is the absorber mode cross-section and where \( A_A \) is the differential absorption. The absorption of the absorber before arrival of the pulse equals \( q_0 \). If the absorber fully relaxes between the pulse arrivals, we can replace \( q_0 \) by \( q_0 \), defined by
\[
q_0 = \Gamma A_A \left( N_{t_A} - \frac{I_{t_A}}{eV} \right), \tag{29}
\]
where \( N_{t_A} \) is the carrier density needed to obtain transparency in the absorber, and \( \tau_A \) is the absorber relaxation time. Taking the length of the absorber \( l_A \), we find for the evolution of the field envelope passing through the absorber
\[
d_m'(\tau) = e^{(1/2)(1-j\alpha_A)} h_A(\tau) a_m(\tau), \tag{30}
\]
where \( \alpha_A \) the linewidth enhancement factor of the absorber and where the transfer function \( h_A(\tau) \) of the absorber is defined by
\[
h_A(\tau) = Q(\omega_p) l_A \left[ 1 - \left( \frac{\Delta \omega}{\omega_L} \right)^2 - j \frac{\Delta \omega}{\omega_L} \right] \left[ \frac{1}{\omega_L^2} \frac{d^2}{d\tau^2} a_m(\tau) \right]. \tag{31}
\]

where we have taken the bandwidth limitation of the absorber equal to the one of the amplifier. The bandwidth limitation has to be dropped again in order to find analytical solutions to the mode-locking problem. Using (28), we now find for the transfer of the field envelope (30), with the transfer function \( h_A(\tau) \) of the absorber defined by
\[
h_A(\tau) = -q(\tau) l_A \tag{32}
\]
where \( q(\tau) \) defined according to (28). Typical values for the saturable absorber are \( \tau_A = 10-15 \) ps, \( A_A = 10^{10} \cdot 10^{-16} \).
cm², $\sigma_A = 0.3-1\ \mu m^2$, $N_{iA} = 0.5-1 \cdot 10^{18} \ cm^4$, $E_{satA} = 1-10\ \text{pJ}$, and $\alpha_A = 2-6$.

In the same way as we did for the saturable amplifier, we can derive the following relation between the small-signal absorption $q_0$ and the absorption before arrival of the pulse $q_i$:

\[
q_0 = q_i \frac{\exp \left( \frac{T_p^A}{\tau_A} \right) - \exp \left( \frac{-E_{satA}}{E_{satA}} \right)}{\exp \left( \frac{T_p^A}{\tau_A} \right) - 1},
\]

where $T_p^A$ is the time between two succeeding pulse arrivals in the saturable absorber. Note that this time does not need to equal the time between two pulse arrivals in the saturable amplifier but depends on the position of the amplifier and absorber in the cavity and the number of pulses inside the cavity.

### C. Bandwidth Limiting Element

To include the bandwidth limitation of the gain and absorber, we have introduced a bandwidth-limiting element in the mode-locking system. In this bandwidth-limiting element, the loss of the cavity is also included. As we have seen in (24), the bandwidth limitation of the amplifier causes a shift in time and a diffusion in time of the field envelope of the electric field. The transfer function of the bandwidth limiting element can be represented by

\[
h_B(\tau) = -\alpha_C l_C \left[ 1 + \left( \frac{\Delta \omega}{\omega_L} \right)^2 + j \frac{\Delta \omega}{\omega_L} \right] \left( 1 - 2j \frac{\Delta \omega}{\omega_L} \right) \frac{d}{d\tau} \frac{d^2}{d\tau^2},
\]

where $\alpha_C$ is the cavity loss of the system and $l_C$ the cavity length and where we have used

\[
G(\omega)L = Q(\omega)l_A + \alpha_C l_C.
\]

The change of the field envelope by the bandwidth-limiting element is given by

\[
a_m^1(\tau) = e^{(1/2)h_B(\tau)}a_m(\tau),
\]

where the transfer function $h_B(\tau)$ is defined by (34).

### D. Dispersive Element

In order to take into account the phase changes of the field envelope and dispersion of the system, we introduce a dispersive element. As mentioned before, we assume that the dispersion of the system is determined by the wavelength and material dispersion introduced by the different system elements. The phase function $\psi(\omega)$ of the dispersive element in the frequency domain can be represented by

\[
\psi(\omega) = \varphi_0 + D(\omega - \omega_0)^2,
\]

where $\varphi_0$ is a constant phase shift, and $D$ is a measure for the velocity dispersion in the system, defined by

\[
D = l_L \left( \frac{d^2k}{d\omega^2} \right) + l_A \left( \frac{d^2k}{d\omega^2} \right) + l_C \left( \frac{d^2k}{d\omega^2} \right).
\]

The pulse repetition time is thus determined by the time needed to travel through the cavity.

