Exceeding the leading spike intensity and fluence limits in backward Raman amplifiers

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V. M. Malkin,[1] Z. Toroker,[2] and N. J. Fisch[1]
[1)] *Department of Astrophysical Sciences, Princeton University, Princeton, NJ USA 08540*
[2)] *Department of Electrical Engineering, Technion Israel Institute of Technology, Haifa 32000, Israel*

The leading amplified spike in backward Raman amplifiers can reach nearly relativistic intensities before the saturation by the relativistic electron nonlinearity. The saturation sets an upper limit to the largest achievable leading spike intensity. It is shown here that this limit can be substantially exceeded by the initially subdominant spikes, which surprisingly outgrow the leading spike after its nonlinear saturation. Furthermore, an initially negligible group velocity dispersion of the amplified pulse in strongly undercritical plasma appears to be capable of delaying the longitudinal filamentation instability in the nonlinear saturation regime. This enables further amplification of the pulse to even larger output fluences.

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I. INTRODUCTION

The backward Raman amplification (BRA) of laser pulses in plasmas[1] is potentially capable of producing laser powers about a million times larger than the chirped pulse amplification (CPA)[2] at the same wavelengths within the same size devices[3,4]. The BRA advantage can be even greater at laser wavelengths much shorter than 1/3 micron, where material gratings used by CPA cannot operate. The possibility of reaching nearly relativistic unfocused intensities in backward Raman amplifiers has been in principle demonstrated experimentally as well[5–11].

The major physical processes that may affect BRA include the amplified pulse filamentation and detuning due to the relativistic electron nonlinearity[1,4,12–16], parasitic Raman scattering of the pump and amplified pulses by plasma noise[1,3,4,17–19], generation of superluminous precursors of the amplified pulse[20], pulse scattering by plasma density inhomogeneities[21], pulse depletion and plasma heating through inverse bremsstrahlung[22–25], the resonant Langmuir wave Landau damping[22,24,26–31] or breaking[1,3,4,32,33], and other processes (see for example [34–37]). Most of these deleterious processes can be mitigated by appropriate preparation of laser pulses and plasmas, choosing parameter ranges and selective detuning of the Raman resonance. Ultimately, the output intensity limit appears to be imposed primarily by the relativistic electron nonlinearity, causing saturation of the dominant leading spike growth[16]. The major goal of this paper is to explore the possibility of extending BRA beyond this theoretical limit for the largest achievable unfocused intensity and fluence of the output pulses.

II. BASIC EQUATIONS

To capture the effects of interest here, one needs to take into account, apart from the resonant Raman backscattering, the relativistic electron nonlinearity (REN) and group velocity dispersion (GVD) of the amplified laser pulse. This is because the amplified pulse reaches nearly relativistic intensities and contracts to a duration of just a few plasma periods. In contrast to this, the pump laser pulse has a non-relativistic intensity and long duration, so that REN and GVD effects are negligible for the pump.

Noteworthy that the sufficiently long pump might experience non-relativistic filamentation instabilities, either ponderomotive[38–40] or thermal[41]. However, these instabilities develop on longer time scales, so they can be avoided by using pump pulses having correlation times shorter than the instability times. Such pumps can be produced by using pulse randomization techniques developed to reduce nonuniformities in direct irradiation of inertially confined targets for nuclear fusion[42–44]. Similar randomization techniques can also suppress the amplified pulse transverse relativistic filamentation instability associated with REN. Thus, to assess the largest output intensity, a one-dimensional model may be adequate, with REN and GVD effects included in the equation for the amplified pulse, but not in the equation for the pump pulse.

The REN and GVD effects can also be neglected in the equation for the Langmuir wave which mediates the energy transfer from the pump to the amplified pulse. Namely, the GVD is negligible because the Langmuir wave group velocity itself is negligible, and the REN is negligible because it produces just a relatively small shift in the Langmuir wave frequency. This shift has not time enough to noticeably affect the Langmuir wave within the short duration of the amplified pulse, while it does not matter here how the Langmuir wave evolves in a given plasma location after the amplified pulse passing this location.

Furthemore, in the regimes well below the wavebreaking, which are of major interest here due to the high BRA efficiency, the hydrodynamic nonlinearity of the Langmuir wave itself may be neglected[45]. For cold enough plasmas, the Langmuir wave kinetic nonlinearity associated with trapped electrons[26,46] is also small and has not time enough to noticeably affect the Langmuir wave within the short duration of the amplified pulse.

