Detuned Raman Amplification of Short Laser Pulses in Plasma

V. M. Malkin, G. Shvets, and N. J. Fisch
Department of Astrophysical Sciences, Princeton University, Princeton, New Jersey 08540
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The recently proposed scheme of so-called "fast compression" of laser pulses in plasma can increase peak laser intensities by 10^5 [Phys. Rev. Lett. **82**, 4448 (1999)]. The compression mechanism is the transient stimulated Raman backscattering, which outruns the fastest filamentation instabilities of the pumped pulse even at highly overcritical powers. This Letter proposes a novel nonlinear filtering effect that suppresses premature backscattering of the pump in a noisy plasma layer, while the desired amplification of a sufficiently intense seed persists with a high efficiency. The effect is of basic interest and also makes it robust to noise the simplest technologically fast compression scheme.

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The recently proposed superradiant [1] or fast compression [2] schemes can substantially increase the peak laser intensities, which are currently achievable through the chirped pulse amplification technique [3]. The peak laser powers can also be substantially increased, while the sizes and costs of ultrapowerful laser compressor amplifiers can decrease remarkably. Ultrahigh energy, short laser pulses might have a variety of applications [3,4].

The fast compression scheme [2] employs transient stimulated Raman backscattering in plasma for pumping a seed laser pulse to an extremely high intensity before filamentation instabilities develop. These instabilities grow very fast at high laser intensities. Yet, 100 PW/cm^2 power densities and kJ/cm^2 fluences of 1 μm wavelength radiation can be achieved through fast compression, making feasible compact devices for 100 PW kJ laser pulses. No material except plasma is capable of withstanding the huge electric fields of such laser pulses.

This Letter proposes a novel scheme for transporting a laser pump to a seed pulse precisely through the amplifying plasma layer. Although desirable technologically, this direct pump path is complicated by the same extreme efficiency of Raman backscattering that makes possible the fast compression. As the pump traverses the plasma layer towards the seed pulse, fast Raman backscattering of the pump by thermal Langmuir waves may lead to premature pump depletion. The problem is aggravated by the fact that the linear Raman backscattering instability of the pump (responsible for unwanted noise amplification) has a larger growth rate than its nonlinear counterpart (responsible for the useful amplification of the seed laser pulse). Nevertheless, through an interesting nonlinear filtering mechanism identified here, it appears to be possible to suppress the unwanted Raman backscattering of the pump by noise, while not suppressing the desirable seed pulse amplification.

The filtering effect occurs because, in the nonlinear regime, the pumped pulse duration decreases inversely proportional to the pulse amplitude. The stronger the pulse, the faster the pump depletion. The pulse frequency bandwidth increases with the pulse amplitude, so that the nonlinear instability, as it grows, can tolerate larger and larger external detuning from the backscattering resonance. On the other hand, the same bandwidth broadening slows down the nonlinear instability by increasing the effective internal detuning. Since the linear instability has a narrower frequency bandwidth, the filtering of the desired signal can be achieved by arranging for an appropriate combination of detuning and nonlinear effects.

The proposed detuning scheme can be used to enable laser intensities even exceeding the theoretical limit of the original fast compression scheme. The idea is shown in Fig. 1. As seen, the total external detuning, $\delta\omega_{\text{detuning}} \equiv \delta\omega$ (which governs the pump backscattering), can be made less than the plasma frequency detuning $\delta\omega_{\text{plasma}} \equiv \delta\omega_p$ (which governs the pumped pulse forward scattering). The larger plasma frequency gradient ω_p' can suppress Raman near-forward scattering of the pumped pulse into the Stokes pulse downshifted by the plasma frequency. The detuning gradient $\delta\omega'$ can be selected to suppress the pump Raman backscattering by noise. Thus, an appropriate combination of two different detuning mechanisms, plasma density gradient and pump chirping, can remove one of the major limitations on the output pulse intensities