Making the transformation back to the time domain for $\varphi(\omega)$, we find for the transfer of the field envelope through the dispersive element

\[
a_m^1(\tau) = e^{(1/2)h_B(\tau)}a_m(\tau),
\]

where the transfer function $h_B(\tau)$ is defined by

\[
h_B(\tau) = -j \left( \varphi_0 - D \frac{d^2}{d\tau^2} \right).
\]

### III. MODE-LOCKING EQUATION

Knowing the influences of the cavity elements on the field envelope $a_m(\tau)$, we are able to derive the mode-locking equation for the system. For the pulse envelope $a_{m+1}(\tau)$ after one cavity round-trip we find

\[
a_{m+1}(\tau) = \exp \left[ \frac{1}{2} \left( 1 - j\alpha_A \right) q(\tau) l_A + \frac{1}{2} \left( 1 - j\omega_L \right) q(\tau) l_L \right.

\left. - \frac{1}{2} \alpha_C l_C \left[ 1 + \left( \frac{\Delta \omega}{\omega_L} \right)^2 + j \frac{\Delta \omega}{\omega_L} \right] \left( 1 - 2j \frac{\Delta \omega}{\omega_L} \right) \frac{d}{d\tau} \frac{d^2}{d\tau^2} \right] \left( \varphi_0 - D \frac{d^2}{d\tau^2} \right) a_m(\tau).
\]

We assume that a pulse propagating through any of the elements of the mode-locking system, amplifier, saturable absorber, etc., is modified only slightly (say 20% gain or loss) on one round-trip. This assumption allows us to expand the exponential in (42) to first order in the argument, leading to

\[
a_{m+1}(\tau) = \left\{ 1 - \frac{1}{2} \left( 1 - j\alpha_A \right) q(\tau) l_A + \frac{1}{2} \left( 1 - j\omega_L \right) q(\tau) l_L \right.

\left. - \frac{1}{2} \alpha_C l_C \left[ 1 + \left( \frac{\Delta \omega}{\omega_L} \right)^2 + j \frac{\Delta \omega}{\omega_L} \right] \left( 1 - 2j \frac{\Delta \omega}{\omega_L} \right) \frac{d}{d\tau} \frac{d^2}{d\tau^2} \right] \left( \varphi_0 - D \frac{d^2}{d\tau^2} \right) a_m(\tau).
\]

Now $a_{m+1}(\tau)$ does not need to be equal to $a_m(\tau)$ because some of the pulse shaping may lead to delays or advances of the pulse.
the pulse. For the round-trip condition, we then find
\[ a_{m+1}(\tau) = a_m(\tau + \Delta T), \tag{44} \]
where \( \Delta T \) represents the time shift of the pulse. By this definition, a positive value of \( \Delta T \) represents a shift backward in time. If we now expand \( a_m(\tau + \Delta T) \) to first order in \( \Delta T \), we find
\[ a_{m+1}(\tau) = a_m(\tau) + \Delta T \frac{da_m(\tau)}{d\tau}. \tag{45} \]
Substituting (17), (28) and (45) into (43) and omitting the subscript \( m \), we find the passive mode-locking equation
\[
\begin{aligned}
1 + (1 - j\alpha_A)Q_i \exp \left[ -\frac{E(\tau)}{E_{\text{sat}A}} \right] \\
- (1 - j\alpha_L)G_i \exp \left[ -\frac{E(\tau)}{E_{\text{sat}A}} \right] \\
+ \xi^2 + j(\xi + \psi) + (1 + \delta - 2j\xi) \frac{1}{\omega_L} \frac{d}{d\tau} \\
- (1 + j\delta) \frac{1}{\omega_L} \frac{d^2}{d\tau^2}
\end{aligned}
\]
where we have introduced normalized amplifier and absorber parameters
\[
Q(\tau) = \frac{q(\tau)L}{\alpha_{CL}}
\quad Q_k = \frac{q_{kL}}{\alpha_{CL}} \quad \text{with} \quad k = 0, i, f
\quad G(\tau) = \frac{g(\tau)L}{\alpha_{CL}}
\quad G_k = \frac{g_{kL}}{\alpha_{CL}} \quad \text{with} \quad k = 0, i, f
\]
and normalized system parameters
\[
\begin{aligned}
\mathcal{D} & = \frac{\omega_L^2 D}{2\alpha_{CL}} \\
\psi & = \frac{\tau_0}{2\alpha_{CL}} \\
\delta & = \frac{\omega_L \Delta T}{2\alpha_{CL}} \\
\xi & = \frac{\Delta \omega}{\omega_L}.
\end{aligned}
\tag{47}
\]
Typical values are \( \alpha_C = 1-2 \text{ cm}^{-1}, l_C = 0.1-1 \text{ cm}, \omega_L = 10-50 \times 10^{12} \text{ rad/s}, \) and \( D = 1-10 \times 10^{-26} \text{ s}^2 \), leading to \( \mathcal{D} = 5-25 \). Note that
\[ G(\tau) = -1 - Q(\tau) + G(\tau) \tag{49} \]
is the net gain parameter. This parameter plays a very important role in the mode-locking problem.

