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Increasing the modulation instability threshold in a direct-diode-pumped Kerr-lens mode-locked Ti:Sa oscillator

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At high pump powers and intracavity laser power of Kerr-lens mode-locked lasers, modulation instability inevitably arises, limiting pulse energy and duration; it has been studied for Diode Pumped Solid-State pumping but not for direct-diode pumping. We implement a direct-diode-pumped Ti:Sa oscillator using two GaN diodes at 455/465 nm. Varying the pump-lens focal length F suppresses modulation instability. For $F = 80$ mm, we obtain stable 12 fs pulses (~ 100 nm bandwidth) with average power up to 300 mW at 8 W of pump; further pump increase leads to noise and breakup into sub-pulses. For $F = 60$ mm, the spectrum broadens to 110–120 nm with 450 mW at 11 W. For $F = 50$ mm, the spectrum is 127–140 nm (10 fs pulse duration) and the power is 500 mW at 11 W, with MI absent. This result provides a practical route to improving the efficiency and stability of direct-diode-pumped Kerr-lens mode-locked Ti:Sa oscillators.

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High-power few-cycle femtosecond (fs) pulses are in strong demand across many areas of science and technology [1]. The most popular source of such pulses is the Ti:Sa oscillator due to its broad gain bandwidth, which enables robust generation of sub-10-fs pulses [1].

Today commercial Ti:Sa oscillators pumped by diode-pumped solid-state (DPSS) lasers are very expensive. A natural alternative is direct diode pumping. However, the demonstrated average output powers are typically relatively low [2].

The reasons for the low output power of such pulses include the onset of modulation instability (MI). For Kerr-lens mode-locked (KLM) oscillators, MI is a key operational boundary. When certain thresholds are exceeded (pump power, dispersion, etc.), the oscillator can enter an unstable regime, where the fs pulse splits into sub-pulses, bifurcations appear and noise grows. Such regimes limit the achievable average power and pulse energy, may lead to a loss of mode locking. In case of diode pumping with large divergence and poor beam quality, changing the focusing of diode beam in the crystal can strongly affect the Kerr soft aperture, and consequently the output parameters of the fs radiation and the onset of MI.

In this work, we demonstrate that changing the focal length F of the pump lens modifies the spatial overlap between the diode pump and the intracavity laser mode. This affects the effective soft Kerr aperture. Here we employ the wave equation ($\xi = z - v_{gr}t, t' = t$, v_{gr} is the pulse group velocity) to analyze the dependence of the MI threshold and laser output power on F:

$$\frac{\partial E}{\partial t'} = (D_r + iD_i)\frac{\partial^2 E}{\partial \xi^2} + \frac{g}{2}E.$$

Here D_i and D_r describe the frequency dispersion of the refractive index and the gain bandwidth. The net gain parameter g includes both the active-medium gain and diffraction losses caused by the Kerr-lens aperture [3]:

$$g = \frac{GN^e}{1 + (G/\gamma)\,\overline{|E|}^2} + \frac{p}{1 + |E|^2/I_d}\,|E|^2.$$

Where G is the Einstein coefficient for stimulated emission, γ is the inversion relaxation time, N^e is the inversion in the absence of radiation, $\overline{|E|}^2$ is the spatially averaged intracavity intensity, $p = \sigma_1/\sigma_0$ is the ratio of linear diffraction losses to total linear losses, and I_d is the characteristic intensity at which the diffraction losses decrease by a factor of two.

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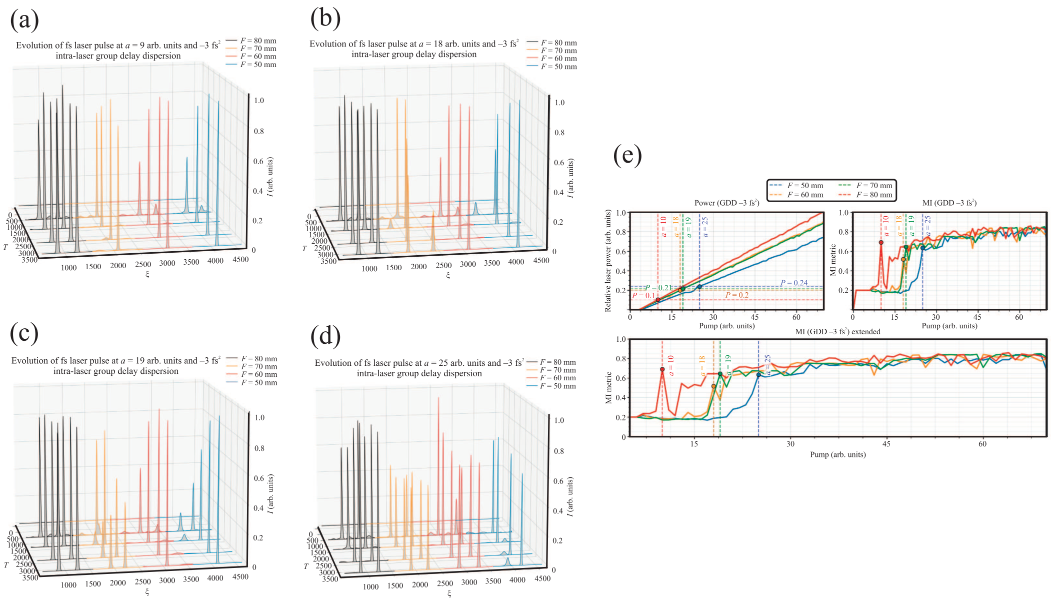