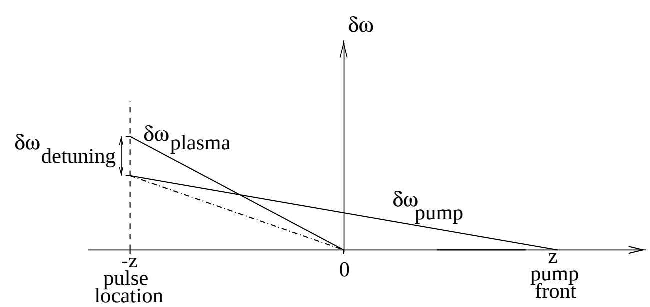

FIG. 1. Combined effect of plasma density gradient and pump chirp on external detuning in Raman backward amplifier. The detuning gradient is $\delta\omega' = \omega_p' - 2\omega_a'$, where ω_p' and ω_a' are derivatives of plasma and pump frequencies at the pulse location over distances from locations of exact resonance (z) and pump front ($2z$), respectively. The detuning gradient is characterized below by dimensionless parameter $q = 2(\omega_p' - 2\omega_a')c/\omega_p\omega_a a_0^2$, where a_0 is the nondepleted amplitude of pump vector potential in units of $m_e c^2/e$.

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in fast laser compressors by suppressing the Raman near-forward scattering instability of the pumped pulse. Remarkably, this can be achieved while still maintaining a highly efficient amplification of the pulse.

The new scheme retains the major advantages of conventional backward Raman amplification in gases: the pumped pulse reaches intensities much higher than that of the pump [5] and requirements on the pump laser quality are very modest (due to the averaging over the pumped pulse path of fluctuations in the pump intensity) [6]. Raman compression in gases, however, is made difficult by the dominance of forward scattering [6], a problem recently addressed in [7]. This Letter makes use of some of the Raman techniques in gases such as chirped pump [8]. It also employs plasma, where backward scattering dominates [9]. However, here the backward Raman amplification is transient rather than stationary, i.e., the damping of the Langmuir wave may be neglected. Compared to the results reported, e.g., by [7], the pumped pulses contemplated here have durations smaller by 10^3 and intensities higher by 10^6.

The slightly detuned transient 3-wave interaction for a Raman backward amplifier can be described by equations,

$$a_t + ca_z = \omega_p f b, \qquad b_t - cb_z = -\omega_p f^* a,$$
$$f_t + i\delta\omega f = -\omega a b^*/2, \qquad \delta\omega = \omega_p + \omega_b - \omega_a. \tag{1}$$

Here a and b are vector potential envelopes of the pump and pulse, respectively, in units of $m_e c^2/e \approx 5 \times 10^5$ V; f is the envelope of the electrostatic field of the Langmuir wave, normalized to $m_e c\omega_p/e = c\sqrt{4\pi m_e n_e} \approx \sqrt{n_e[\mathrm{cm}^{-3}]}$ V/cm; ω_p, ω_b, and ω_a are the plasma, laser-seed and laser-pump frequencies; $\delta\omega$ is the detuning from the 3-wave resonance; subscripts t and z denote time and space derivatives. The pulse duration is larger than ω_p^{-1}. Both lasers are circularly polarized. Self-nonlinearities of lasers and the Langmuir wave are neglected. Plasma ions are assumed to be immobile. For $\omega_b \gg \omega_p$, one may assume $\omega_a \approx \omega_b = \omega$ (except in calculating $\delta\omega$).

It is convenient to count time from the seed-laser front arrival to a given point $-z$ and to rescale variables,

$$\zeta = (t + z/c)\sqrt{\omega\omega_p}/2, \qquad \tau = -z\sqrt{\omega\omega_p}/c,$$
$$f = \bar{f}\sqrt{\omega/\omega_p}, \qquad \delta\omega = \sqrt{\omega\omega_p}\,\delta/2, \tag{2}$$

so that Eqs. (1) take the form

$$a_\zeta - a_\tau = \bar{f} b, \quad b_\tau = -\bar{f}^* a, \quad \bar{f}_\zeta + i\bar{f}\delta = -ab^*. \tag{3}$$

A smoothly varying detuning can be linearized near the resonance $\tau = 0$, where $\delta \approx \delta'\tau \equiv qa_0^2\tau$.

