Suppression of Superluminous Precursors in High-Power Backward Raman Amplifiers

Yu. A. Tsidulko, V. M. Malkin, and N. J. Fisch
Department of Astrophysical Sciences, Princeton University, Princeton, New Jersey 08544
(Received 9 July 2001; published 28 May 2002)

The very promising scheme for producing ultrapowerful laser pulses through Raman backscattering of pump lasers in plasmas can be jeopardized by superluminous precursors. Growing from the leading part of the Raman pumped seed pulse, these precursors can disturb the plasma and the pump ahead of the pumped pulse. These ruinous effects, however, might be averted by a detuning, large enough to suppress the precursors, yet small enough to allow the desired backscatter effect.

DOI: 10.1103/PhysRevLett.88.235004 PACS numbers: 52.38.Bv, 42.65.Yj, 52.35.Mw

The fast compression of a long laser pump into a short pumped pulse by means of stimulated backward Raman scattering (SBRS) in a plasma may lead to ultrahigh laser powers [1]. In a 1D model, a very efficient attractor regime exists (known as the π-pulse regime), in which nearly all incident pump power goes into a short seed pulse, which contracts as its amplitude grows. Two-dimensional numerical simulations now confirm essential aspects of the 1D model [2]. Although no experiments have clearly entered this π-pulse regime, experiments have been set up to verify the ultrahigh amplification concept [3,4], and some amplification of a short pulse in a plasma capillary has been reported [3].

In the π-pulse regime, even as the pumped seed pulse grows much more intense than the pump, the energy still flows from the pump to the seed via the three-wave resonant decay of pump photons to lower frequency seed photons and Langmuir plasmons. The opposite process is forbidden so long as the seed encounters fresh plasma layers free of Langmuir waves.

While having a large basin of attraction, this desirable π-pulse regime might be completely upset if the leading part of the seed pulse is not of the desired form. We will refer to such leading parts as precursors. Our intent is to explore their consequences and how to mitigate them. The precursors, also attractor solutions, frequently appear in amplifying, absorbing, and anomalously dispersive media (see, e.g., [5]). We show that, in regimes otherwise suitable for the fast compression, the pump energy can be backscattered into such precursors, each of which, nonlinearly stabilized, consumes only a small portion of the pump energy. What ensues is a long precursor wavetrain, rather than the efficient compression of pump energy into a single intense spike. Moreover, even when some pump power penetrates the precursor wavetrain to the original seed pulse, the desired amplification can be inhibited by Langmuir waves excited by the precursors ahead of the seed.

This paper is organized as follows: First, we demonstrate the deleterious role that the precursors can play. We then show that the formation of precursors is more sensitive to detuning than is the desired amplification. We then identify the conditions for avoiding precursors, yet allowing the compression effect.

To see the role that precursors can play, consider first the precisely resonant backward Raman amplification (BRA) [1] in the form

$$a_t + ca_z = Vbf, \qquad f_t = -Vab^*,$$
$$b_t - cb_z = -Vaf^*, \tag{1}$$

where $V = \sqrt{\omega\omega_e/2}$, with laser frequency $\omega \gg \omega_e$, the electron plasma frequency; the dimensionless space-time envelopes, a and b, of the pump and the seed electron quiver velocities, are normalized so that the pump power density is $I_a = \pi c(m_e c^2/e)^2|a|^2/\lambda^2 = 2.736 \times 10^{18}|a|^2/\lambda^2[\mu\mathrm{m}]$ W/cm^2, and similarly for seed b; f is the normalized Langmuir wave envelope; λ is the laser wavelength, t the time, z the distance in the direction of the pump propagation; subscripts denote derivatives; and c is the speed of light. The normalized wave amplitudes can be renormalized to the pump maximum amplitude a_0, whereupon the time is normalized to the natural growth time $1/Va_0 \equiv 1/\gamma_0$ and length to c/Va_0. What is left to choose is the precursor shape.