IV. SOLUTION TO THE MODE-LOCKING EQUATION

Equation (46) is a nonlinear differential equation that looks quite complicated. A simple solution of (46) is obtained if we expand the exponentials to second order in their argument. Because \( E_{\text{sat}L} > E_{\text{sat}A} \), as will become clear later on, we may break off the expansion of \( \exp \left[ E(\tau)/E_{\text{sat}L} \right] \) with the first-order term, so that we obtain
\[
\begin{aligned}
1 + (1 - j\alpha_A)Q_i - (1 - j\alpha_L)G_i \\
- \left[ (1 - j\alpha_A)Q_i - (1 - j\alpha_L)G_i \right] \frac{E(\tau)}{E_{\text{sat}A}} \\
+ \frac{1}{2} \left( 1 - j\alpha_A \right)Q_i \left[ \frac{E(\tau)}{E_{\text{sat}A}} \right]^2 \\
+ (1 + \delta - 2j\xi) \frac{1}{\omega_L} \frac{d}{d\tau} - (1 + j\delta) \frac{1}{\omega_L} \frac{d^2}{d\tau^2}
\end{aligned}
\]
where we have introduced the stability parameter
\[ s = \frac{E_{\text{sat}L}}{E_{\text{sat}A}} \tag{51} \]
The differential equation (50) can be solved analytically by introducing a guess [22] for the pulse envelope \( a(\tau) \)
\[ a(\tau) = a_0 \left[ \text{sech} \left( \frac{\tau}{\tau_0} \right) \right]^{1+j\beta}, \tag{52} \]
where the amplitude \( a_0 = \sqrt{V_i E_{\text{sat}A}/2\sigma C L} \) with \( V_i = E_i/E_{\text{sat}A} \). The pulse shape denoted by (52) is a hyperbolic secant. As a result of (52), we have
\[
\begin{aligned}
E(\tau) & = \frac{V_i E_{\text{sat}A}}{2} \left[ 1 + \tanh \left( \frac{\tau}{\tau_0} \right) \right] \\
\tau_0 \frac{d}{d\tau} a(\tau) & = - (1 + j\beta) \tanh \left( \frac{\tau}{\tau_0} \right) a(\tau) \\
\tau_0^2 \frac{d^2}{d\tau^2} a(\tau) & = - (2 + 3j\beta - \beta^2) \text{sech}^2 \left( \frac{\tau}{\tau_0} \right) \\
& + (1 + j\beta)^2 a(\tau).
\end{aligned}
\tag{53}
\tag{54}
\tag{55}
\]
Substituting the above equations into (50) and equating the coefficients of the terms \( a(\tau), \frac{d}{d\tau} a(\tau), \) and \( \frac{d^2}{d\tau^2} a(\tau) \), yields three complex algebraic equations. The real and imaginary parts of these three complex equations are
\[
\begin{aligned}
& 2 - \beta^2 \frac{D}{\omega_L^2 \tau_0} - \frac{3\beta D}{\omega_L^2 \tau_0} - \frac{1}{8} \alpha_A Q_i V_i^2 = 0 \\
& \mathcal{D}(2 - \beta^2) \frac{D}{\omega_L^2 \tau_0} + 3\beta \frac{D}{\omega_L^2 \tau_0} + \frac{1}{8} \alpha_A Q_i V_i^2 = 0 \\
& - \frac{1 + \delta}{\omega_L \tau_0} - \frac{2\beta \xi}{\omega_L \tau_0} - \frac{1}{2} \left( \frac{G_i}{s} \right) V_i + \frac{1}{4} \alpha_Q V_i^2 = 0 \\
& - \beta(1 + \delta) + \frac{2\xi}{\omega_L \tau_0} + \frac{1}{2} \left( \alpha_A Q_i - \alpha \frac{G_i}{s} \right) V_i \\
& - \frac{1}{4} \alpha_A Q_i V_i^2 = 0.
\end{aligned}
\tag{56}
\tag{57}
\tag{58}
\tag{59}
\]
1 + Q_i - G_i + \xi^2 - \frac{1}{2} \left( Q_i - \frac{G_i}{s} \right) V_i + \frac{1}{4} Q_i V_i^2
- \frac{1 - \beta^2}{\omega^2 L^2 \gamma_0^2} + \frac{2 \beta D}{\omega^2 L^2 \gamma_0^2} = 0
(60)

\psi + \xi - \frac{2 \beta}{\omega^2 L^2 \gamma_0^2} \frac{D(1 - \beta^2)}{\omega^2 L^2 \gamma_0^2} - \alpha_A Q_i + \alpha_L G_i
+ \frac{1}{2} \left( \alpha_A Q_i - \alpha_L \frac{G_i}{s} \right) V_i - \frac{1}{4} \alpha_A Q_i V_i^2 = 0
(61)

These six equations have six unknowns: \tau_0, V_i, \beta, \delta, \xi, and \psi.
In principle, the system can be solved consistently for different
Q_i and G_i.

V. MODE-LOCKING CONDITIONS

In this section, we derive the boundaries for stable pulse
mode-locking solutions. These boundaries are graphically
shown in the \( G_0-Q_0 \) plane. We start with describing the
conditions for stable mode-locking.

