A route to extreme peak power and energy scaling in the mid-IR chirped-pulse oscillator-amplifier laser systems

ALEXANDER RUDENKOV,1,* VLADIMIR L. KALASHNIKOV,1 EVGENI SOROKIN,2,3 MAKSIM DEMESH,1 AND IRINA T. SOROKINA1,3

1Department of Physics, Norwegian University of Science and Technology, N-7491 Trondheim, Norway
2Photonics Institute, Vienna University of Technology, 1040 Vienna, Austria
3ATLA Lasers AS, Richard Birkelands vei 2B, 7034 Trondheim, Norway
* alexander.rudenkov@ntnu.no

Abstract: The paper introduces a new route toward the ultrafast laser peak power and energy scaling in a hybrid mid-IR chirped pulse oscillator-amplifier (CPO-CPA) system, without sacrificing neither the pulse duration nor energy. The method is based on using a CPO as a seed source allowing the beneficial implementation of a dissipative soliton (DS) energy scaling approach, coupled with a universal CPA technique in a simple single-path oscillator-amplifier system. The key is avoiding a destructive nonlinearity in the final stages of an amplifier and compressor elements by using a chirped high-fidelity pulse from CPO as a seed for a single-pass CPA. Our main intention is to realize this approach in a Cr2+:ZnS-based CPO as a source of energy-scalable DSs with well-controllable phase characteristics for a single-pass Cr2+:ZnS amplifier. A qualitative comparison of experimental and theoretical results provides a road map for the development and energy scaling of the hybrid CPO-CPA laser systems, without compromising pulse duration. The suggested technique opens up a route towards extremely intense ultra-short pulses and frequency combs, in particular, in the mid-IR spectral range from 1 to 20 μm.

© 2022 Optica Publishing Group under the terms of the Optica Publishing Group Open Access Publishing Agreement

1. Introduction

In the last decades, many front-end achievements in modern science have been made possible by the progress in femtosecond laser technology [1, 2]. These advances have provided breakthroughs, for example, in attosecond and relativistic physics [3, 4] and ultrafast spectroscopy [5], nano-photonics [6], particle beam acceleration [7], material processing [8], and many other fields of science and technology. Such ultra-intense ultrashort light pulses opened the doors for the discovery of many new fundamental physical phenomena, allowing us to look inside the atom, deliver the power to objects in outer space, produce ultraprecise measurements, make new discoveries in astrophysics, quantum electrodynamics, open the doors to relativistic phenomena and even question some of the fundamental constants. The current status of femtosecond technology is PW-level peak power lasers. That peak power corresponds to irradiances in the range of $10^{22}$–$10^{23}$ W/cm$^2$ [9], with the highest irradiances currently amounting to $10^{23}$ W/cm$^2$ [10]. Ion motion becomes relativistic [11] at intensities above $10^{25}$ W cm$^{-2}$. Such intensity regimes have been recently demonstrated with very-large-scale lasers like the HERCULES laser at the University of Michigan in the USA or ELI (the Extreme Light Infrastructure project) across Europe. Above $10^{30}$ W cm$^{-2}$, nonlinear quantum electrodynamics becomes within reach. However, it is also important that the next generation high-intensity lasers become more compact and user-friendly, thus, opening a wider avenue towards numerous scientific and industrial applications. The extreme peak powers in combination with ultrashort pulse widths, open the way for high-field physics using merely university lab tabletop lasers [12]. In the future compact high-intensity lasers will also make
numerous societal applications a reality, with proton therapy in medicine, a clean (neutron- and nuclear waste-less) nuclear energy production with particularly compact femtosecond laser drivers, nuclear waste transmutation, just to name a few.

To date, the solid-state mode-locked lasers have allowed the generation of femtosecond high-peak power pulses directly from an oscillator with high (>MHz) repetition rates [13]. In contrast to a classical chirped-pulse amplification scheme [14], such a high pulse repetition rate promises an extreme signal rate improvement factor of $10^3$-$10^4$. However, the main trouble in further energy scaling in parallel with peak power for such systems is still the destructive contribution of nonlinearities. The alternative approach is provided by the integration of a chirped-pulse amplifier (CPA) and the laser into a single oscillator, chirped pulse oscillator (CPO) [15, 16], which utilizes the remarkable property of a dissipative soliton (DS) to combine high stability with the energy scalability [17].

Simultaneously, the deep insight into the mechanisms of ultrashort pulse generation allows the development of a “smart laser” in which a transparent algorithm of optimization and, ultimately, self-optimization of a laser device would be realized [18,19]. In particular, the concepts of a dissipative soliton resonance [20] and a master diagram [21] provide us with the tools for such an optimization aimed at a DS energy scaling (e.g., see [22,23]).