A short seed pulse gives rise to the compression solution, with no precursors [1]. But a longer, even if less intense, prepulse can produce precursors. Figure 1 shows the solution of Eqs. (1) for an initial seed comprised of two left-going pulses $b = B + \tilde{b}$, where B is a desired short seed and $\tilde{b}$ is a longer, presumably unwanted, prepulse signal. The desired seed is a Gaussian, $B(0,t) = 4a_0/\{5\sqrt{\pi}\exp[(10\gamma_0 t)^2]\}$, accompanied by a longer prepulse $\tilde{b}$ of three different forms: In Fig. 1a, the prepulse is a longer Gaussian $\tilde{b}(0,t) = a_0/\{5\sqrt{\pi}\exp[(2\gamma_0 t)^2]\}$. The sequence of snapshots, viewed in the seed frame, reveals the evolution of a nearly π-pulse solution, with pump depletion. Figure 1b shows a hyperbolic secant shaped prepulse, $\tilde{b}(0,t) = a_0/[5\sqrt{\pi}\cosh(2\gamma_0 t)]$, generating a wavetrain of quasistationary superluminous precursors ahead of the short seed. This wavetrain only partially depletes the pump, excites Langmuir waves ahead of the seed, and results in little useful BRA. Figure 1c shows a Lorentzian-shaped prepulse, $\tilde{b}(0,t) = a_0/\{5\sqrt{\pi}[1 + (2\gamma_0 t)^2]\}$, generating an even longer precursor wavetrain that completely destroys the useful high-power compression effect. Some energy is backscattered, but not compressed.

 © 2002 The American Physical Society

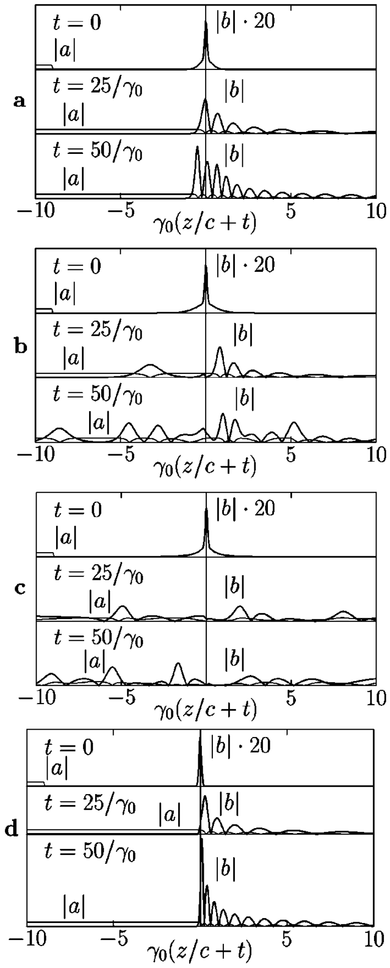

FIG. 1. For exactly resonant BRA, evolution of a short Gaussian seed accompanied by a longer prepulse of (a) Gaussian, (b) hyperbolic secant, and (c) Lorentz form. For comparison, (d) no prepulse.

As seen in Fig. 1, the generation of precursors is very sensitive to the form of the prepulse. This is because the seed front at first evolves within the linearized version of Eqs. (1), with constant pump ($a = a_0$). The SBRS instability then has the local growth rate,

$$\gamma = cp/2 + \sqrt{c^2p^2/4 + \gamma_0^2}, \tag{2}$$

where $\gamma_0 = Va_0$ is the growth rate for the monochromatic wave instability, and $p = \partial\,\mathrm{Re}S/\partial z$ is the local logarithm slope of the seed front $b = e^S$. Note also that the seed amplitude (i.e., the location where $\mathrm{Re}S = \mathrm{const} \Rightarrow d\,\mathrm{Re}S = 0$) propagates with the superluminous velocity $v(p) = -\gamma/p$. Since the p-derivative of front velocity $v = -\gamma/p$ is positive, the sign of v_z coincides with that of p_z. It follows that the front with $p_z > 0$ stretches, becoming relatively less steep with time, while the front, with $p_z < 0$, steepens with time. The formation of precursors outrunning the bulk of the seed front should then be expected in the case of the nonsteepened front, $p_z \geq 0$, as indeed observed numerically.

Note that the superluminous dynamics, when $|v| > c$, is produced by simple reshaping of the signal as it is amplified, and not by information transfer with a speed exceeding c. Causality is respected, since the leading part of the front is not affected by the dynamics behind it.

The leading front grows exponentially until limited by pump depletion. Then, for a nonsteepened front, $p_z \geq 0$, nonlinear stabilization should occur. This is because the pump depletion depends on the encountered area under the seed envelope, so that this area cannot grow after substantial pump depletion. Hence, for a nonsteepened front, the leading precursor amplitude also saturates. The stabilized precursor should be approximated well by the classical 2π-pulse stationary moving with velocity v, corresponding to the given front slope p,

$$u = 4\arctan e^{p(z-vt-z_0)}, \quad b = \frac{a_0\sqrt{\gamma(\gamma + pc)}}{\gamma_0\cosh[p(z - vt - z_0)]}. \tag{3}$$

The dimensionless "area" u under the pulse envelope tends to 2π at large z, hence the name "2π-pulse."