The net gain as a function of the normalized energy \( U(\tau) = E_i(\tau)/E_{sat} \) has the dependence
\[ G_T[U(\tau)] = -(1 + Q_i - G_i)
+ \left( Q_i - \frac{G_i}{s} \right) U(\tau) - \frac{1}{2} Q_i U^2(\tau).
(62) \]

In order to avoid that perturbations preceding the pulse can
grow, the net gain preceding the pulse \( (U(\tau) = 0) \) must be
negative, leading to
\[ G_i < 1 + Q_i.
(63) \]

Note that in order to start the mode-locking system, we must have
\[ G_0 > 1 + Q_0.
(64) \]

It is important to note that if the amplifier relaxes fully
between two incoming pulses, i.e. \( G_i = G_0 \), we can not fulfill
(63). As \( Q_i \leq Q_0 \) and \( G_i = G_0 \), we have \( G_i > 1 + Q_i \), which
is not in agreement with (63). So (63) together with (64) leads
to the mode-locking condition \( \tau_L/T_p \geq 1 \), i.e., the amplifier
should not recover completely between two succeeding pulse
arrivals. In order to avoid that perturbations following the pulse
can grow, we must satisfy the stability criterion that the net
gain is negative after passage of the pulse \( (U(\tau) = V_i) \), leading to
\[ -(1 + Q_i - G_i) + \left( Q_i - \frac{G_i}{s} \right) V_i - \frac{1}{2} Q_i V_i^2 < 0.
(65) \]

A last condition for the pulse is that the loss must be converted
into gain with rising pulse energy, which means that the
coefficient of \((Q_i - G_i/s)\) in (62) must be positive, leading to
\[ G_i < s Q_i.
(66) \]

First, we consider the mode-locking system when dispersion
and self-phase modulation are absent.

If dispersion and self-phase modulation are absent \((D = 0,
\alpha_L = 0, \text{and} \ \alpha_A = 0)\), we find that the pulse is not chirped

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{Stability boundary (grey area between dashed lines), laser threshold (dash-dotted line), and equipulse energy lines (solid lines) as a function of normalized small signal gain and absorption for \( \tau_L/T_p = 5, \tau_A/T_p = 0.1, \ s = 6, \ D = 0, \ \alpha_A = 0, \ \text{and} \ \alpha_L = 0. \)
\end{figure}

\( (\beta = 0)\), that no phase shift is present on one round-trip \( (\psi = 0)\), and that the carrier frequency coincides with the peak gain
frequency \((\xi = 0)\).

In Fig. 2, we have drawn the stability boundaries as a
function of the normalized small signal gain \( G_0 \) and normalized
small signal absorption \( Q_0 \).

The area in which stable mode-locking is possible is given
by the grey area between the two dashed lines. The laser
threshold \( G_0 = 1 + Q_0 \) is given by the dash-dotted line. We
have also plotted the normalized pulse energy \( V_i \) in the \( G_0-Q_0 \)
plane. The solid lines are equipulse energy lines, i.e., they
connect points in the \( G_0-Q_0 \) plane that produce mode-locked
pulses with the same energy.

As we can see in Fig. 2, the energy of the pulse increases if
the gain is raised. For each value of the normalized small signal
absorption \( Q_0 \), there is a minimum pulse energy required
to obtain steady state mode-locking. This is understandable
because pulses with low energy can only pull down the gain
by small amounts. However, the gain must be pulled down
below the loss at the wings of the pulse. Otherwise, the pulse
is unstable to noise perturbations preceding and following
the pulse. Consequently, stable pulse operation requires a
minimum pulse energy, which can be derived from \( G_i = 1 + Q_i \)
and from
\begin{align*}
G_0 &= G_i \left\{ \frac{\exp \left( \frac{T_p}{\tau_L} \right) - \exp \left( \frac{-E_i}{\tau_{sat}} \right)}{\exp \left( \frac{T_p}{\tau_L} \right) - 1} \right\}.
(67) \]
\end{align*}

These equations lead to a minimum pulse energy
\begin{align*}
E_{\text{min}}^{\text{sat}} &= - \ln \left\{ \exp \left( \frac{T_p}{\tau_L} \right) - \frac{G_0}{1 + Q_i} \left[ \exp \left( \frac{T_p}{\tau_L} \right) - 1 \right] \right\}.
(68) \]
\end{align*}
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Fig. 3. Stability boundary (grey area between dashed lines), laser threshold (dash-dotted line) and equipulse width lines (solid lines) as a function of normalized small signal gain and absorption for $\tau_L/\tau_p = 5$, $\tau_A/\tau_p = 0.1$, $s = 6$, $\mathcal{D} = 0$, $\alpha_A = 0$, and $\alpha_L = 0$.