In our work, we theoretically and experimentally consider the ways of ultrafast laser power and energy scaling in a hybrid mid-IR CPO-CPA system. Our main intention is to use a CPA based on a Cr$^{2+}$:ZnS-based mode-locked laser [24,25] as a source of energy-scalable DS with well-controllable phase characteristics for a single-pass Cr$^{2+}$:ZnS polycrystalline ceramic amplifier. Using a CPO as a seed source allowed us to implement a DS energy scaling approach and incorporate such laser to the universal concept of CPA, beneficially without using a pulse stretcher and reducing the amplification stage that lowers the system complexity and avoids a destructive nonlinearity contribution. DS energy scaling can be performed by changing several laser parameters such as pump power, mode area, and pulse repetition frequency (cavity period). This issue will be discussed in more detail later in this work.

Cr$^{2+}$:ZnS crystals are characterized by the high stimulated emission cross sections and one of the broadest among all existing lasers gain [26], which makes it possible to obtain sufficient amplification of few optical cycle pulses even in a single-pass amplifier configuration. A combination of more than one amplification stages can provide comparatively high pulse energy still at full repetition rate, while the nonlinearity in the amplifier media is controlled by scaling the mode area, thus operating the amplifier close to linear regime. What is particularly important is that using a CPO provides an effective pre-amplification without degradation of amplification efficiency due to operation in the vicinity of the maximum fidelity range, where chirp is almost spectrally-independent.

In this paper, we propose a novel technique to provide hybrid CPO-CPA pulse energy scaling while not having to sacrifice the pulse duration (maximizing peak power) and preserving the output spectrum’s high-fidelity compression. In the following two sections, we describe the first steps in this direction, both experimentally and theoretically, demonstrating the viability of the proposed concept.

2. Experimental setup

For the experimental part of our work, we developed a setup consisting of a seed laser, single pass amplifier and a prism compressor (Figure 1). For pulse parameter control we used the home-made interferometric autocorrelator, capable to measure pulse duration up to 50 ps, APE waveScan Extended IR spectrometer (800-2600 nm) and usual sensors like power meters and photodetectors.
We start the design of the seed laser from a 70 MHz soliton-like pulse Kerr-lens mode-locked laser operating in anomalous dispersion regime, with 2.6 mm thick Cr$^{2+}$:ZnS Brewster-cut single crystal as an active medium. Beam radius in the active element was about 40 µm. As a pump source we employ a 5 W Er-fiber laser operating at 1610 nm. Intracavity dispersion balance was maintained by using specially designed chirped mirrors which compensated simultaneously the second (group delay dispersion, GDD) and the third (third-order dispersion, TOD) orders of dispersion created by gain medium and the YAG wedges, used for fine tuning of the GDD. Highly reflective mirrors and an output coupler were also designed for broadband operation and flat dispersion profile with low absolute values. Such approach allowed obtaining soliton-like 5 nJ pulses with about 36 fs duration, assuming a sech$^2$ shape. Spectra and autocorrelation traces are shown in Figs. 2 and 3. However, direct amplification of such pulses inevitably causes strong nonlinear effects in the amplifier medium [30].

We therefore tune the dispersion into normal regime by introducing additional HR mirrors with small amount of positive dispersion distributed over a broad spectral range.

![Fig. 1. Experimental setup. Chirped pulse oscillator (CPO), single pass amplifier (AMP), prism compressor (CMP), concave mirror (CM), chirped mirror (ChM), wedges for fine dispersion compensation (ChW), output coupler (OC), dichroic mirrors (DM), erbium doped fiber laser (EDFL), fused silica prisms (FS), highly reflective mirrors (HR).](image1)

![Fig. 2. Soliton-like pulse laser spectra.](image2)

![Fig. 3. Interferometric autocorrelation traces of soliton-like pulses.](image3)
Simultaneously the laser cavity was modified for operation at 12.345 MHz pulse repetition frequency to scale the pulse energy. Intracavity dispersion profile of the CPO regime is shown in Figure 4.

![Fig. 4. CPO intracavity dispersion profile.](image)

Single pass amplifier used a Cr$^{2+}$:ZnS polycrystalline gain element with 13.5x8x2.7 mm$^3$ dimensions and Cr$^{2+}$ concentration of 2.3·10$^{18}$ cm$^{-3}$ (IPG Photonics). Beam radius in the amplifier gain element was about 60 µm. After the amplifier we install a prism compressor made of very dry fused silica material. Total dispersion profile of the amplifier and compressor elements is shown in Figure 5.

![Fig. 5. Combined dispersion of the amplifier and the prism compressor.](image)

3. Experimental results

During the experiment we first characterize the laser parameters in continuous wave (CW) regime of operation. Then we switch to the mode-locking (ML) regime. ML operation was provided by a Kerr-lens mechanism only, that is better suitable for energy scaling purposes due to higher laser induced damage threshold compared to using absorbing material modulators such as SESAMs or graphene mirrors [27-29].

Output power characteristics of the laser in both CW and ML regimes are shown in Figure 6.

![Fig. 6. Output parameters of the oscillator in different operation modes.](image)
Maximum average output power in the ML regime is slightly higher than in CW – 273 mW versus 259 mW with optical-to-optical efficiencies of 12.2% and 11.6% respectively, which is due to the soft-aperture based self-amplitude modulation. Maximum pulse energy at the CPO output was about 22 nJ. In addition to power/energy characteristics of the CPO we also measured pulse duration (blue curve) and spectral width (green curve in Figure 7).