Note that small varying parts of ω_p and ω_a could have been included in the respective envelope definitions, which would reduce Eqs. (1) to the same form as for zero frequency detuning (but with more complicated boundary conditions at large $-z$ or $ct - z$). Such a transformation differs from that of [12], which removed the *wavelength* detuning from equations for the near-resonant 3-wave interaction. Indeed, when the group velocity of one wave is zero, the transformation [12] is reduced just to a *time-independent* phase shift for this wave. Such a shift cannot remove the *frequency* detuning from Eqs. (1). Thus, the *frequency* detuning is not covered in general by the transformation [12]. Also, the linear theory for parametric instabilities in inhomogeneous medium, suggested in [13], is not directly applicable to the *frequency* detuned Eqs. (1).

For the linear stage of the backscattering instability, when the pump depletion is negligible, $a \approx a_0 = $ const, the solution of (3) with $\delta \approx qa_0^2\tau$ can be written as

$$b(\zeta,\tau) = \frac{\partial}{\partial\zeta}\int d\zeta_1\, G(\zeta - \zeta_1, \tau) b(\zeta_1, 0),$$
$$G(\zeta,\tau) = \frac{1}{2\pi i}\int_C \frac{dp}{p}\exp\left[p\eta + \frac{i}{q}\ln\left(1 - \frac{iq}{p}\right)\right], \tag{4}$$

where $\eta \equiv a_0^2\zeta\tau$ and the integration contour C in the complex plane p encompasses in the positive direction singularities at $p = 0$ and $p = iq$.

For $|q|\sqrt{\eta} \ll 1$, the Green's function G reduces to that in a uniform plasma, where $q = 0$ and $G = I_0(2\sqrt{\eta})$, in agreement with linear theory for backscattering in uniform media [10]. Here I_0 is the modified Bessel function. In the domain $\eta \gg 1$ (but still $|q|\sqrt{\eta} \ll 1$, which is possible when $|q| \ll 1$), one has $G \approx \exp(2\sqrt{\eta})/2\sqrt{\pi\sqrt{\eta}}$. In original variables, $\eta = a_0^2\omega\omega_p(t + z/c)(-z)/2c$, so that the maximum of G, reached at $z = -ct/2$, moves with the speed $-c/2$ and increases with the peak growth rate for the monochromatic wave instability $a_0\sqrt{\omega\omega_p/2}$.

The effect of detuning becomes noticeable at $|q|\sqrt{\eta} \sim 1$, when the backscattering instability makes about $|q|^{-1}$ exponentiations. For $|q| \ll 1$, the Green's function (4) can be evaluated by the method of steepest descent. It increases exponentially with η up to the point $\eta_M = 4/q^2 \gg 1$. In the domain $\eta \gg 1$, but $\eta_M - \eta \gg 1$, G can be approximated by the formula

$$G = \frac{\exp[\sqrt{\eta(1 - \eta/\eta_M)} + (\sqrt{\eta_M} + i)\,\mathrm{arcctg}\sqrt{\eta_M/\eta - 1} - i\eta/\sqrt{\eta_M}\,]}{2\sqrt{\pi\sqrt{\eta(1 - \eta/\eta_M)}}}. \tag{5}$$

For $\eta \ll \eta_M$, Eq. (5) reduces to the exact resonance case. At the applicability limit $\eta_M - \eta \sim 1$, $|G|$ attains its maximum value, $\max_\eta |G| \sim \exp(\pi/|q|)$. For larger η, $|G|$ drops abruptly. Hence, the maximum amplification factor for the integrated amplitude, $u \equiv \int^\zeta d\zeta\, b(\zeta,\tau)$, of a small narrow seed b is $\sim \exp(\pi/|q|)$.