For an exact 2π-pulse, the pump amplitude $a = a_0\cos(u/2)$ is completely restored after the pulse passes. No Langmuir wave $f \propto \sin(u/2)$ is left behind. But the 2π-pulse tail is unstable, so a second trailing precursor is formed, eventually resulting in a long wavetrain of precursors. Each of these precursors slightly depletes the pump and leaves some Langmuir waves behind. Both of these effects hinder the efficient amplification of the original seed pulse. A train of 2π-precursors can be seen clearly in Figs. 1b and 1c, corresponding to initially nonsteepened prepulse fronts that obey the condition for precursor production, namely $p_z \geq 0$.

In contrast to the examples given in Figs. 1b and 1c, the Gaussian prepulse shown in Fig. 1a satisfies the opposite condition $p_z < 0$. This example features amplification much resembling a π-pulse. Despite the resemblance, the compression may not be as efficient, because the π-pulse is actually seeded by the prepulse itself, which is weaker and longer, rather than by the carefully prepared short seed. It is effective integrated initial amplitude of the seed that sets the efficiency [1]. The π-pulse seeded by the prepulse, the prepulse shadow the well-prepared seed. Since the integrated amplitude of the prepulse, $\int \tilde{b}(0,t)\,dt$, is small, the energy compression is less efficient, with the leading spike capturing a smaller portion of the pump energy.

Note that the nonlinear stage of the instability is amenable, in principle, to analytic approaches. Equations (1) are integrable by the inverse scattering transform [6]. However, analytical solutions, even in the form [7], are sufficiently cumbersome that the dynamics of interest is hard to extract. Recent work [8], which takes into account wave damping and studies stationary moving structures, also leaves unaddressed the questions of interest here, in particular, the extent to which these precursors might inhibit the useful amplification and compression.

Since the precursors are not efficiently compressed, and hence have smaller bandwidth than the compression solutions, it may be expected that the precursors can be selectively suppressed, like noise amplification in [9], by detuning the resonance. To evaluate this possibility, the frequency detuning is accounted for by modifying just one of the Eqs. (1), say the equation for f, to read

$$f_t + \imath\delta\omega f = -Vab^*. \quad (4)$$

We now linearize the frequency detuning $\delta\omega$ over z near the exact resonance point, $z = z_r$, and introduce the dimensionless detuning gradient q, $\delta\omega = q\gamma_0^2\tilde{z}/c$, $\tilde{z} = z - z_r$, in the same way as in [9]. In practice, this detuning gradient can be introduced either by chirping the pump pulse or by introducing an axial density (Raman frequency) gradient.

In the detuned BRA regimes, an unstable small signal makes just a finite number of exponentiations before leaving the resonant domain. Provided there is sufficient contrast between the desired short seed and the precursor background, the detuning may stabilize the prepulse at a benign level, while allowing the useful seed to enter the nonlinear regime. As the useful seed is amplified and compressed, it may be expected to tolerate the growing detuning. Moreover, even for insufficient seed-background contrast, detuning should suppress the generation of precursors. This is because the leading precursor no longer moves with a superluminous velocity once it leaves the resonant domain. Therefore, the distance between the leading precursor and the useful seed no longer increases. This limits the number of precursors that could be generated ahead of the desired seed pulse.

These anticipated effects are confirmed by numerical solutions. Figures 2a–2c show the amplification for the same seed and prepulses as in Figs. 1, but now with a detuning gradient. For comparison, we show in Fig. 2d the detuned amplification for the desired short seed $b(0,t) = a_0/\{\sqrt{\pi}\exp[(10\gamma_0 t)^2]\}$, in the absence of a prepulse but in the presence of the detuning gradient. These cases all feature what essentially are compression solutions. Figure 2a shows that the detuning suppresses substantially, although not entirely, the BRA produced by the precursor Gaussian seed. Figures 2b and 2c show that the detuning indeed suppresses precursor wavetrains generated by the hyperbolic secant and Lorentz prepulse seeds, respectively. These precursors no longer deplete the pump noticeably. While Langmuir waves are excited ahead of the useful seed, they do not distort appreciably the usefully amplified pulses nor diminish appreciably the amplification efficiency. For instance, at $t = 50/\gamma_0$, the largest spike in Fig. 2b contains just 10% less energy than the largest spike would contain if the precursor background were absent, as given in Fig. 2d. The width, however, of the largest spike in Fig. 2b is about twice that in Fig. 2d. Thus, in the presence of the detuning gradient, good compression solutions are obtained for each of the precursor backgrounds.

