Mode locking of a continuous wave Nd:glass laser pumped by a multistripe diode laser

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The performance of a continuous wave actively mode-locked Nd:phosphate glass laser longitudinally pumped by a multistripe diode laser is described. The laser operates at 1.054 μm; the pump threshold and the slope efficiency are found to be 120 mW and 11%, respectively. The shortest pulse duration is 7 ps, which appears to be approximately twice as short as predicted by the theory of amplitude modulation mode locking. We explain the improved performance by additional frequency modulation due to the nonlinear index of the active material.

Short pulse generation with solid-state lasers traditionally has been realized by active mode locking of cw Nd:YAG lasers and by passive mode locking of pulsed Nd:glass lasers. The advantage of shorter pulses from glass lasers was significantly impaired by their low repetition rate and poor stability. To combine the advantages of cw active mode locking with the large bandwidth of a glass laser became possible by using an argon laser (514 nm) instead of lamps as a pump source.[1–3] The appearance of GaAlAs diode lasers offered the possibility to further improve the performance of laser-pumped glass lasers.

Diode lasers are compact, efficient, and reliable sources of coherent radiation. Their use as a pump source provides significant advantages over argon lasers: mainly higher efficiency, reduced thermal problems, and improved stability. Recently, diode-pumped mode-locked neodymium lasers have been built using both crystals (YAG,[4] YLF[5]) and glass[6] as a host material. As expected, the glass laser generates the shortest pulses. However, because of the low pump power delivered by the single stripe diode laser the pulse energy was too low to measure the pulse duration accurately.[6] Stimulated by this work we started to develop a multistripe diode-pumped Nd:glass laser with tolerable threshold.

The pump threshold of an end-pumped laser can be greatly reduced by tight focusing of the cavity mode and pump beam in the active medium. The smallest spot sizes achievable depend on the shortest active length possible and thereby on the absorption length of the laser material at the pump wavelength. In a glass laser concentration quenching and/or thermal damage set a limit to the dopant concentration, the quantity controlling absorption.[7] Careful design preceded the experiments to find the most suitable glass host and optimum values for dopant concentration and spot sizes of pump and cavity modes in the active material. Besides threshold and slope efficiency thermal effects such as lensing, birefringence, and damage were considered.

As a result of these investigations we found a phosphate glass (Schott LG 760) with 8% Nd^{3+} concentration to be the optimum choice for pump powers < 500 mW and $\lambda_{pump} = 0.8$ μm.[7] The absorption coefficient is $\alpha_{pump} = 18\ cm^{-1}$, allowing an absorption efficiency of $\eta_a = 0.9$ in an active medium as short as $L_a = 1.2$ mm. The requirements of a long resonator and small beam waist were best satisfied by a folded cavity commonly used in dye lasers.[8]

Figure 1 shows a schematic diagram of our laser cavity. The 1-mm-thick plate of LG 760 Nd:glass with an 8% doping level was inserted at Brewster's angle into the resonator. The laser medium was placed at the beam waist of the cavity mode. To facilitate heat removal, the glass plate was sandwiched between two copper plates contacted to a heat sink. The astigmatism of the oblique-incidence mirror $M2$ was compensated for by that of the active material at $\vartheta = 9°$.[8] The use of the plane mirror $M3$ for folding the long arm of the cavity made a compact setup possible.

The pump source was a 500 mW ten-stripe GaAlAs diode array (SDL 2432) temperature tuned by a Peltier cooler for operation at 798 nm. The emitted radiation was collected using a compound lens of 8 mm focal length and a numerical aperture of 0.5 (Melles Griot 06GLC002). This lens focused the beam through the mirror $M1$ onto the active medium in the plane *perpendicular* to the array. An additional cylindrical lens of focal length 19 mm was employed to provide tight focusing along the direction *parallel* to the array. This has been shown to be the most efficient way to obtain a small spot and good overlap with the cavity mode in the active medium.[9]

