The Transient Evolution of AM Mode-Locking a TEA CO₂ Laser

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Abstract—The evolution of the pulse in an AM mode-locked TEA CO₂ laser has been investigated. The experiments have been performed by injecting the mode-locked pulses in a high-pressure slave oscillator at various time intervals after the initiation of the mode-lock process.

This technique allows the measurements of the pulse widths independent on the pulse energies. A numerical solution of a dynamic model for the mode-locking process accurately predicts the transient evolution. It is shown that the build-up time to reach steady state can be, depending on the modulation depth, considerably larger than the duration of the pulse train.

I. INTRODUCTION

A commonly used method to generate nanosecond or sub-nanosecond pulses in a CO₂ laser system is active mode locking with an acoustooptic modulator. In principle, the pulse duration is limited by the bandwidth of the active medium. A theoretical analysis of this mode-locking process under stationary conditions has been worked out in the past [1]. It can be calculated with this theory that for 1 atmosphere CO₂ and 40 MHz modulation frequency, the pulse width is about 1 ns, and that the pulse duration decreases with pressure p as $p^{-1/2}$. These short pulses at pressures of 1 atmosphere and higher in CO₂ mixtures can only be generated during pulsed discharges. However, when dealing with UV preionized TEA CO₂ discharges with relatively short gain periods, there may be insufficient time for the mode-locking process to reach quasi steady state. Moreover, the generated pulses are unstable in shape and amplitude. Although these instabilities can be eliminated by applying an intracavity low-pressure section [2] or by injecting monochromatic CW radiation [3], there still remains the observation that the duration of each pulse is longer than what can be expected from the steady-state theory.

This discrepancy increases with pressure because for shorter pulses, the mode-locking process takes more time. Furthermore, the gain period after the discharge is shorter for higher pressures.

The purpose of the present contribution is to study the transient behavior of the mode-locking process in CO₂ systems and to investigate the time period for reaching steady state.

II. EXPERIMENTAL PROCEDURE

Our experimental technique is essentially injection mode locking [4]. The schematic representation is shown in Fig. 1. The mode-locking process to be studied occurs in a cavity

Fig. 1. Experimental configuration.

containing a single discharge TEA CO₂ laser [5], a low-pressure CW CO₂ discharge operating below threshold, and a modulator. The TEA CO₂ discharge has a volume of $20 \times 1 \times 1 \text{ cm}^3$ sealed with ZnSe Brewster windows and operating with a 1:1:3 = CO₂:N₂:He mixture at 1 atmosphere. The modulator is an acoustooptic germanium crystal with Brewster angles and driven by a 40 MHz RF power oscillator. The modulation depth as a function of RF power was measured by observing the response on a smooth, 200 ns long CO₂ laser pulse from a single mode, hybrid TEA CO₂ laser.

The low-pressure section is a 55 cm long sealed-off tube filled with a 1:2:5 = CO₂:N₂:He mixture at 20 torr; the bore is 15 mm. The pulse train coming from the AM mode-locked master oscillator is injected in the slave oscillator through the germanium outcoupling mirror. The discharge of this laser is a UV preionized multiatmosphere TE CO₂ system operating with a 1:1:10 = CO₂:N₂:He mixture at 5.5 atmospheres with a length of 30 cm [6]. We used this high pressure to increase the bandwidth of the slave oscillator in order to minimize the effect of broadening on the injected pulses by the dispersive gain medium [7]. Because we injected the full train of pulses in the slave oscillator, the cavity lengths of the two oscillators were made equal. An optimal matching was achieved by tuning the cavities for minimum pulse width from the slave oscillator. The amplification of an injected pulse by the slave oscillator is sufficiently large to avoid distortion of the selected pulse by the following injected pulses. To avoid additional pulses in the slave oscillator that originate from reflections on the outcoupling mirror of the master oscillator, we used for the slave oscillator a 2.5 m radius curved outcoupling mirror with a flat surface on the other side. Reflected pulses from this outcoupler are strongly diverged and practically blocked by an iris in front of the beam splitter (see Fig. 1). The slave oscillator produces a train of pulses which all, because of the large bandwidth of the active medium, have the same width equal to the injected one. The energies of the detectable pulses in this train do not depend