A further concern is that the precursor background might fluctuate randomly in the directions transverse to the propagation direction in a manner that would affect the output focusability, even if not the amplification itself. We determine, therefore, the levels of prepulse background for which such distortions do not occur.

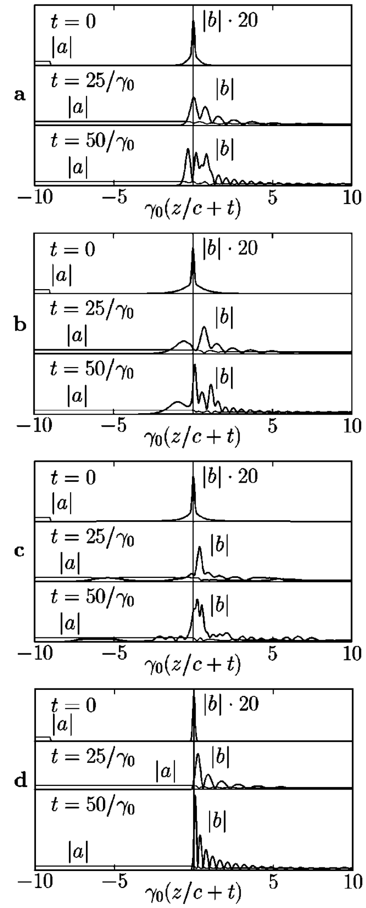

FIG. 2. Under detuning $q = 0.25$, evolution of the short Gaussian seed accompanied by a longer prepulse seed of (a) Gaussian, (b) Poisson, and (c) Lorentz form. For comparison, (d) no prepulse.

Note first that there is a further helpful effect due to detuning that favors the amplification of the short useful seed over the longer unwanted prepulse. A short small seed pulse makes $\pi/|q|$ exponentiations between the exact resonance point z_r and off-resonant domain. For z_r located precisely at the boundary where the seed enters the plasma, $z_r = 0$, this is the actual number of exponentiations that the short small seed makes. (However, for z_r located inside the amplifying plasma layer, not too close to its boundaries, the total number of exponentiations of an originally short small pulse is $2\pi/|q|$, since it traverses twice the distance between z_r and the off-resonance domain.) Within the resonant domain, a short pulse evolution is less affected by the detuning, so that, during the linear stage of instability, the pulse maximum moves with speed half the front speed. Therefore, the maximum spends in the resonant domain twice the time than the front does. Thus, the pulse maximum experiences twice the number of exponentiations, than it would experience were it to move at the front speed. However, the long prepulse and any incipient precursors move with approximately the same speed as the seed front. Thus, between $z_r = 0$ and where all resonance is completely detuned, the prepulse should experience only half the number of exponentiations ($\pi/2|q|$) that the pulse maximum experiences. This explains how even modest detuning can result in precursor suppression together with useful pulse amplification, as shown in Fig. 2, even for poor initial seed-background contrasts.

To describe this effect formally, consider the following exact solution of the linearized detuned BRA equations: $b \propto e^S$, with

$$S = \gamma(t + \tilde{z}/c) + \imath q^{-1} \ln(1 + \imath q \gamma_0^2 \tilde{z}/c\gamma), \quad (5)$$

which can be verified by substitution. Close to the resonance, $|q\gamma_0^2\tilde{z}/c\gamma| \ll 1$, Eq. (5) simplifies to $S = \gamma t + p\tilde{z}$ with $p = (\gamma - \gamma_0^2/\gamma)/c$, i.e., the detuned linear solution (5) reduces to the exactly resonant solution describing the front of exponential slope p moving with the superluminous velocity $v = -\gamma/p$, where γ is given by the formula (2). According to Eq. (5), $\mathrm{Re}S = q^{-1}\arctan(q\gamma_0^2\tilde{z}/c\gamma) + \gamma(t + \tilde{z}/c)$, so that, at a fixed distance from the front $t = -\tilde{z}/c$, the pulse amplitude $|b| \propto e^{\mathrm{Re}S}$ indeed makes $\pi/2|q|$ exponentiations between the resonance point $\tilde{z} = 0$ and the off-resonance domain $|q\gamma_0^2\tilde{z}/c\gamma| \gg 1$.