As one can see from the data, the CPO pulse duration (uncompressed, blue curve on Figure 7) – spectral width (green curve at Figure 7) behavior can be subdivided into two zones. First zone with pulse energies from 5 to 15 nJ, where pulse duration shortens and spectral width increases, and the second zone with pulse energies from 15 to 22 nJ, where spectral width stays nearly the same, but pulse duration approximately linearly increases. The first zone behavior has some similarities to the Schrodinger soliton, but the second zone is more interesting for our power scaling purposes. In this case the DS spectral width is clamped in a certain range, defined by the CPO design parameters, so the pulse must increase its duration in order to accommodate the energy growth into the CPO area theorem. Such behavior is called “Dissipative soliton resonance” (DSR) in the literature [19] and is achieved by increasing the pulse phase modulation (chirp) thus making the pulse longer. In Fig. 7, the red line shows the amplified pulse duration after the compressor, which was optimized for the highest energy. As expected, the autocorrelation trace durations indicate overcompensation at lower energies, where the chirp is lower.

The autocorrelation traces and pulse spectra for three characteristic points (A – before DS resonance, B – transition period and C – DS resonance, on the Figure 7) are shown in Figs. 8 and 9.
A short comment on the spectral width estimation technique. DSs at the CPO output have specific shape of the spectrum namely, a spectrum with sharply cut edges. In addition, the incompletely compensated TOD in our laser leads to a slope in the upper part of the spectrum, which introduces additional ambiguity in the case of using the width at half maximum criterion. We therefore used the following width definition: we calculated the first derivative of the spectral curve and estimated the distance (spectrum width) between the peaks of the derivative (see Figure 10).

Input pulse energy was 20 nJ resulting in output pulse energy of 62 nJ under amplifier incident pump power of about 12.7 W (energy gain ratio 3.1). The compressor transmission of about 80% reduces the output pulse energy to 50 nJ that correspond to 278 kW power.

It should be noted, that using the CPO seed effectively turns this setup to a chirped-pulse amplifier scheme with strongly reduced nonlinearity. The spectra of the amplified and seed pulses are nearly identical (as opposed to soliton amplification [30]) indicating negligible spectral phase distortion and suggesting further energy scalability of this scheme.
4. DS scalability and compressibility

The advantage of CPO as a source for further pulse amplification is its energy scalability provided by a significant DS chirp [16]. It is well-known [31,32] that a DS developing in the normal GDD regime ($\beta_2 > 0$) can be described as a soliton-like solution of the complex cubic-quintic non-linear Ginzburg-Landau equation [33]:

$$\frac{\partial a(z,t)}{\partial z} = \left\{ i \left( \beta_2 \frac{\partial^2}{\partial t^2} - \gamma |a(z,t)|^2 \right) + \left[ -\sigma + \alpha \frac{\partial^2}{\partial t^2} + (\kappa - \zeta |a(z,t)|^2) |a(z,t)|^2 \right] a(z,t) \right\}. \quad (1)$$

Here, $\beta_2$ is a GDD coefficient, $\alpha$ is a squared inverse bandwidth of a spectral filter (e.g., a gain bandwidth), $\gamma$ is a self-phase modulation (SPM) coefficient, $\sigma$ is a saturated net-loss coefficient, $\kappa$, and $\zeta$ are the coefficients describing a saturable nonlinear gain due to soft-aperture mode-locking [34].

Adiabatic theory of a strongly-chirped DS [21,31] allowed a unified description of DS by mapping a system parametric space to the so-called master diagram (Figure 11) spanned by two dimensionless coordinates: $c = 2\gamma \beta_2 \kappa$ and $E^* = E \times \left( \kappa \sqrt{\zeta / \beta_2 \gamma} \right)$, where the first parameter $c$ describes a relative contribution of spectral dissipation, GDD, self-phase and self-amplitude modulations. The second coordinate $E^*$ is a dimensionless DS energy $E$ [31]:

$$E = \frac{6\nu \zeta}{\gamma} \frac{E}{\beta_2} \frac{1}{\zeta} \arctan \frac{\Delta}{\xi}. \quad (2)$$

The complex spectral amplitude $\varepsilon(\omega)$ in the absence of higher-order GDD has the following form [35]:

$$\varepsilon(\omega) = \frac{6\pi \gamma \exp \left\{ \frac{\alpha^2 (\omega^2 - \omega_{0}^2)}{2 \xi} \right\} \text{Heaviside}(\Delta^2 - \omega^2)}{\sqrt{\Delta^2 + \omega^2}}. \quad (3)$$