This gives the threshold for stabilizing the pump by the detuning gradient q in a noisy backward amplifier. If the initial value of the integrated amplitude u is much smaller than $\exp(-\pi/|q|)$, it remains forever smaller than 1. Such a small seed never reaches the nonlinear stage of instability. The pump depletion similarly remains small.

Consider now the nonlinear evolution of the desired signal, a seed pulse with the initial integrated amplitude u larger than $\exp(-\pi/|q|)$, which is sufficient to deplete the pump before making $\pi/|q|$ exponentiations. In the nonlinear stage of amplification, the pump depletion scale decreases, while the amplification time increases, since a fixed pump acts relatively more slowly on larger signals. When the ζ scale becomes much smaller than the τ scale, the term $|a_\tau| \ll |a_\zeta|$ in (3) can be neglected. Then the resulting set of nonlinear equations has a new symmetry, allowing the self-similar substitution $\eta = a_0^2 \tau \zeta, a(\zeta,\tau) = a_0 \tilde{a}(\eta), \bar{f}(\zeta,\tau) = a_0 \tilde{f}(\eta)$, $b(\zeta,\tau) = a_0^2 \tau \tilde{b}(\eta)$, with new functions satisfying the nonlinear ordinary differential equation (ODE)

$$\tilde{a}_\eta = \tilde{f}\tilde{b}, \quad \tilde{f}_\eta + iq\tilde{f} = -\tilde{a}\tilde{b}^*, \quad (\eta\tilde{b})_\eta = -\tilde{a}\tilde{f}^*, \tag{6}$$

and the initial condition $\tilde{a} \to 1$ at $\eta \to +0$. The solution depends on a single parameter, say $\tilde{b}(+0) = \epsilon_1$.

To analyze Eqs. (6) for arbitrary pump depletion, it is useful to rewrite them in real form introducing real amplitudes and phases by $\tilde{a} = Ae^{i\alpha}, \tilde{b} = Be^{i\beta}, \tilde{f} = Fe^{i\phi}$, to obtain a set of six real first-order equations,

$$\begin{aligned} &A_\eta = FBcs, \quad F_\eta = -ABcs, \quad (\eta B)_\eta = -AFcs\,, \\ &\alpha_\eta = FBs/A, \quad \phi_\eta = ABs/F - q, \quad \beta_\eta = AFs/B\eta\,, \\ &cs \equiv \cos\Delta, \quad s \equiv \sin\Delta, \quad \Delta \equiv \beta + \phi - \alpha\,. \end{aligned} \tag{7}$$

These six equations can be reduced to a set of just two real first-order equations. Reduction of order by 2 is due to the invariance of (7) to the 2-parameter group of constant phase shifts that do not change Δ. An extra reduction by one is due to the integral of (7), $A^2 + F^2 = C^2 = \text{const}$, corresponding to the conservation of the joint number of the pump and Langmuir quanta in the 3-wave decay interaction, which allows one to introduce function U, such that $A = C\cos(U/2), F = C\sin(U/2)$. To proceed with the reduction, the identity $(\eta^2 B^2 \beta_\eta)_\eta = (\eta^2 B^2)_\eta q/2$ can be derived from Eqs. (7). Taking into account the regularity of all fields at $\eta \to +0$, it follows that $\beta_\eta = q/2$. Thus, the amplified pulse acquires a linear frequency chirp exactly half of the external detuning gradient, while the central pulse frequency drift compensates for half of the external detuning.