Using the same solution (5) of the linearized detuned equations, we can calculate the amplitude of resonant Langmuir wave excited ahead of the short useful seed by the prepulse with the front growth rate γ. The useful amplified pulse is not significantly distorted when the resonant Langmuir wave ahead of the pulse is much smaller than the Langmuir wave excited by the useful pulse itself. This condition restricts the amplitude of the prepulse seed $\tilde{b}_0$ to

$$\tilde{b}_0 \ll a_0(\sqrt{|q|} + \gamma/\gamma_0|q|)\exp(-\pi/2|q|). \quad (6)$$

Numerical solutions of the nonlinear detuned equations with $z_r = 0$ indeed confirm that prepulse backgrounds satisfying condition (6) do not affect the efficient useful amplification shown in Fig. 2d. For the exact resonance point z_r located well inside the plasma layer, the exponent $-\pi/2|q|$ in (6) should be replaced by $-\pi/|q|$, because of the doubling of the effective resonant domain.

The conditions identified here for inhibition of the BRA by superluminous precursors, both for the resonant and the detuned cases, are important in interpreting existing experiments and in prescribing how seed pulses must be prepared for accessing the favorable BRA regimes. In particular, backscattering experiments such as [3,4] must show an amplification that is sufficiently time resolved, so that amplification not be confused with mere pump reflection. Also, experiments documenting backward Raman scattering of laser pulses by thermal noise, such as [10], as opposed to experiments employing properly seeded counterpropagating pulses, must clearly be interpreted in a very different manner. The former are very unlikely to exhibit compression solutions. Note that the requirements for properly prepared seed pulses are technologically modest. However, the conditions derived here could challenge a self-seeded scheme of BRA.

The 1D model employed here is justified in the counterpropagating geometry, since all the processes of our interest develop much faster than transverse filamentation instabilities. Also, so long as Langmuir wave breaking regimes are avoided, the fluid description should be adequate. For the wave breaking regimes, that require a kinetic description of the plasma, the results obtained here may be useful as a guide [1].

In summary, we identified how superluminous precursors inhibit ultrapowerful laser compression in plasma. We derived the conditions on the seed profile that must be avoided. We further identified helpful detuning effects, specifying quantitatively the requirements needed for highly efficient BRA in plasma. Moreover, we identified conditions for which the amplified wave fronts are not distorted by prepulse-generated Langmuir waves, so that the output can be well focusable. The effects identified and calculated here are critical for understanding the high-power Raman pulse compression effect in plasma, which may serve as the basis for a new generation of more powerful lasers. However, the results here may be much more generally applied, since the general 3-wave equations that describe Raman backscattering in plasma, both including and not including detuning, also describe effects in a variety of Raman gases, liquids, and fibers.

This work is supported under DOE DEFG030-98DP00210 and DARPA.

[1] V. M. Malkin, G. Shvets, and N. J. Fisch, Phys. Rev. Lett. **82**, 4448 (1999).
[2] G. M. Fraiman *et al.,* Phys. Plasmas (to be published); I. Y. Dodin *et al.,* Sov. Phys. JETP (to be published); P. Mardahl *et al.,* Phys. Lett. A (to be published).
[3] Y. Ping *et al.,* Phys. Rev. E **62**, R4532 (2000).
[4] R. Kirkwood *et al.,* Bull. APS, DPP (2001) (paper KP1.105); M. Dreher *et al.,* in Proceedings of the 2nd International Conference on Superstrong Fields in Plasmas, Varenna, Italy, 2001.
[5] G. L. Lamb, Jr., Phys. Lett. A **29**, 507 (1969); Rev. Mod. Phys. **43**, 99 (1971); C. G. B. Garrett and D. E. McCumber, Phys. Rev. A **1**, 305 (1970); S. Chu and S. Wong, Phys. Rev. Lett. **48**, 738 (1982); E. Picholle *et al.,* Phys. Rev. Lett. **66**, 1454 (1991); R. Y. Chiao, A. E. Kozhekin, and G. Kurizki, Phys. Rev. Lett. **77**, 1254 (1996); A. Yu. Andreev and D. A. Kirzhnits, Usp. Fiz. Nauk **166**, 1135 (1996); A. N. Oraevsky, Usp. Fiz. Nauk **168**, 1311 (1998); L. J. Wang, A. Kuzmich, and A. Dogariu, Nature (London) **406**, 277 (2000).
[6] V. E. Zakharov and S. V. Manakov, Sov. Phys. JETP Lett. **18**, 243 (1973); Sov. Phys. JETP **42**, 842 (1976); D. J. Kaup, Stud. Appl. Math. **55**, 9 (1976).
[7] S. C. Chiu, J. Math. Phys. **19**, 168 (1978).
[8] C. Montes, A. Picozzi, and D. Bahloul, Phys. Rev. E **55**, 1092 (1997).
[9] V. M. Malkin, G. Shvets, and N. J. Fisch, Phys. Rev. Lett. **84**, 1208 (2000).
[10] C. Rousseaux *et al.,* Phys. Rev. Lett. **74**, 4655 (1995).