Fig. 11. Master diagram representing DS stability threshold (black curve), spectral half-width $\Delta^*$ (black dotted curve) and DS $T^*$ width (black dashed curve) in dependence on the normalized energy $E^*$. Magenta curve confines a region of energy scalable DS with finger-like spectra. Normalizations are: $E^* = E \times \left( \kappa \sqrt{\zeta / \beta_2 \gamma} \right)$, $T^* = T \times \left( \kappa \sqrt{\gamma \zeta \beta_2} \right)$, and $\Delta^* = \Delta \times \sqrt{\beta_2 \zeta \gamma}$. The parameter $\Delta = \sqrt{P_0 / \beta_2}$, where $P_0$ is a DS peak power, defines the spectrum width, i.e., the cut-off frequency. Formally, such a cut-off can be explained by a resonance between DS and background radiation with the wave-numbers $\gamma P_0$ and $\beta_2 \Delta^2$, respectively [24,36].
A Lorentz-like central profile ("finger") of the spectrum is defined by a "width" $\Xi = \sqrt{\gamma \left( 1 + c \right) / \zeta - \frac{5P_0}{3}} / \beta_2$. The characteristic spectral power profile from (3) is

$$ p(\omega) = |\varepsilon(\omega)|^2 = \frac{6\pi \gamma \text{Heaviside}(\Delta^2 - \omega^2)}{\zeta \left( \Xi^2 + \omega^2 \right)} $$

(4)

and the corresponding spectral chirp is

$$ Q(\omega) = \frac{1}{2} \frac{d^2 \phi(\omega)}{d\omega^2} = \frac{3\gamma^2}{2\beta_2 \kappa' \xi} \left( \frac{\Delta^2 - \omega^2}{\Xi^2 + \omega^2} \right)^{-1}. $$

(5)

Fig. 11 (a) demonstrates that the DS energy is scalable along the curve (black) $\sigma = 0$ defining the DS stability. In a same way as for a Schrödinger soliton, the negativity of net-gain $\sigma$ is a necessary condition of DS stability against the background noise. In a steady-state, this parameter can be described as $\sigma \approx \delta (E/E_{cw} - 1)$, where $E_{cw}$ is an energy of CW-radiation, i.e., a mean power multiplied by a cavity period in the CW-regime, and $\delta = l^2 / g_0$ ($l$ is a net-loss coefficient, $g_0 = l (E_{cw}/E_0 + 1)$ is a small-signal gain, and $E_0$ is a gain saturation energy) [21]. The implicit expression for the stability threshold $\sigma = 0$ presented by the black solid curve in Fig. 11 (a) follows from Eq. (2), when $\Delta$ and $\Xi$ are taken for a limiting value $P_0 = (3/2)(1 - c/2)^{-1}$.

of all). From Eq. (2), such scalability (i.e., $\lim_{\Xi \to 0} E = \infty$) exists within the interval of $c = (5/6) \sqrt{1 - 4b} - 1/6$, where $b = \zeta \sigma / \kappa \in [0, 1/4]$. This energy-scaling is provided by the DS broadening (see the dashed black curve in Fig. 11) with an asymptotically constant spectral half-width $\Delta \to \frac{3}{2} \left( 1 - \zeta \right) \gamma / \beta_2 \zeta$ (compare Figs. 7 and 11 (dashed red curve)), $\Xi \to \frac{\left( \frac{9c}{4} - \frac{3}{2} \right) \gamma / \beta_2 \zeta$, and peak power $P_0 \to \frac{3}{2} \left( 1 - \zeta \right) \zeta^{-1}$. Thus, we would summarize the DSR conditions obtained from the adiabatic theory:

$$ E \to \infty \Leftrightarrow \begin{cases} P_0 \to \frac{3}{2} \left( 1 - \zeta \right) \zeta^{-1}, \\
\Delta \to \frac{3}{2} \left( 1 - \zeta \right) \gamma / \beta_2 \zeta, \\
\Xi \to \sqrt{\frac{\left( \frac{9c}{4} - \frac{3}{2} \right) \gamma / \beta_2 \zeta}, \\
c = \frac{5}{6} \sqrt{1 - 4b} - \frac{1}{6}, \\
b \in \left[ 0, \frac{1}{4} \right]. \end{cases} $$

(6)

Figs. 7, 11, and Eqs. (6) demonstrate the hallmarks of a transition to DSR: 1) change of pulse squeezing to its broadening with growing energy, 2) constant spectral width $\Delta$, and 3) appearance of a visible spike at the spectrum center. The finger-like spectrum ($\Xi < \Delta$) (see Eq. (5) and [31,37]) is clearly visible in Figure 9, $\text{SP}_C$. The first hallmark is crucial for DS energy harvesting: its peak power is fixed (Eq. (6)) but the energy grows by the pulse stretching $\propto 1/\Xi$ (Figs. 7, 11) due to a chirp scaling [21]: $\psi \approx (\gamma \Delta / 4\pi) E$.