The average spot size of the pump beam in the active medium was calculated to be $(\overline{w_p^2})^{1/2} \simeq 45\,\mu$m.[10] Due to the limited spatial coherence of a multistripe diode laser this value is about twice the minimum value for a diffraction-limited Gaussian beam at this wavelength. To achieve high slope efficiency $\overline{w_p^2}$ might not exceed the cavity mode spot size $\overline{w_l^2}$ considerably. We met this requirement by appropriate choice of the radius of curvature of the folding mirror

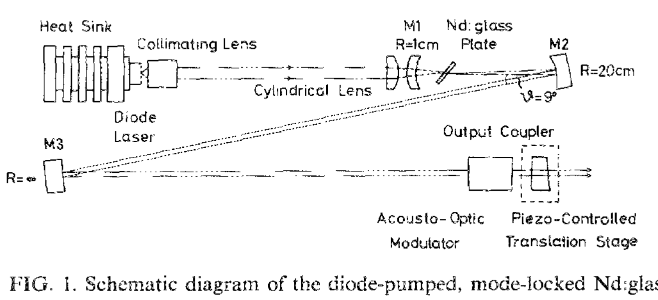

FIG. 1. Schematic diagram of the diode-pumped, mode-locked Nd:glass laser.

$M2$. $R_2 = 20$ cm and a cavity length of $L_c \approx 2$ m yielded $(\bar{w}_l^2)^{1/2} \approx 40\,\mu$m. This value refers to the case when the short leg of the cavity is adjusted to the middle of the stability range, which amounted to 5 mm in our case. The long cavity length reduces the pulse repetition rate and thereby facilitates switching single pulses into a regenerative amplifier.

Active mode locking of the thin plate Nd:glass laser was accomplished by an acousto-optic standing-wave Bragg modulator (IntraAction ML-385J) inserted close to the output coupler into the cavity. Both windows of the modulator are AR coated at 1.05 μm. The modulator was driven at about $\nu_m = 38$ MHz with 5 W of rf power. The round trip amplitude transmission of a standing-wave Bragg modulator is given by

$$a(t) = \cos^2(\theta_m \sin \omega_m t), \tag{1}$$

where $\omega_m = 2\pi\nu_m$ and the modulation depth θ_m was measured to be 0.8 for a dissipated rf power of 5 W. The ultrastable frequency synthesizer generating the rf signal was tuned near the cavity resonance and locked to avoid frequency drift. Fine adjustment was performed by positioning the output coupler with a piezo translator. With a 2.5% output coupler the pump threshold of the mode-locked laser described above was measured to be ≈ 120 mW. Although the output was not exactly a linear function of the input power[7] we defined an effective slope efficiency, which was found to be $\approx 11\%$.

The mode-locked laser pulses were monitored by a picosecond diode and sampling oscilloscope. Simultaneously, the pulse widths were measured by the standard collinear intensity autocorrelation method. The autocorrelator scanned a range of 120 ps with a repetition rate of 3 Hz. The pulse width resolution was <100 fs. The spectra of the pulses were investigated by a 1 m double grating spectrometer with a resolution of 0.1 Å. The shortest pulses were generated at pumping levels so low ($<2P_{\rm th}$) that stable pulses just developed. The autocorrelation trace at an average output power of ≈ 5 mW is shown in Fig. 2. The full width at half maximum (FWHM) of the trace is 10 ps. Assuming a Gaussian pulse shape we obtained the pulse duration $\tau_p = (10/\sqrt{2})\,{\rm ps} \approx 7$ ps. The laser bandwidth (FWHM) was measured to be 3.8 Å, and revealed a modulated structure.

The measured data yielded a time-bandwidth product of 0.71, a value which is 1.6 times that for a bandwidth-limited Gaussian pulse (0.44). The coherence-spike-free autocorrelation trace and the modulated spectrum of the laser suggest that the pulses are phase modulated. It should be noted that stable short pulse generation tolerates only a very small amount of cavity detuning: 1 μm change in the cavity length implies a very noisy autocorrelation trace and a considerably increased pulse width. The noise sitting on the autocorrelation trace in Fig. 2 is mainly attributed to small variations of the cavity length in the absence of a feedback loop providing stabilization.