Fig. 4. Stability boundary (grey area between dashed lines), laser threshold (dash-dotted line) and equipulse energy lines (solid lines) as a function of normalized small signal gain and absorption for $\tau_L/\tau_p = 5$, $\tau_A/\tau_p = 0.1$, $s = 6$, $\mathcal{D} = 10$, $\alpha_A = 0$, and $\alpha_L = 0$.

where $Q_i$ is given by

$$Q_i = Q_0 \left[ \frac{\exp \left( \frac{T_p^A}{\tau_A} \right) - 1}{\exp \left( \frac{T_p^A}{\tau_A} \right) - \exp \left( \frac{-E_t}{E_{sat_A}} \right) \text{boundary}} \right].$$

(69)

Of course there is also a maximum value of the pulse energy if we want to obtain stable mode-locking. If the energy is raised above this value, then there is a positive net gain $G_T$ for perturbations preceding the pulse. This means that the system will eventually work under continuous wave (the net gain $G_T$ is always positive).

Furthermore, we can draw a very important conclusion from the analysis. The stability regime has an intersection point with the threshold line. This intersection point is given by $Q_0 = 1/(s - 1)$ and $Q_i = s/(s - 1)$. We see that for values of $s > 1$, this point moves toward the origin. For values of $s$ approaching 1, the intersection point moves toward infinity. Values of $s < 1$ are not allowed because these do not fulfill (66). If we thus want to obtain a stability area that is not too far from the origin, we have to make sure that $s > 1$. Physically, this means that the amplifier must be harder to saturate than the absorber. Another important advantage of a larger value for $s$ is that the stable mode-locking area becomes larger and that the area moves toward the threshold line. So in order to have a mode-locking system that has a wide range of stable operation and a low threshold, we need a value of $s$ as large as possible. In Fig. 3, we have drawn, in the same way as we did in Fig. 2, equipulse width lines in the $Q_0$-$Q_i$ plane.

The pulse width $\tau_\text{w}$ is mainly determined by the bandwidth of the system. From the normalized pulse width $1/\omega_L\tau_0$, we find that possible pulse widths in the order of 50–100 fs can be obtained.

Finally, we note that we have to check if the system is self-starting and stable against self-pulsations. For this analysis, we refer to a paper by Haus [26].

VI. EFFECTS OF DISPERSION AND SELF-PHASE MODULATION

A. Dispersion

We now investigate the influence of dispersion on the mode-locking system. If dispersion is not neglected ($\mathcal{D} \neq 0$), we find that the pulse is chirped ($\beta \neq 0$) and that a phase shift is present on one round-trip ($\psi \neq 0$). The carrier frequency of the electric field does also not coincide with the frequency at peak gain ($\xi \neq 0$). In Figs. 4 and 5, we give the equipulse energy lines and equipulse width lines in the $Q_0$-$Q_i$ plane.

As we can see in Fig. 6, the pulse is broadened if dispersion is present in the mode-locking system. Dispersion not only
causes the pulse width to increase but simultaneously poses a chirp on the pulse. Since we have assumed that self-phase modulation in the amplifier and absorber are absent ($Q_A = 0$ and $A_A = 0$), the chirp caused by dispersion cannot be compensated by the amplifier nor by the absorber. The pulse width is minimum and unchirped only if dispersion is absent.

As can be seen from Fig. 6, a normalized dispersion $D = 20$ reduces the achievable pulse width with a factor of about 4. Typical values for pulse widths that can be achieved are 200–500 fs.

B. Self-Phase Modulation

We now introduce self-phase modulation in the system, by setting $\alpha_A \neq 0$ and $\alpha_L \neq 0$. We have drawn again the normalized pulse energy, pulse width, and the chirp parameter as a function of the dispersion in Fig. 7 for $\alpha_A = \alpha_L = 2$ and $\alpha_A = \alpha_L = 6$.

From Fig. 7, we see that the self-phase modulation of the amplifier and the absorber causes a serious chirp on the pulse if the dispersion is positive. This chirp increases with increasing self-phase modulation of the amplifier and absorber. Due to the fact that for positive dispersion, the pulse exhibits a large chirp, the bandwidth of the system is used less efficiently. This leads to a larger value for pulse width for positive dispersion, as can be seen in Fig. 7.

For negative dispersion, almost no chirp is imposed on the pulse, because the chirp that is introduced by negative dispersion is compensated by the chirp caused by the amplifier and absorber. Thus, the presence of both dispersion and self-phase modulation do not necessarily have a negative impact on the mode-locking system.

Finally, it is to be noted that for the chosen parameters ($Q_1 = 2$ and $\alpha_1 = 2.8$) no mode-locking solutions are found for specific ranges of the dispersion. These areas are indicated by the grey areas in Fig. 7.

VII. EFFECTS OF PULSE COLLISIONS

Colliding-pulse mode-locking is an improvement in mode-locked lasers. The new feature of colliding pulse mode-locking is two counter-propagating pulses that are synchronized to precisely overlap in the saturable absorber. These overlapping pulses create a transient standing wave pattern in the optical field and, consequently, a transient grating in the absorption of the absorber that shortens the optical pulses in a very effective manner. In this section, we describe the colliding-pulse mode-locking principle, and we show that it is very easy to make use of this principle in ring cavity lasers and Fabry–Pérot cavity lasers. We prove that a ring configuration of the mode-locking system leads to a higher stability parameter than a conventional Fabry–Pérot cavity and thus, to a better performance of the mode-locking system.