The latter factor allows energy re-distribution inside a DS preserving its integrity. Using the approximating aproach for a DS profile (Figure 12)
\[ a(t) = a_0 \text{sech} \left( \frac{t}{T} \right) \exp[i(\phi(x) + \Omega t + \theta \tanh \left( \frac{t}{T} \right) + \psi \log \text{sech} \left( \frac{t}{T} \right) + \chi \log^2 \text{sech} \left( \frac{t}{T} \right)] \] (7)

(\phi\) is a phase, \(\Omega\) is a frequency shift from a gain band centrum, \(\psi\) is a chirp, \(\theta\) and \(\chi\) are the phase distortions caused by TOD and FOD, respectively, allows defining the energy flow \(j(t')\) inside the soliton \(t' = t/T, T\) is a DS width) [38,39]:

\[
j(t') \equiv i \frac{1}{2} (a \partial_{t'} a^* - a^* \partial_{t'} a) = \frac{a_0^2}{2} \text{sech}^4(t') \left[ \left( i \psi + 2 \chi \log \text{sech} \left( \frac{t'}{T} \right) \right) \sinh(2t') - 2\theta \right] \] (8)

Figure 13 demonstrates such a stabilizing flow, which increases with \(\psi\).

Fig. 12. Dimensionless spectral profiles \(p(\omega')\) for \(\psi = 5, \theta = \chi = \Omega = 0\) (black curve), \(\psi = 5, \theta = 2, \chi = \Omega = 0\) (red curve) and \(\psi = 5, \theta = 2, \chi = 1, \Omega = 0\) in Eq. (7).

Fig. 13. Dimensionless energy flux \(j(t')\) in dependence on the chirp \(\psi, \theta = \chi = 0\) in Eq. (7).

The next practically important phenomenon demonstrated by the master diagram is the transition to an energy-scalable regime at \(c = \sqrt{1 - 16b/3}\) when \(\Delta = \Xi\) (magenta curve in Figure 11). That corresponds to a maximally flat \(Q(\omega)\) around the DS spectral center (Eq. (5)), which means the maximal DS compressibility (or maximum compression fidelity) by a subsequent chirp compensation in a compressor [40]. Generally speaking, two magenta curves in Fig. 11 represent the domain of maximal DS energy scalability. The corresponding DS parameters are:
Eqs. (9) provide us with a guide to approximate estimation of unmeasurable parameters of the self-amplitude modulation \( \kappa \) and \( \zeta \) in Eq. (1). Fig. 9, SPc demonstrates a transition to a finger-like spectrum and could be considered as a manifestation of crossing the maximal fidelity curve in Fig. 11. Then, the knowledge of experimental spectral width \( 2\Delta = 2.32 \times 10^{13} \) s (127 nm at 2.27 \( \mu \)m central wavelength), intracavity energy \( E = 90 \) nJ (for 14.6\% net-loss), SPM coefficient \( \gamma = 2.3 \) MW\(^{-1}\), squared inverse gain bandwidth \( \alpha \approx 10 \) fs\(^2\), and GDD coefficient \( \beta = 687 \) fs\(^2\) leads to \( \kappa \approx 2.16 \) MW\(^{-1}\) and \( \zeta \approx 2.4 \) MW\(^{-1}\) that is close to the theoretical estimations obtained from the theory of soft-aperture KLM [34].

To completely quantify the master diagram, we need to estimate the saturated net-gain \( \sigma \approx \delta(E/E_{cw} - 1) \). This parameter can be derived from the comparison of CW and mode-locking energies for different pump levels, as in Fig. 6. Fig. 14 shows the results of the calculations. Then, using the Eqs. (6) would allow fitting the experimental data in the master diagram.

![Fig. 14. Dependencies of the saturated net-gain \( \sigma \) (black curve) and \( \delta \) (blue curve) parameters on the intracavity CW-energy \( E_{cw} \) calculated from the experimental data. Dashed red curve is smoothed dependence for \( \sigma \).](image)

As Figure 12 demonstrates, the observed asymmetry of spectra (Figs. 9, 10) is a direct result of the TOD (Figure 4). This factor could be an obstacle to the energy scaling and needs elimination by a fine GDD control in the CPO. An additional factor affecting the DS spectrum is FOD causing the formation of an M-like spectral shape (Figure 12) [41].

5. Discussion

A qualitative comparison of experimental and theoretical results provides a road map for a hybrid CPO-CPA pulse energy scaling. As the master diagram (Figure 11) demonstrates, the operating parameters of a CPO have several distinct areas separated by the stability and
maximum fidelity curves. Excluding the area of instability, we can focus on the two regions on the master diagram divided by the maximal pulse fidelity curve.

The characteristic feature of this division is the behavior of the pulse spectrum, which changes its shape from a flat-top to a finger-like one when passing through the maximum fidelity curve in the direction of increasing pulse energy. Such behavior is clearly seen during our experiment - spectral shape changes from flat-top (Figure 9, SP A) to finger-like (Figure 9, SP C), which confirms the CPO operation near the maximum fidelity curve. Then, crossing a fidelity border accompanies a change of the DS width behavior with the energy growth - from decreasing to rising (Figs. 7, 11b) with the spectral width finally reaching the asymptotic. Such qualitative behavior of the sound and measurable CPO characteristics provide a path to CPO design and optimization.