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on the energy of the injected one so that the widths of all pulses in the mode-locked TEA laser, independent on the pulse energies, can be observed. The pulses are detected with two photon drag detectors and a Tektronix 7912 AD transient digitizer. The response time of the detection system is estimated to be approximately 700 ps. The delay between the two discharges can be varied electronically. We measured the delay by observing the light coming from the last spark gap of the Marx generators of the two discharges. These light signals are transported by two optical fibers to photodiodes. In Fig. 2 we display some examples of measurements of parts of the pulse train coming from the slave oscillator. Generally, we observed broad pulses at a short delay time and small pulses at large delays between the discharges. As expected, the pulse widths depend on the modulation depth and whether the low pressure section in the master cavity was switched on or off.

In Figs. 3–5 we show the results of measurements of the width of pulses from the slave oscillator as a function of the delay between the two discharges. Figs. 3 and 4 show the measurements for the modulation depth of $\delta_m = 0.03$ and 0.15, respectively, with the low-pressure section switched on. Fig. 5 shows the results in the case where the low-pressure section is switched off and $\delta_m = 0.15$.

From these measurements, it is seen that with the low-pressure section active, the pulse evolves from a large value of the pulse width to a steady state in about 2 μ s. In the case that the low-pressure section was switched off, we could not measure the pulse width for low values of the delay with acceptable accuracy because in this area, the observed pulse envelopes were not stable, i.e., pulses with substructures and double pulses were often observed. The pulses observed at a large delay, $>2 \mu$ s, on the other hand, were found to be very stable. With the low-pressure section on, we observed longer pulses than without this section. For larger delays, say after 3 μ s, the pulses broaden again. This is because the pulses in the tail of the train are too weak to saturate the low-pressure gain. The pulses are then broadened by the low-pressure section.

In Fig. 4 it looks like a fast evolution to a width of about 3 ns in less than 1 μ s followed by a slower evolution to 1 ns in about 2 μ s. This can be explained as follows. In the beginning, the oscillations from the gain of the continuous section together with the weak gain of the pulsed section are just above threshold and are forced by the strong modulation to produce pulses which are mainly determined by the continuous section with small bandwidth. Later on, the gain of the pulsed discharge dominates and further compression occurs. For a smaller modulation depth (see Fig. 3), the pulse-forming mechanism is slower and there is not enough time for quasi-stationary pulse forming on the continuous discharge section.

To check whether we can interpret the observations as accurate measurements of the width of pulses in the train of the 1 atmosphere TEA CO₂ master oscillator, we also performed with a second photon drag detector direct measurements on the pulses in the train of the master oscillator. As mentioned before, the weak pulses during the build-up time and also in the tail of the train are not detectable. From the master oscil-

Fig. 2. Examples of observed pulses for different values of the delay between the master and slave discharge.

Fig. 3. Width of pulses coming from the slave oscillator as a function of the delay between the master and slave discharge. Full line = numerical solution; broken line = analytical solution [10].

lator, we only observe the strong pulses during the pulse width of the envelope. In Fig. 6 the widths of about 40 successive pulses have been displayed. The oscilloscope was triggered by the first detectable pulse in the train. In Fig. 7 we show the envelopes of the pulse trains for three cases shown in Figs. 3–5, respectively. Here a 20 MHz filter was used to suppress the ripple coming from the individual pulses in the train.

When we compare the results of Figs. 6 and 7 to those of Figs. 3–5, we see that the horizontal axis of Figs. 3–5 have to be shifted to the left in order to increase the delays so that the two measurements are consistent. An explanation for this is the fact that the observed delays come from the time differ-

Fig. 4. As Fig. 3, but with higher modulation depth.

Fig. 5. As Fig. 4, but with inactive low-pressure section.

ence with the last spark gap of the slave oscillator, whereas the gain of the slave to amplify the injected pulse appears somewhat later. The difference can be estimated as only a few tenths of a microsecond. We must add this small time delay to the delay observed by means of the light signals of the Marx generator gaps.

III. THEORETICAL MODEL

In this section, the dynamic laser processes and the pulse widths of a mode-locked TEA CO₂ laser are described numerically. We start from a set of rate equations based on a simplified four-level system as previously done by Gilbert *et al.* [8]. It is assumed that the rotational distribution is always ther-

Fig. 6. Direct measurements of the width of successive pulses from the AM mode-locked master oscillator.