A. Colliding-Pulse Configurations

We start from a basic noncolliding pulse cavity configuration, which occurs if we have a unidirectional laser cavity. This unidirectional mode-locking (UM) configuration is displayed in Fig. 8, where the dispersion and bandwidth-limiting element have been omitted for convenience. The unidirectional behavior can, of course, only occur in ring cavities.

In the UM configuration of Fig. 8, no pulse collisions can occur inside the saturable absorber because only one pulse is circulating inside the cavity. This single pulse saturates the absorber and the amplifier section. If we change the UM ring configuration of Fig. 8 into a conventional Fabry–Pérot cavity and thus, to a better performance of the mode-locking system.
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has advantages compared to the UM configuration, due to the grating that is formed inside the absorber.

A final and even better improvement can be obtained if we use a bidirectional ring configuration in which two counter-propagating pulses collide in the saturable absorber. This so-called colliding pulse mode-locking (CPM) configuration is shown in Fig. 10. The two counter-propagating pulses in this ring laser configuration always meet in the saturable absorber because the intracavity loss is sharply reduced for that condition. This sharp reduction in intracavity loss occurs because the absorber is saturated by two pulses at the same time.

Colliding pulse mode-locking can also be realized in conventional Fabry-Perot cavities. In such cavities, the saturable absorber must be placed at a submultiple $m \geq 2$ of the cavity length, and the amplifier gain must be adjusted to support $m$ pulses. Under these conditions, two counter-propagating pulses always simultaneously saturate the absorber. In Fig. 11, we have shown the configurations for two, three, and four intracavity pulses. The special case where $m = 1$ gives again the SPCM configuration shown in Fig. 9.

In the next sections, we examine the effects of pulse collision on the pulse shaping performed by the absorber.

**B. Pulse Shaping by an Absorption Grating**

The effects of pulse collisions in saturable absorbers have been studied in dye lasers [23], [24]. If two counter-propagating pulses collide in a saturable absorber (SCPM or CPM configuration), they form a standing wave in the absorber. At the peaks of the standing wave, the power of the electrical field is maximum, and the absorption is saturated to a high extent. The net remaining absorption therefore has a minimum at the peaks of the standing wave. At the minima of the standing wave, the absorption is not as much saturated as at the peaks, and thus, the remaining absorption is higher at those points. This leads to an absorption grating inside the saturable absorber. Since the absorption does not have to be lowered to transparency everywhere in the saturable absorber in order to obtain a low loss state, it is very likely that the saturable absorber bleaches at a lower pulse energy than would be necessary if the standing wave did not exist.

In our analysis of pulse collisions, we assume that the electromagnetic field inside the amplifier can be represented by two harmonic plane waves, one traveling in the positive $z$-direction and one traveling in the negative $z$-direction, both with an optical carrier frequency $\omega_0$ and with a slow varying envelope $a_R(z, t)$ and $a_L(z, t)$, where the subscript $R$ and $L$ refer to the wave traveling in the positive and negative $z$-direction, respectively.

Generally, the saturable absorbers used in semiconductor lasers are slow absorbers, so we may use

$$\frac{\partial q(z, t)}{\partial t} = -q(z, t) \frac{P(z, t)}{E_{sat}}. \quad (70)$$
We now spatially expand the absorption \( q(z,t) \) in Fourier elements to second order

\[
q(z, t) = q_A(z, t) + q_B(z, t) (e^{2ikz} + e^{-2ikz}) = q_A(z, t) + 2q_B(z, t) \cos (2kz).
\]

(71)

By substituting (71) into (70) and neglecting fourth-order terms, we find the following equations:

\[
\frac{\partial q_A}{\partial t} = - \frac{\sigma C}{E_{\text{sat}} A} [q_A(|a_R|^2 + |a_L|^2)] + q_B(a_R a_L^* + a_R^* a_L) + \frac{1}{2} q_A(a_R a_L + a_R^* a_L^*).
\]

(72)

\[
\frac{\partial q_B}{\partial t} = - \frac{\sigma C}{E_{\text{sat}} A} [q_B(|a_R|^2 + |a_L|^2)] + \frac{1}{2} q_A(a_R a_L + a_R^* a_L^*).
\]

(73)

The change of the electric field within the absorber is described by Maxwell’s equations. Using the rotating wave approximation [23], we find for the amplitudes of the harmonic waves traveling in the positive and negative \( z \)-direction

\[
\frac{\partial a_R}{\partial z} + \frac{1}{v_g} \frac{\partial a_R}{\partial t} = - \frac{1}{2} (q_A a_R + q_B a_L),
\]

(74)

\[
\frac{\partial a_L}{\partial z} + \frac{1}{v_g} \frac{\partial a_L}{\partial t} = - \frac{1}{2} (q_A a_L + q_B a_R).
\]

(75)

The amplitude of the pulse is diminished by absorption and by scattering into the opposite direction and is supplemented by scattering from the counter-propagating pulse.