These hallmarks follow the idea of DS energy scaling with maximum fidelity (compressibility) and using such extremely broad and smooth gain bandwidth media as Cr:ZnS without a significant spectral phase degradation in CPA. That could be a road to developing the high-power ultrafast laser system with ultrashort pulse durations working at high pulse repetition rates.

One must remark on the main obstacles to realizing the perfect DS power/energy harvesting in a CPO-CPA system. The first is a destructive contribution of higher-order dispersions (TOD, FOD, etc.). They result in the DS spectrum distortion and squeezing and may cause a strong destabilization, including chaotization and even DS fission [25, 41, 42]. Especially, the TOD is maximally destructive in this way that requires applying the elaborated dispersion compensation techniques.

There also exists the risk of possible fidelity degradation with advancing into the domain of perfect energy scalability, which needs an additional study. The issue could be illustrated by a close analogy between DS and turbulence [43]. The energy harvesting results from a soliton width growth $\Lambda_t \propto 1/\Xi$. The last parameter $\Lambda_t$ plays a role of a “long-range” correlation scale [44]. Simultaneously, the “cut-off” frequency $\Delta$ defines a “short-range” correlation scale $\Lambda_\Delta \propto 1/\Delta$. Under this angle of view, it is the coincidence of these scales $\Lambda_t = \Lambda_\Delta$ that provides maximal fidelity. However, the growing discrepancy of these correlation scales with the $\Xi$-decrease can lead to a DS “decoupling”, that is to the emergence of “soliton gas” at a “distance” of $\Lambda_\Delta$ [45,46] trapped by a collective “potential” with a characteristic width $\propto \Lambda_t$. Such partially coherent DS becomes more sensitive to perturbations, which are clearly visible in the spectrum (see Fig. 9, SP C and [47]).

Finally, from a thermodynamical point of view, the expression (4) is analogous to the Rayleigh-Jeans distribution [48,49] with a negative “chemical potential” $\mu = \Xi^2$ and a “temperature” $T = 6\pi^2 / \kappa$ [43]. The DS energy growth can be treated as a result of the particle number $N$ (mass) increase by analogy with the Bose-Einstein condensation from a “basin” [50]. The negativity of chemical potential preserves the entropy $S$ in such a process when a temperature is constant and the Gibbs free energy $dF = dU - TdS - \mu dN$ decreases, in agreement with the fundamental thermodynamic relation $dU = TdS + \mu dN$, where $U$ is a total energy of a system. However, $\mu$ tends to zero in parallel with such a process (Eq. (6) for $c = 2/3$). This “wave condensation” counteracts the further particle number growth in parallel with the $F$ minimization. By analogy with a Bose-Einstein condensate, it means that a further condensation (i.e., DS energy growth) becomes impossible [51], and DS tends to the internal decoupling (“fission”) due to $\Lambda_\Delta / \Lambda_t \rightarrow 0$. However, the detailed analysis of the transit to turbulence awaits further study.

### 6. Conclusion

We suggested and show feasibility to avoid the destructive nonlinearities in the final stages of amplifiers and compressor elements by using a chirped high-fidelity pulse from a CPO as a seed for a single-pass CPA. Such scheme eliminates the necessity of stretchers and expensive
and troublesome electrooptical instrumentation, being suitable for pulse generation and amplification systems operating at high pulse repetition rates. A qualitative comparison of experimental and theoretical results provides an algorithmically realizable road map for a hybrid CPO-CPA pulse energy scaling while not having to sacrifice the output pulse duration by preserving the output spectrum's high-fidelity compression. Experimentally, we demonstrate a Cr$^{2+}$:ZnS laser system that provides 180 fs pulses with 50 nJ pulse energy at 12.345 MHz pulse repetition frequency.

The proposed methodology promises to open up a way to compact table-top laser instruments that combine high repetition rate and high peak power in the mid-IR, enabling productivity increase in applications and novel scientific approaches, that are yet inaccessible in university laboratories.