Substitution of the above A, F, and integral $\beta_\eta = q/2$ in Eq. (7) gives a closed set of two first-order ODE,

$$U_\eta = -2B\cos\Delta, \qquad 2(\eta B)_\eta = -C^2 \sin U \cos\Delta\,, \tag{8}$$

$$q\eta B = C^2 \sin U \sin\Delta\,. \tag{9}$$

The initial value of U is $U(+0) = -2\arctan\epsilon_1 \equiv \epsilon$ [since $A(+0) = 1$ and $F(+0) = -B(+0) = -\epsilon_1$].

For $q = 0$, when one can take $\Delta = 0$, Eqs. (8) are equivalent to the second order ODE $(\eta U_\eta)_\eta = C^2 \sin U$. With $\eta = \xi^2/4C^2$, it reduces to Eq. (6) of [2], $U_{\xi\xi} + U_\xi/\xi = \sin U$, corresponding to the exactly resonant interaction.

For $q \neq 0$, the equations can also be put in the form $U_\eta = -\sin(2\Delta)\sin U/q\eta, U(+0) = \epsilon, \Delta_\eta = 2\sin^2\Delta \times \cos U/q\eta - q/2, \Delta(+0) = 0$. The limit $q \to 0$ has a peculiarity. For $q \neq 0$, U varies only inside $(0,\pi)$ interval, while for $q = 0$, $U \equiv U^o$ varies inside $(0, 2\pi)$ interval and tends to the limit π at $\xi \to \infty$ oscillating around this value ("π-pulse" solution). The points where $U^o = \pi$ are not zeros of B, but are rather close to maximums of the pulse intensity B^2. For a small $q \neq 0$, one can neglect the variation of $B \approx B_*$ near a point $\eta = \eta_*$, where $U^o = \pi$ and, integrating there (8) with Δ from (9), to get $\tilde{U} \equiv \pi - U \approx |B_*|\sqrt{q^2\eta_*^2/C^4 + 4(\eta - \eta_*)^2}$.

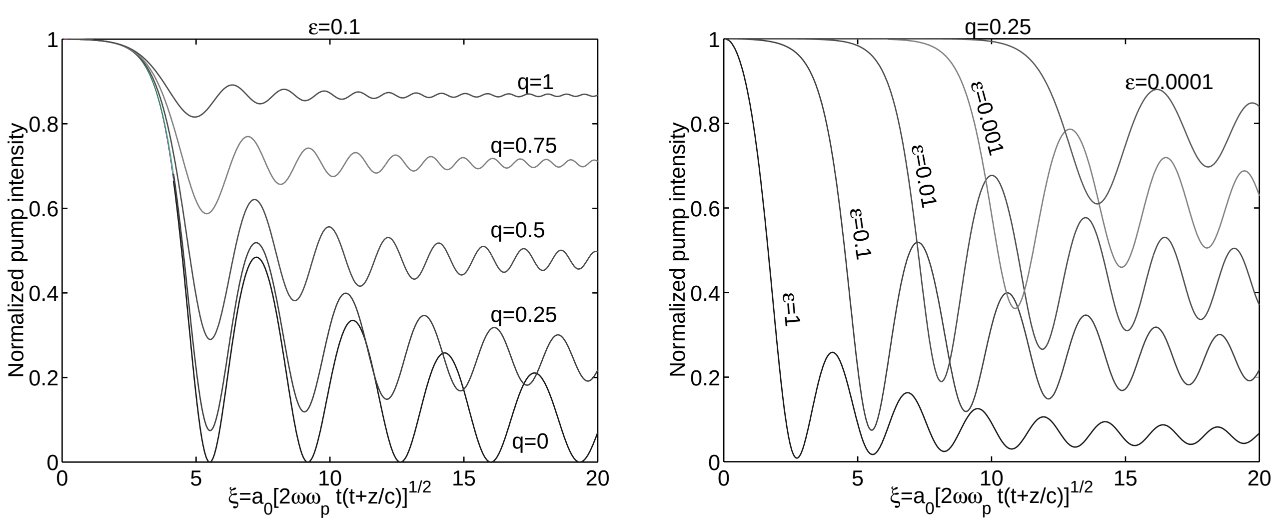

FIG. 2. Normalized pump intensity A^2 for different values of integrated seed amplitude $\epsilon_1 = -\tan(\epsilon/2)$ and detuning gradient q. The pump depletion increases with the increase of $|\epsilon|$ and decreases with the increase of $|q|$.