Fig. 7. Measurements of the envelope of the pulse train coming from the AM mode-locked master oscillator.

malized because of the fast rotational relaxation. Thus, although the laser pulses may originate from a single rotational transition, the full vibrational inversion feeds the laser pulses.

For the densities of the lower laser level N_a , the upper laser level N_b , and the vibrational excited nitrogen N_c , one may write the following rate equations:

$$\frac{dN_a}{dt} = \rho\sigma c(N_b - N_a) + \gamma_{ba}N_b - \gamma_{ao}N_a + W_a \quad (1)$$

$$\begin{aligned} \frac{dN_b}{dt} = & -\rho\sigma c(N_b - N_a) - \gamma_{ba}N_b - \gamma_{bc}N_b \\ & + \gamma_{cb}N_c + W_b \end{aligned} \quad (2)$$

$$\frac{dN_c}{dt} = -\gamma_{cb}N_c + \gamma_{bc}N_b + W_c \quad (3)$$

where ρ is the photon density, $\rho\sigma c$ is the rate constant for stimulated emission, γ_{ij} is the exchange rate from level i to j , and W_i is the pumping source of level i . The various collisional rates [9] in the above equations are

$$\begin{aligned}\gamma_{ao} &= (10x + 4.6y + 0.56z) \mu s^{-1} \\ \gamma_{ba} &= (0.27x + 0.08y + 0.06z) \mu s^{-1} \\ \gamma_{bc} &= 14.5y \mu s^{-1}, \quad \gamma_{cb} = 13.2x \mu s^{-1}\end{aligned}$$

where x , y , and z are, respectively, the CO₂, N₂, and He partial pressures in atmospheres. According to Gilbert's model, the pumping rates can be approximated as

$$W_a = W_b = 0.4W_c = W_o t \exp(-t/\tau_p) \quad (4)$$

where τ_p represents the pumping pulse duration parameter and W_o is chosen to reproduce the specified small-signal gain of 3 percent/cm in a 1:1:3 = CO₂:N₂:He mixture. The coupling of the population to the field is given by

$$g_{TEA} = \sigma L' [N_b(t) - N_a(t)] \quad (5)$$

where σ is the cross section for stimulated emission and L' is the length of the active medium.

In the following, we calculate the evolution of an arbitrary pulse starting from noise or of a very weak signal from the continuous low-pressure discharge within a time window given by the roundtrip time $T = 2L/c$. It is assumed, and this can be verified, that the variations of the interacting quantities are very small during such a time interval. The changes of the densities during this time interval are therefore calculated according to (1), (2), and (3) for constant photon density and constant inversion densities. The used photon density in (1) and (2) is an average value which is related to the electric field by

$$\rho = \frac{c}{L} \int_{-\infty}^{\infty} |E|^2 dt \quad (6)$$

where the field is defined in such a way that $|E|^2$ represents the photon density. The change of the electric field during one roundtrip on its turn is calculated for constant inversion density.

The electric field $E_i(t)$ after the i th round trip is related to $E_{i-1}(t)$ by

$$E_i(t) = rT_i(t) \cdot F^{-1}[G_{TEA}(f) \cdot G_{CW}(f) \cdot F\{E_{i-1}(t)\}] \quad (7)$$

where $T_i(t)$ is the time-dependent transmission function of the acoustooptic modulator, r is the amplitude reflection coefficient including all other linear resonator losses, $G_{TEA}(f)$ and $G_{CW}(f)$ are the frequency-dependent gain of the TEA section and the low-pressure section, respectively, and F is the Fourier transform operator.

For the acoustooptic modulator, the periodic losses with frequency f_m can be expressed by

$$T_i(t) = \exp[-2\delta_m \sin^2\{-2\pi f_m(t - \delta\tau(i-1))\}] \quad (8)$$

where $\delta\tau = 1/2f_m - T$ is the detuning. The active media of both the TEA and the low-pressure section are assumed to be Lorentzian, i.e.,

$$\begin{aligned}G_{TEA}(f) &= \exp\left(\frac{g_{TEA}(i)}{1 + 2j(f - f_o)/\Delta f_{TEA}}\right) \\ G_{CW}(f) &= \exp\left(\frac{g_{CW}(i)}{1 + 2j(f - f_o)/\Delta f_{CW}}\right)\end{aligned} \quad (9)$$

where g_{TEA} and g_{CW} are the center frequency gain of the TEA and the low-pressure section, respectively.