**Funding.** The work is supported by the Norwegian Research Council projects #303347 (UNLOCK), #326503 (MIR).

**Disclosures.** The authors declare no conflicts of interest.

**References**

1. T. Brabec F. Krausz “Intense few-cycle laser fields: frontiers of nonlinear optics.” Rev Mod Phys. 72, 545–591 (2000).
2. Z. Chang, P. B Corkum, S. R. Leone. “Attosecond optics and technology: progress to date and future prospects” [Invited]. J Opt Soc Am B. 33, 1081–1097 (2016).
3. G. A. Mourou, T. Tajima, S. V. Bulanov. Rev. Mod. Phys. 78, 309–371 (2006).
4. F. Krausz, M. Ivanov. “Attosecond physics.” Rev Mod Phys. 81, 163–234 (2009).
5. T. Kobayashi. “Development of ultrashort pulse lasers for ultrafast spectroscopy.” Photonics. 5, 11 (2018).
6. P. Dombi, Z. Papa, J. Vogelsang, et al. “Strong-field nano-optics.” Rev Mod Phys. 92, 025003 (2020).
7. D. Guénot, D. Gustas, A. Vernier, et al. “Relativistic electron beams driven by kHz single-cycle light pulses.” Nat Photon. 11, 293–296 (2017).
8. A. Nejadmalayeri, P. Herman, J. Burghoff, et al., “Inscription of optical waveguides in crystalline silicon by mid-infrared femtosecond laser pulses” Opt. Lett. 30, 964-966 (2005).
9. Ch. Radier, O. Chalas, M. Charbonneau, et al. “10 PW peak power femtosecond laser pulses at ELI-NP”, High Power Laser Science and Engineering, 10, e21, 5 pages (2022). doi:10.1017/hpl.2022.11
10. J.W. Yoon, Y.G. Kim, I.W. Choi, J.H. Sung, H.W. Lee, S.K. Lee, and C.H. Nam, Optica 8, 630 (2021).
11. T. Esirkepov, M. Borgesi, S. Bulanov, G. Mourou, and T. Tajima, Phys. Rev. Lett. 92, 175003 (2004).
12. T. Südmeyer, S. Marchese, S. Hashimoto, et al. Femtosecond laser oscillators for high-field science. Nature Photon. 2, 599–604 (2008).
13. C. Baer, O. Heckl, C. Saraceno, et al. “Frontiers in passively mode-locked high-power thin disk laser oscillators” Opt. Express 20, 7054-7065 (2012).
14. D. Strickland, G. Mourou. “Compression of amplified chirped optical pulses.” Optics Commun. 55, 447–449 (1985).
15. A. Fernandez, T. Fuji, A. Poppe, A. Fürbach, F. Krausz, and A. Apolonski. “Chirped-pulse oscillators: a route to high-power femtosecond pulses without external amplification” Opt. Lett. 29, 1366-1368 (2004).
16. S. Naumov, A. Fernandez, R. Graf, P. Dombi, F. Krausz, and A. Apolonski. “Approaching the microjoule frontier with femtosecond laser oscillators.” New Journal of Physics 7, 216 (2005).
17. P. Grelu, N. Akhmediev. “Dissipative solitons for mode-locked lasers” Nature Photon 6, 84–92 (2012).
18. T. Baumeister, S.L. Brunton, and J.N. Kutz. “Deep learning and model predictive control for self-tuning mode-locked lasers.” JOSA B 35, 617-626 (2018).
19. R. I. Woodward, and E.J.R. Kelleher. "Towards ‘smart lasers’: self-optimisation of an ultrafast pulse source using a genetic algorithm." Scientific reports 6, 1-9 (2016).
20. W. Chang, A. Ankiewicz, J. M. Soto-Crespo, and N. Akhmediev. “Dissipative soliton resonances.” Physical Review A 78, 023830 (2008).
21. V. L. Kalashnikov, E. Podivilov, A. Chernykh, and A. Apolonski. “Chirped-pulse oscillators: theory and experiment.” Applied Physics B 83, 503–510 (2006).
22. Wei Lin, Wang Simin, Xu Shanhui, et al. "Analytical identification of soliton dynamics in normal-dispersion passively mode-locked fiber lasers: from dissipative soliton to dissipative soliton resonance." Optics Express 23, 14860-14875 (2015).
23. W. H. Renninger and F. W. Wise. “Fundamental Limits to Mode-Locked Lasers: Toward Terawatt Peak Powers.” IEEE Sel. Top. Quantum Electron. 21, 1100208 (2015).
24. E. Sorokin, N. Tolstik, and I.T. Sorokina, “1 Watt femtosecond mid-IR CrZnS laser,” Proc. SPIE 8599, 859916 (2013).
25. E. Sorokin, N. Tolstik, V. Kalashnikov, and I. Sorokina, "Chaotic chirped-pulse oscillators," Opt. Express 21, 29567-29577 (2013).
26. I. T. Sorokina, E. Sorokin, "Femtosecond Cr2+-based Lasers," IEEE J. Sel. Topics Quantum Electron. 21, 1-19 (2015). Doi: 10.1109/JSTQE.2014.2341589
27. N. Tolstik, I. T. Sorokina, E. Sorokin, "Graphene Mode-locked Cr:ZnS Chirped-pulse Oscillator," in Advanced Solid-State Lasers Congress, M. Ebrahim-Zadeh and I. Sorokina, eds. (Optical Society of America, 2013), p. MW1C.2. doi: 10.1364/MICS.2013.MW1C.2
28. N. Tolstik, A. Pospischil, E. Sorokin, I. T. Sorokina, "Graphene mode-locked Cr:ZnS chirped-pulse oscillator," Opt. Expr. 22, 7284-7289 (2014). doi: 10.1364/OE.22.007284