The function has a minimum $\tilde{U}_* \approx |B_* q|\eta_*/C^2$ at η_*. The corresponding $\sin\Delta \approx \tilde{U}_*/\tilde{U}$ tends to zero as $q \to 0$ outside a narrow ($|\eta - \eta_*| \leq q\eta_*$) vicinity of η_*. In the outer domain, the solution is close to the π pulse up to the terms $\propto q^2$.

The pump intensity has a minimum $A_*^2 \approx C^2\tilde{U}_*^2/4 \approx B_*^2 q^2 \eta_*^2/4C^2$ at η_*. The zero-order values of B_* and η_* can be taken from the π-pulse solution. For $\epsilon \ll 1$, the integrated amplitude of the leading spike of the π-pulse wave train is close to the classical 2π-pulse solution of the sine-Gordon equation, $U^o \approx 4\arctan(\epsilon e^{\xi}/4\sqrt{2\pi\xi})$ [2]. The point $U^o = \pi$ is located at $\xi_* \approx \ln(4\sqrt{2\pi\xi_*}/\epsilon)$ ($\eta_* = \xi_*^2/4$). Calculating a small deviation from the 2π pulse, one can show that the pulse intensity is $B_*^2 \approx 4/(\xi_* + 1)^2$. Then, $A_*^2 \approx q^2\xi_*^4/16(\xi_* + 1)^2$. This simple asymptotic ($q \to 0, \epsilon \to 0$) formula agrees with numerical results presented in Fig. 2. For $q = 0.25, \epsilon = 0.1$, both the analytical and numerical solutions give $A_*^2 \approx 8\%$. Even for $q = 0.5, \epsilon = 0.1$, the agreement is still reasonable: $A_*^2 \approx 33\%$ analytical versus $A_*^2 \approx 29\%$ numerical.

The numerical solution of Eqs. (8) confirms the linear theory prediction that detuning suppresses the pump instability to noise: as seen from Fig. 2, very small seeds virtually do not deplete the pump. Yet, it is possible to maintain a high efficiency of the useful amplification process that starts from a moderately small initial seed and can tolerate a large enough detuning gradient q.

After the pumped pulse passes, there is no further depletion. The final pump depletion depends on parameters ϵ and q and may take any value in the $(0, 1)$ interval. It corresponds to $U(+\infty)$ taking any value in the $(0, \pi)$ interval. Thus, all "less-than-π pulses" can appear in detuned Raman amplifiers, in contrast to the well-known exactly resonant case where just π pulses appear [11].

Figure 3 demonstrates that the self-similar solution is an attractor. It also shows that the final pump depletion (say, 76% for $\epsilon = 0.1, q = 0.25$, or 52% for $\epsilon = 0.1, q = 0.5$) can be somewhat increased by working near the threshold of the Langmuir wave breaking. The near-threshold breaking occurs near the leading maximum of the pumped pulse intensity, which prevents the pulse energy from scattering back to the pump. It suppresses the second and further spikes in the amplified pulse wave train.

In summary, a self-similar attractor solution is found for the detuned Raman backscattering of a laser pump into a short counterpropagating seed laser pulse. The amplified pulse acquires a frequency shift exactly equal to half the imposed detuning. A nonlinear filtering effect was identified. The solution generalizes the classical π-pulse regime solution for the transient exactly resonant 3-wave interaction.

This generalized solution of the π-pulse regime should be applicable to a broad range of phenomena [11]. However, the application discussed here, the selective suppression of unwanted Raman instabilities through a combination of detuning and nonlinear filtering, makes robust the effect of fast compression. This application itself may make feasible a new generation of compact ultrahigh energy compressors of short laser pulses.