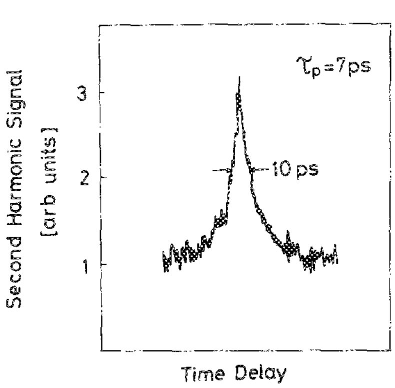

FIG. 2. Autocorrelation trace of the laser pulses at an average output power of approximately 5 mW.

Let us compare these results with predictions of the self-consistent field theory of active mode locking developed by Kuizenga and Siegman.[11] The round trip amplitude transmission experienced by a short pulse passing through the modulator can be written without loss of generality as

$$a(t') = \exp\{-[\delta_0 + \delta_1(\omega_m t') + \delta_2(\omega_m t')^2 + \cdots]\}, \tag{2}$$

where t' is a restarted time so that the center of the pulse is at $t' = 0$ and the constants are to be evaluated for particular modulation. They depend on the type of modulator, the modulator depth, and the phase of the modulation $\psi = \omega_m t_0$ when the pulse passes through the modulator. For an amplitude modulator all the constants are real. Assuming a Gaussian pulse $E_0 \exp(-\gamma t'^2) \exp(-i\omega_l t')$ to circulate in the cavity and a cavity length tuned exactly to the driving frequency (i.e., $\delta_0 = \delta_1 = 0$) the Kuizenga theory yields

$$\gamma \approx \pi^2 \nu_m \Delta\nu \sqrt{\delta_2/g}, \tag{3}$$

describing the self-consistent solution. $\omega_l = 2\pi\nu_l$, where ν_l is the center frequency of the laser radiation, $\Delta\nu$ is the gain bandwidth (FWHM), and $g \approx 1/2 \ln(1/R_{\rm eff})$ is the round trip amplitude gain, where $R_{\rm eff}$ includes all cavity losses. Generally, $\gamma = \alpha + i\beta$ may be complex, where $\alpha = 2\ln 2/\tau_p^2$ and β is the chirp parameter describing a linear frequency sweep during the pulse. Owing to the real δ_2 for an amplitude modulator Eq. (3) predicts a real γ and thereby bandwidth-limited pulse generation from an AM mode-locked laser.

Expanding (1) into power series at $\psi = 0$ yields $\delta_2 = \delta_a = \theta_m^2$. In our case $\delta_2 = \delta_a = 0.64$, $\Delta\nu = 6$ THz, and $g = 0.05$, substituting these numbers into (3) we obtained $\tau_p = 13$ ps. However, instead of bandwidth-limited pulses of 13 ps duration the laser generates phase-modulated pulses of 7 ps duration. Although the theory rigorously applies only to homogeneously broadened lasers it should be a good approximation also for a phosphate glass laser, in which the width of a so-called effective homogeneous line including overlapping Stark transitions is found to be comparable to that of the fluorescence line including inhomogeneous broadening as well.[7]

Besides, group velocity dispersion (GVD) and self-phase modulation (SPM) by intracavity elements are possible sources of the discrepancy. Whereas GVD is expected to be small in a picosecond laser, noticeable SPM occurs in the glass plate, where the cavity mode is tightly focused. The round trip phase shift due to SPM is given by

$$\Delta\phi_{\rm SPM}(t') = (2\pi/\lambda_l) n_2 I(t'), \tag{4}$$

where λ_l is the laser wavelength, n_2 represents the nonlinear index of the glass host, and $I(t')$ is the pulse intensity along the cavity axis in the active material. Approximating the Gaussian pulse by a parabolic pulse shape the round trip amplitude transmission is modified due to SPM by the factor $\exp[-i\delta_{\rm SPM}(\omega_m t')^2]$, where the time-independent phase shift is ignored and the modulation index by SPM is given by

$$\delta_{\rm SPM} = \frac{4\ln 2}{\pi}\left(\frac{L_a}{\lambda_l}\right)(n_2 I_0)\left(\frac{1}{\tau_p \nu_m}\right)^2, \quad (5)$$

where $I_0 = I(t'=0)$. Thus, the resulting modulation index δ_2 becomes complex: $\delta_2 = \delta_a + i\delta_{\rm SPM}$. Although the round trip phase shift due to SPM $4\pi(L_a/\lambda_l)(n_2 I_0)$ is very small, $\delta_{\rm SPM}$ may take considerably large values owing to the factor $(1/\tau_p \nu_m)^2$. In other words, the large "curvature" of the phase change makes it possible, in spite of its small magnitude, to significantly influence pulse shaping in the actively mode-locked laser.