The bandwidths of the TEA and low-pressure section are, respectively, Δf_{TEA} and Δf_{CW} ; f_o is the center frequency of the transition.

The maximum gain of both media saturates with laser intensity. For the continuous low-pressure section, it can be expressed as

$$g_{CW}(i) = \frac{g_{CW}^o(i)}{1 + I(i)/I_s} \quad (10)$$

where g_{CW}^o is the small-signal gain and I_s is the saturation photon density. The parameters used throughout the calculation are listed in Table I.

The results of the calculations are shown in Figs. 3–5 by the full lines. In the case of the CW section active, we start with a single mode. We see in Fig. 4 the discontinuity of the decreasing pulse width which we explained in the previous section. The model also predicts the increase in pulse width after the pulsed gain has disappeared.

In the absence of the continuous section, we start from spontaneous noise. It turns out that the results are practically independent on the shape and size of this noise signal.

It is seen that the evolution of the pulse width reaches a minimum which can be considered as the steady-state width. This steady state is, as expected, independent on the start pulse.

In the cases where the gain of the low-pressure section has been saturated by the high intensity inside the resonator or when the low-pressure section is switched off, the analytical solution of the pulse-forming process as introduced by Siegman and Kuizenga [10] can be applied:

$$\tau_p(t) = \tau_{p0} / [\tanh(t/t_0)]^{1/2} \quad (11)$$

where

$$\tau_{p0} = \frac{(2^{1/2} \ln 2)^{1/2}}{\pi} \left(\frac{g}{\delta_m}\right)^{1/4} \left(\frac{1}{f_m \Delta f_{TEA}}\right)^{1/2}$$

and

$$t_0 = \frac{\Delta f_{TEA}}{8\sqrt{2}g^{1/2}f_m^2\delta_m^{1/2}}.$$

This solution can be deduced from (7) by assuming Gaussian-shaped pulses and by approaching the gain function and the transfer function of the modulator also as Gaussian.

In Figs. 3–5 we plotted with the broken lines this analytical solution as well by starting the calculation from the moment the low-pressure gain saturates. We used the parameters listed in Table I. The gain has been chosen equal to the steady-state gain, i.e., $g = 1/2 \ln 1/R$. The small deviations from the numerical solution can be attributed mainly to gain switching during the pulse evolution.

TABLE I
PARAMETERS USED FOR THE CALCULATION ON THE AM TEA CO₂ LASER

Δf_{CW}	0.15 GHz
g_{CW}^0	0.25 (0.4 percent $\cdot$ cm ⁻¹)
I_s	5 W $\cdot$ cm ⁻²
Δf_{TEA}	5 GHz
$\tau = 2L/c$	12.5 ns
c	30 cm $\cdot$ ns ⁻¹
f_{mod}	40 MHz
$R = r^2$	0.6
τ_p	50 ns
W_0	1.4 10^{14} cm ⁻³ $\cdot$ ns ⁻²
σ	5.8 10^{-20} cm ²
CO ₂ :N ₂ :He	1:1:3 (1 atm. total pressure)

IV. DISCUSSION

It is shown that the build-up time for reaching short pulses by mode locking a TEA CO₂ laser can be longer than the pulse train.

For a modulation depth of $\delta = 0.15$ in a 1 atmosphere 1:1:3 = CO₂:N₂:He system, the formation time is about 2 μ s, whereas the envelope of the pulse train is between 1.25 and 1.75 μ s after the discharge. The build-up time decreases with increasing modulation depth. So it depends on the modulation depth whether the steady-state conditions will be reached during the pulse train; one has to be careful by applying a simple analytic expression for determining the pulse width in a TEA CO₂ system. This will be especially the case for multi-atmosphere systems because for a higher gas pressure with larger line width, the number of roundtrips to reach full compression of the pulse will be larger, and also because the duration of the inversion after the pulsed discharge in a multiatmosphere system will be shorter due to the faster relaxation phenomena at higher pressure. Further, it is found that a numerical solution of the mode-locking process using a dynamic model for the interactions accurately predicts the transient evolution of mode locking a TEA CO₂ laser.

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