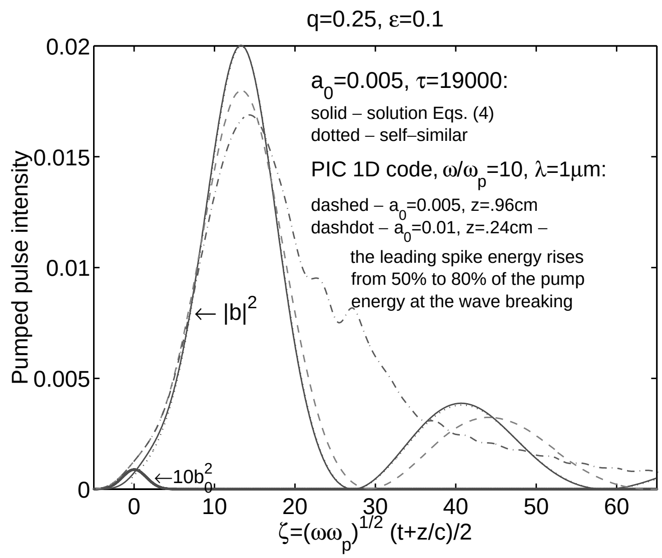

FIG. 3. Coincidence of the solid and dotted lines indicates that the self-similar solution is the attractor for Eqs. (3) with small and narrow initial seed pulse (small spike at $\zeta = 0$, where the intensity b_0^2 is multiplied by 10 to make it distinguishable). The dashed and dashed-doted lines, based on 1D particles-in-cell simulation, show that the Langmuir wave breaking (occurring at $a_0 \approx 0.008$ for $\omega/\omega_p = 10$) raises the total energy extraction from the pump and suppresses secondary spikes.

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[1] G. Shvets *et al.*, Phys. Rev. Lett. **81**, 4879 (1998).
[2] V. M. Malkin, G. Shvets and N. J. Fisch, Phys. Rev. Lett. **82**, 4448 (1999).
[3] G. A. Mourou, C. P. J. Barty, and M. D. Perry, Phys. Today **51**, No. 1 22 (1998).
[4] M. Tabak *et al.*, Phys. Plasmas **1**, 1626 (1994).
[5] M. Maier, W. Kaiser, and J. A. Giordmaine, Phys. Rev. Lett. **17**, 1275 (1966); Phys. Rev. **177**, 589 (1969).
[6] J. R. Murray *et al.*, IEEE J. Quantum Electron. **QE-15**, 342 (1979).
[7] H. Nishioka *et al.*, IEEE J. Quantum Electron. **29**, 2251 (1993); E. Takahashi *et al.*, Fusion Eng. Des. **44**, 133 (1999).
[8] J. A. Caird, IEEE J. Quantum Electron. **QE-16**, 489 (1980).
[9] C. E. Capjack, C. R. James, and J. N. McMullin, J. Appl. Phys. **53**, 4046 (1982).
[10] D. L. Bobroff and H. A. Haus, J. Appl. Phys. **38**, 390 (1967).
[11] G. L. Lamb, Jr., Phys. Lett. **29A**, 507 (1969); Rev. Mod. Phys. **43**, 99 (1971); V. E. Zakharov, JETP Lett. **32**, 589 (1980); S. V. Manakov, JETP Lett. **35**, 237 (1982); V. A. Gorbunov and V. R. Startsev, Bull. Acad. Sci. USSR **48**, 120 (1984); J. Coste and C. Montes, Phys. Rev. A **34**, 3940 (1986).
[12] A. H. Reiman, A. Bers, and D. J. Kaup, Phys. Rev. Lett. **39**, 245 (1977); A. Reiman, Rev. Mod. Phys. **51**, 311 (1979).
[13] M. N. Rosenbluth, Phys. Rev. Lett. **29**, 565 (1972).