29. N. Tolstik, C. S. J. Lee, E. Sorokin, and I. T. Sorokina, "6 MHz Extended Cavity Cr:ZnS Chirped-pulse Oscillator," in Conference on Lasers and Electro-Optics, OSA Technical Digest (online) (Optical Society of America, 2018), paper SF1N.2.
30. S. Vasilyev, I. Moskalev, M. Mirov, S. Mirov, and V. Gapontsev, "Multi-Watt mid-IR femtosecond polycrystalline Cr2+:ZnS and Cr2+:ZnSe laser amplifiers with the spectrum spanning 2.0-2.6 µm," Opt. Expr. 24, 1616-1623 (2016).
31. E. Podivilov and V. L. Kalashnikov. “Heavily-chirped solitary pulses in the normal dispersion region: new solution of the cubic-quintic complex Ginzburg-Landau equation.” Journal of Experimental and Theoretical Physics Letters 82, 467–471 (2005).
32. W. H. Renninger, A. Chong, and F. W. Wise. “Dissipative solitons in normal-dispersion fiber lasers.” Phys. Rev. A 77, 023814 (2008).
33. J. D. Moores, “On the Ginzburg-Landau laser mode-locking model with fifth-order saturable absorber term,” Opt. Commun. 96, 65–70 (1993).
34. J. Hermann. “Theory of Kerr-lens mode locking: role of self-focusing and radially varying gain”. JOSA B 11, 498–512 (1994).
35. V. L. Kalashnikov. “Dissipative solitons: perturbations and chaos formation.”, pp. 199-206 in Ch. Skiafas and I. Dimotikalis, eds. Chaos Theory: Modeling, Simulation and Applications: Selected Papers from the 3rd Chaotic Modeling and Simulation Conference (CHAOS2010), Chania, Crete, Greece, 1-4 June 2010. World Scientific, 2011.
36. V. L. Kalashnikov and S. Sergeyev, “Dissipative Solitons in Fibre Lasers,” in Fiber Laser, Mukul Chandra Paul, (Ed.), pp. 165-210 (ISBN: 978-953-51-4615-5, InTechOpen, 2016).
37. V. L. Kalashnikov, E. Podivilov, A Chernykh, et al. "Approaching the microjoule frontier with femtosecond laser oscillators: theory and comparison with experiment." New J. Phys. 7, 217 (2005).
38. A. Maimistov. "Evolution of solitary waves which are approximately solutions of a nonlinear Schrödinger equation." Sov. J. Exp. Theor. Phys. 77, 727-731 (1993).
39. N. Akhmediev and A. Ankiewicz. Dissipative solitons: from optics to biology and medicine (Lecture notes in physics, vol. 751) (Heidelberg, Springer, 2008).
40. L. Zhu, A. J. Verhoef, K. G. Jespersen, V. L. Kalashnikov, L. Grüner, D. Lorenc, A. Baltuska, and A. Fernández. "Generation of high fidelity 62 fs, 7-nJ pulses at 1035 nm from a net normal-dispersion Yb-fiber amplifier with anomalous dispersion higher-order-mode fiber." Optics express 21, 16255-16262 (2013).
41. V. Kalashnikov, A. Fernández, and A. Apolonski. "High-order dispersion in chirped-pulse oscillators," Opt. Express 16, 4206-4216 (2008).
42. E. Sorokin, V. L. Kalashnikov, J. Mandon, et al. “Cr3+:YAG chirped-pulse oscillator.” New J. Phys. 10 083022 (2008).
43. V. L. Kalashnikov and E. Sorokin, "Self-organization, coherence and turbulence in laser optics." In Complexity in biological and physical systems, R. Lópeez-Ruiz, (Ed.), pp. 97-112 (ISBN: 978-1-78923-051-2, InTechOpen, 2018).
44. A. Picozzi, B. Barviau, B. Kibler, and S. Rica. “Thermalization of incoherent nonlinear waves,” Eur. Phys. J. Special Topics 173, 313-340 (2009).
45. A. I. D’yachenko, V. E. Zakharov, A. N. Pushkarev, et al. “Soliton turbulence in nonintegrable wave systems.” Soviet Physics – JETP 69, 1144-1147 (1990).
46. S. Smirnov, S. Koptsev, S. Kukarin, and A. Ivanenko. “Three key regimes of single pulse generation per round trip of all-normal-dispersion fiber lasers mode-locked with nonlinear polarization rotation.” Optics Express 20, 27447-27453 (2012).
47. V. L. Kalashnikov, A. Chernykh. “Spectral anomalies and stability of chirped-pulse oscillators.” Phys. Rev. A 75, 033820 (2007).
48. V.E. Zakharov, V.S. L’vov, G. Falkovich, Kolmogorov Spectra of Turbulence I (Springer, Berlin, 1992).
49. A. Picozzi. “Towards a nonequilibrium thermodynamic description of incoherent nonlinear optics.” Optics Express 15, 9063-9083 (2007).
50. V. L. Kalashnikov, S. Wabnitz. A "metaphorical nonlinear multimode fiber laser approach to weakly dissipative Bose-Einstein condensates." EPL 133, 34002 (2001).
51. G. Düring, A. Picozzi, S. Rica. “Breakdown of weak-turbulence and nonlinear wave condensation.” Nonlinear Phenomena 238, 1524 (2009).