Fig. 1. (Color online) Calculated evolution of the fs pulse for GDD $= -3$ fs^2. Panels (a)–(d) show the evolution of $I(\xi, T)$ at pump values $a = 9$, 18, 19, and 25 for different focal lengths. Panel (e) shows the relative output power $P_L(a)$ and the MI metric MI(a)

After normalization: $t = t'\sigma_0, \xi = \xi'\sqrt{\frac{\sigma_0}{D_r}}$ and $\theta = \frac{D_i}{D_r}, b = \frac{I_d}{LI_g}\sqrt{\frac{D_r}{\sigma_0}}, q = -\frac{2\omega_0 n_2 I_d l}{n_0\sigma_0 L}$ with $\delta\sigma = \frac{p}{1+|E|^2}|E|^2, E = \frac{E'}{\sqrt{I_d}}$ the equation takes form [3]:

$$\frac{\partial E}{\partial t} = (1 + i\theta)\frac{\partial^2 E}{\partial \xi^2} + \frac{1}{2}\left(\frac{1+a}{1 + b\int |E|^2 d\xi} - 1 + \delta\sigma + iq|E|^2\right)E.$$

The model accounts for how pump focusing affects the parameters b, p, and q. To quantify the pump–laser overlap, we use the factor $\eta(\xi, F)$. The normalized pump $a(\xi, F)$ is written as $a_0\eta(\xi, F)$:

$$\eta(\xi, F) = \frac{\int I_P(x, y, \xi, F)\, I_L(x, y, \xi)\, dx\, dy}{\sqrt{\left(\int I_P^2(x, y, \xi, F)\, dx\, dy\right)\left(\int I_L^2(x, y, \xi)\, dx\, dy\right)}}.$$

Here, $I_P(x, y, \xi, F)$ and $I_L(x, y, \xi)$ are the spatial distributions of the pump and intracavity laser intensity.

The evolution equation for $E(\xi, t)$ was solved using a symmetric Strang split-step Fourier scheme. We quantify the MI by detecting the local peaks of side subpulses and calculating the fraction of pulse energy located outside the main central peak. If $0 \leq \mathrm{MI} \leq 0.1$, the pulse remains single and MI is absent. For $0.1 \leq \mathrm{MI} \leq 0.3$, weak asymmetry and small side structures appear in $I(\xi, T)$. For $0.3 \leq \mathrm{MI} \leq 0.6$, the side sub-pulses become pronounced. For $\mathrm{MI} \geq 0.6$, the main pulse evolves into a multi-pulse regime.

The temporal evolution of the $I(\xi, T)$ was obtained for the intracavity group-delay dispersion, GDD $= -3$ fs^2, and for different focal lengths of the pump lens. In Figure 1, the coordinate ξ denotes the retarded time, T is the number of cavity round-trips, and the vertical axis shows the relative intracavity intensity.

The onset of MI strongly depends on the pump focusing. The earliest transition occurs for $F = 80$ mm, where the MI rises at $a \approx 10$. For $F = 60$ and 70 mm, the transition is shifted to $a \approx 18$ and $a \approx 19$, whereas for $F = 50$ mm it occurs only at $a \approx 25$. Thus tighter pump focusing suppresses MI and extends the single-pulse operation range, allowing a larger increase of the pump. This trend is directly visible in panels (a)–(d): at low pump all cases remain nearly single-pulse, then pulse breakup first appears for $F = 80$ mm, while the $F = 50$ mm case stays stable over the widest range.

The pump source consists of two GaN laser diodes operating at 455/465 nm. Their radiation is combined by a dichroic mirror. Each channel is shaped by a cylindrical telescope. The pump waist inside the Ti:Sa crystal is formed by a spherical lens L with different focal lengths $F = 80, 60$ and 50 mm. For each F, the cavity was optimized for the shortest pulse duration without noise signatures in the pulse train and splitting. The laser output was sent through an external prism compressor to compensate the GDD introduced by the cavity optics. With a pump lens of $F = 80$ mm, stable operation at the minimum pulse duration (12 fs) was limited to 300 mW of average power at 8 W of pump, a further pump increase led to noise and pulse splitting. In contrast, for $F = 50$ mm, the instability threshold was shifted to higher pump power, which enabled stable generation of 10 fs pulses with an average output power of 500 mW at 11 W of pump.

Thus, controlling the pump beam geometry by variation of the pump lens L provides a practical way to suppress MI, increase the output power, and broaden the spectrum in diode-pumped KLM Ti:Sa oscillators.

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