With $n_2 = 2.8\times 10^{-10}$ cm^2/MW, $I_0 = 15$ MW/cm^2, and $\tau_p = 13$ ps, $\delta_{\rm SPM} = 8.3$ is obtained, i.e., the phase modulation dominates in the pulse-shaping process under these conditions. Substituting $\delta_2 = 0.64 + 8.3i$ into (3) the steady-state pulse parameters can be calculated by successive approximation because of the dependence of $\delta_{\rm SPM}$ on τ_p. After a few iterations we obtained $\tau_p \simeq 5$ ps and a time-bandwidth product of 0.62. As the quadratic phase shift approximation somewhat overestimates the extent of pulse shaping by SPM the agreement with experimental results can be considered satisfactory.

One would expect further pulse shortening at higher output levels due to an increased value of $\delta_{\rm SPM}$. However, numerical calculations of Haus and Silberberg[12] have shown that this kind of pulse shaping becomes unstable when the pulse shortening with respect to the pulse width without SPM exceeds a factor of approximately 2. It is quite understandable considering the fact that SPM acts as an ideal frequency modulator only during the center part of the pulse. Although no signs of instabilities were observed in our laser at increased pump powers, significant pulse broadening occurred resulting in pulse widths of 12–20 ps for 10–40 mW of output power ($P_{\rm pump} \leqslant 400$ mW).

Cavity (or modulator) detuning effects can also be treated by the self-consistent field theory. When the cavity length (or modulation frequency) is detuned from its ideal value, the pulse passes through the modulator at some phase angle ψ different from zero. As a consequence, the pulses get longer and the output power decreases. For a low-loss oscillator the theory predicts the latter effect to be more severe. The relative frequency shift reducing the output power by a factor of 2 is calculated for a standing-wave Bragg modulator to be

$$\left(\frac{\Delta\nu_m}{\nu_m}\right)_{\rm crit} \simeq \frac{8}{\pi}\frac{\nu_m}{\Delta\nu} g, \quad (6)$$

where the expressions $\delta_0 = \theta_m^2 \sin^2\psi$, $\delta_1 = \theta_m^2 \sin 2\psi$, and $\delta_2 = \theta_m^2 \cos 2\psi$ were used. Equation (6) holds as long as $g/\theta_m^2 \ll 1$. Interestingly, the relative critical detuning is directly proportional to the modulator frequency and inversely proportional to the gain bandwidth. As a result $(\Delta\nu_m/\nu_m)_{\rm crit}$ is less than 10^{-6} in our case leading to a cavity length tolerance shorter than 1 μm. Apparently, active stabilization of L_c (or ν_m) by a feedback loop is needed to maintain long-term stability of mode locking. A Fabry–Perot étalon reducing the gain bandwidth provides another means of increasing stability at the expense of broadened pulses. For the same reason, lasers such as Nd:YAG and Nd:YLF are less sensitive to detuning.[13]

In summary, we developed a cw mode-locked Nd:glass laser pumped by a high-power multistripe diode laser. The complete utilization of the weak concentration quenching in phosphate glass allowed us to achieve a low threshold in spite of the poor spatial coherence of the pump beam. The shortest pulses generated by the laser are about twice as short as that predicted by the theory of active mode locking. We found self-phase modulation in the active medium to be a possible source of additional pulse shortening. This effect might have played an important role also in recent short pulse generation experiments by active mode locking.[3,5,6,14] In comparison with diode-pumped Nd:YAG and Nd:YLF lasers the Nd:glass laser generates considerably shorter pulses with nearly the same efficiency. We conclude that Nd:glass is a promising active material for future diode-pumped solid-state laser systems.

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