The pulse propagation of two counter-propagating pulses colliding with each other in a saturable absorber is numerically calculated. Both pulses and the loss \( q(z, t) \) of the absorber are shown in Fig. 12 at different moments in time. In this figure, we have chosen the length of the absorber comparable to the physical length of the pulse. If the length of the absorber is short compared to the physical length of the pulse, the nonuniform effects of the grating at the edges of the absorber disappear, and the grating is almost uniformly spread along the absorber.

As we can see in Fig. 12, an absorption grating is set up where the two pulses overlap. With increasing saturation, the average loss in the absorber and the amplitude of the grating eventually diminishes. The absorption drops below the zero axis in the last pictures of Fig. 12 because we have expanded the absorption \( q(z, t) \) only to second order in space and neglected the fourth-order terms in our calculations.

C. Pulse Shaping Effects

In this section, we compare the pulse shaping effects of non-colliding and colliding pulse saturable absorbers, see Fig. 13.

A pulse of secant hyperbolic shape (solid line in Fig. 13) is sent through a saturable absorber in three configurations. In each case \( Q_0 = 0.3, E_{\text{sat}} A = 1 \), and \( E_t = 0.5 \) for the UM and SCPM configuration, and \( E_t = 0.25 \) for the CPM configuration. The normalized input pulse shape \( P_{\text{in}} \) is given by the solid line in Fig. 13. In the SCPM configuration, the pulse collides with itself in the absorber, and the resulting outgoing pulse is given by the dash-dotted line. In the CPM configuration, we have two pulses with energy \( E_t = 0.25 \), entering the absorber in opposite direction. The result is given by the dashed line, where the two outgoing pulses have been added.

As we can see, the interaction of the SCPM and CPM configuration cleans up the front of the pulse just as effectively as the UM configuration while the peak intensity after
by the SCPM and CPM is almost identical to the pulse shaping by a UM absorber with a saturation energy that is 51 and 67%, respectively, of the saturation energy of the SCPM and CPM absorber. In other words, identical absorbers give higher pulse powers in colliding pulse configurations than in the UM configuration.

We can, of course, perform the same analysis for pulses that collide in the saturable amplifier. We then find that the pulse shaping by the SCPM and CPM amplifier is almost exactly the same as the pulse shaping by a UM amplifier with a saturation energy that is 59 and 69%, respectively, of the saturation energy of the SCPM and CPM amplifier, respectively. The effective gain of an amplifier with a low saturation energy is less than the gain of one with a higher saturation energy. In other words, identical amplifiers give higher pulse powers if pulse collisions in the amplifier are avoided.

D. Enhanced Mode-Locking Stability

In this section, we compare the stability of a number of configurations for colliding pulse mode-locking. In Fig. 16(a), we have displayed the simplest CPM configuration in a ring cavity. As we have discussed in the preceding section, we only want to have pulse collisions in the saturable absorber and not in the saturable amplifier. For the ring-cavity, this means that the amplifier must not be opposite to the absorber (as displayed in Fig. 16(a)).

By separation of the absorber and amplifier by one quarter of the cavity, the pulses never collide inside the amplifier while the interval between the pulse arrival time is made large and equal for both pulses. In Fig. 16(b), we have shown this configuration.

For Fabry–Pérot cavities, we examine the SCPM configuration where the absorber is at one mirror while the amplifier is in the middle of the cavity in order to avoid pulse collisions inside the amplifier (see Fig. 16(c)). The CPM configuration can be established by putting the absorber in the middle of the cavity and the laser at one mirror end (in order to reach maximum pulse arrival time). For this configuration, however, we can not avoid that the pulses collide with themselves inside the amplifier. This configuration is displayed in Fig. 16(d).

In Table I, we have shown the changes in saturation energies and pulse arrival times for the absorber \(T_{p}^{A}\) and amplifier \(T_{p}^{D}\). All values are normalized to the values of the UM configuration of Fig. 8. Also the resulting change in the stability parameter \(s\) is calculated.

Evaluating Table I, we see that the configuration of Fig. 16(a) behaves almost exactly as the UM configuration.
Fig. 17. Stability boundary (area between drawn lines of the same kind) as a function of normalized small signal gain and absorption, for different mode-locking configurations, while $\tau_A/\tau_R = 0.1$, $\tau_L/\tau_R = 5$, $D = 10$, $\alpha_A = 0$, and $\alpha_L = 0$. For the other system parameters, see Table I.

The new boundaries for the CPM configurations of Fig. 16(b). The new boundaries for the SCPM configuration (Fig. 16(c) and (d)), however, are not improved compared to the UM configuration. This is due to the fact that for the CPM configurations, only one pulse with half the energy of the pulse of the SCPM configuration saturates the amplifier. This leads to an effective saturation energy of the amplifier that is twice as high as the saturation energy of the amplifier of the SCPM configuration. Finally, we want to note that we have assumed that the Fabry–Pérot cavities have perfectly reflecting mirrors. If the absorber is at the end mirror (see Fig. 16(c)), the reflectivity of that mirror should be as high as possible because this lowers the effective saturation energy of the absorber and thus increases the stability parameter $s$. In practice, this can be achieved by using a high reflection coating for this mirror. In contrast to this, if the amplifier is near an end mirror (see Fig. 16(d)), no high reflection coating should be applied to the mirror because this lowers the effective saturation energy of the amplifier and thus the stability parameter $s$. In practice, the reflectivity of the mirror at the amplifier end (see Fig. 16(d)) is about 30%. This leads to a somewhat higher value of the stability parameter $s$ compared to the configuration of Fig. 16(c). So, in practical Fabry–Pérot laser structures, the CPM mode-locking configuration behaves somewhat better than the SCPM configuration.
VIII. CONCLUSION

In this paper, we have proposed an analytical model for passively mode-locked semiconductor laser structures. The pulse shape generated by the passively mode-locked laser structure is a chirped hyperbolic secant with a width and energy that is determined by other system parameters. In the absence of dispersion and if absorber length is not the pulse width limiting factor, the width of the pulse is determined by the system bandwidth and depends thus on the active layer material used. Dispersion causes the pulse width to increase, while self-phase modulation imposes a chirp on the pulse. Under certain conditions, this chirp can be compensated by the interaction of the dispersion and the self-phase modulation of the amplifier and absorber.

We have also considered the conditions under which mode-locking occurs. The most important rule for mode-locking, as described in the papers of Haus [20]–[22], [26], is that the amplifier saturation energy must be higher than the absorber saturation energy. Furthermore, the saturable amplifier must not fully recover between two pulses, while the saturable absorber must recover completely or at least more than the amplifier between two pulses.

The effects of pulse collisions have also been investigated. Pulses, colliding in either a saturable amplifier or absorber, cause a grating due to the interaction of two counter-propagating electric fields. As a consequence of this grating, the saturable amplifier and absorber are more easily saturated compared to the situation where no grating is present. We can profit from colliding pulses, by designing a mode-locked system in such a way that pulse collisions only occur in the saturable absorber, while pulse collisions in the saturable amplifier remain absent.

We have also shown that various cavity configurations can be covered by the present model through adjusting the saturation energies of the saturable amplifier and absorber. Regarding the performance of the mode-locked system, a bidirectional ring cavity with a saturable amplifier and absorber separated by one quarter of the cavity length performs better than Fabry-Pérot cavities. A second inexperience with Fabry-Pérot cavities is that the pulse repetition time, determined by the cavity length, depends on the cleaving position of the end facets of the laser. This can be a rather inaccurate process. In contrast to this, the pulse repetition time for a ring laser, embedded in the production mask, can not be influenced by cleaving the output facets. Another disadvantage of Fabry-Pérot cavities is that for good performance of the CPM configuration, the absorber has to be exactly at a submultiple of the cavity length, which is not easily achieved. For a ring cavity, this problem does not occur. If one chose for a Fabry-Pérot cavity, one may use the CPM or SCPM configuration. If cleaving of the end facets can not be done accurately enough, it is recommended to choose for the SCPM configuration. Otherwise, the CPM configuration should be chosen as, in practical lasers, this configuration performs somewhat better than the SCPM configuration.

We can now formulate a number of design rules and recommendations for passively mode-locked lasers. First of all, a ring configuration gives the largest range of stable mode-locking operation. The system must be designed in such a way that pulse collisions occur in the saturable absorber and not in the saturable amplifier. In a ring cavity, this can be done by separating the amplifier and absorber by one fourth of the cavity. This increases the stability parameter \( s \) by a factor of three. The stability parameter \( s \) is also determined by the ratio \( \sigma_L/\sigma_A \). Stability can be increased by changing this ratio to a higher value. In practice, this can be achieved by designing a larger cross-section for the amplifier than for the absorber. Stability is also influenced by the ratio \( \sigma_L/\sigma_A \). The value of this ratio depends on the material used in the mode-locked laser. In order to increase stability, the value of the ratio \( \sigma_L/\sigma_A \) must be minimized. As this ratio is somewhat lower for quantum-well active layers compared to bulk active layers [33], passively mode-locked lasers constructed from quantum-well active layers might perform better than lasers made of bulk material. Further, we note that the bandwidth of the system must be maximized in order to achieve pulse widths as short as possible, while cavity dispersion must be minimized. Both parameters depend on the material used in the laser. Influencing these parameters, the pulse width and the effect of self-phase modulation of amplifier and absorber can be minimized, so that ultra-short unchirped pulses are obtained from the mode-locking system. A final advise in designing mode-locked lasers is to remember that the total absorption per unit length of an absorber can be about 10–20 times higher than the gain per unit length of an amplifier. In order to compensate all the losses in the mode-locking system and to set the laser above threshold, the length of an amplifier section must be chosen 10–20 times larger as the length of the absorber.

REFERENCES


This paper is part of the work carried out for the M.S. thesis at Eindhoven University of Technology in the Division of Optical Telecommunications. The project was an initiative of Philips Optoelectronics Centre, Eindhoven. After the M.S. degree, he continued working with the Division of Optical Telecommunications in Eindhoven on the subject of crosstalk in multichannel optical network interconnects. Since September 1995, he has been working toward the Ph.D. degree at California Institute of Technology, Pasadena. The subject of the research is mode-locking in semiconductor lasers.

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