Thus, the resulting one-dimensional equations for the resonant 3-wave interaction, taking into account the lowest order relativistic electron nonlinearity and group velocity dispersion effects for the amplified pulse, can be put in the form[13]:

$$a_t + c_a a_z = V_3 f b \, , \quad f_t = -V_3 a b^*, \tag{1}$$
$$b_t - c_b b_z = -V_3 a f^* + \imath R |b|^2 b - \imath \kappa b_{tt} \, . \tag{2}$$

Here a, b and f are envelopes of the pump pulse, counterpropagating shorter pumped pulse and resonant Langmuir wave, respectively; subscripts t and z signify time and space derivatives; c_a and c_b are group velocities of the pump and amplified pulses; V_3 is the 3-wave coupling constant (real for appropriately defined wave envelopes), R is the coefficient of nonlinear frequency shift due to the relativistic electron nonlinearity, $\kappa = c_b'/2c_b$ is the group velocity dispersion coefficient (c_b' is the derivative of the amplified pulse group velocity over the frequency).

The group velocities c_a and c_b are expressed in terms of the respective laser frequencies ω_a and ω_b as follows:

$$c_a = c\sqrt{1 - \frac{\omega_e^2}{\omega_a^2}}, \quad c_b = c\sqrt{1 - \frac{\omega_e^2}{\omega_b^2}}\, , \tag{3}$$

where c is the speed of light in vacuum,

$$\omega_e = \sqrt{\frac{4\pi n_e e^2}{m_e}} \tag{4}$$

is the electron plasma frequency, n_e is the electron plasma concentration, m_e is the electron rest mass and $-e$ is the electron charge, so that

$$2\kappa = \frac{c_b'}{c_b} = \frac{\omega_e^2}{\omega_b(\omega_b^2 - \omega_e^2)} = \frac{\omega_e^2\, c^2}{\omega_b^3\, c_b^2} \, . \tag{5}$$

The pump pulse envelope, a, is further normalized such that the average square of the electron quiver velocity in the pump laser field, measured in units of c^2, is $|a|^2$, so that

$$\overline{v_{ea}^2} = c^2 |a|^2. \tag{6}$$

Then, the average square of the electron quiver velocity in the seed laser field and in the Langmuir wave field are given by

$$\overline{v_{eb}^2} = c^2 |b|^2 \frac{\omega_a}{\omega_b}, \quad \overline{v_{ef}^2} = c^2 |f|^2 \frac{\omega_a}{\omega_f}. \tag{7}$$

The 3-wave coupling constant can be written as[47]

$$V_3 = k_f c \sqrt{\frac{\omega_e}{8\omega_b}} \tag{8}$$

where k_f is the wave number of the resonant Langmuir wave

$$k_f = k_a + k_b, \;\; k_a c = \sqrt{\omega_a^2 - \omega_e^2}, \;\; k_b c = \sqrt{\omega_b^2 - \omega_e^2}\, . \tag{9}$$

The frequency resonance condition is

$$\omega_b + \omega_f = \omega_a \, , \tag{10}$$

where $\omega_f \approx \omega_e$ is the Langmuir wave frequency in a not too hot plasma. The nonlinear frequency shift coefficient R can then be put as[48–50]

$$R = \frac{\omega_e^2 \omega_a}{4\omega_b^2}. \tag{11}$$

This hydrodynamic model is applicable for the pump pulse intensity I_0 smaller than that at the threshold of the resonant Langmuir wave breaking $I_{\rm br}$. The motivation for studying specifically such regimes is that for deep wavebreaking regimes the BRA efficiency is lower[1,3].

III. UNIVERSAL VARIABLES

The above equations will be solved for a small Gaussian initial seed and constant initial pump with a sharp front. After entering the pump depletion stage, the leading amplified spike (propagating directly behind the seed pulse) grows and contracts (since it depletes the pump faster and faster, as it grows). Thus the spike becomes of much shorter duration than the elapsed amplification time, attaining the universal features of a classical π-pulse before the REN becomes important. This prepares universal initial conditions for entering the REN regime. To expose this universality, it is helpful to change z and t variables to dimensionless variables

$$\tau = \left(1 + \frac{c_a}{c_b}\right)^{1/3} R^{1/3} V_3^{2/3} a_0^{4/3}\, \frac{L - z}{c_b}\, , \tag{12}$$
$$\zeta = \left(1 + \frac{c_a}{c_b}\right)^{-1/3} R^{-1/3} V_3^{4/3} a_0^{2/3} \left(t - \frac{L - z}{c_b}\right),$$

where τ measures the elapsed amplification time (or the distance traversed by the original seed front), ζ measures the distance (or delay time) from the original seed front; L is the plasma width and a_0 is the input pump amplitude; the seed is injected into the plasma at $z = L$, $t = 0$ and meets immediately the pump front injected into the plasma at $z = 0$, $t = -L/c_a$.

Then, defining new wave amplitudes a_1, f_1 and b_1 by formulas

$$a = a_0 a_1 \ , \tag{13}$$

$$f = -a_0 \left(1 + \frac{c_a}{c_b}\right)^{1/2} f_1 \ , \tag{14}$$

$$b = \left(\frac{V_3 a_0^2}{R}\right)^{1/3} \left(1 + \frac{c_a}{c_b}\right)^{1/6} b_1 \ , \tag{15}$$

and neglecting the "slow" time derivative of the pump amplitude compared to the "fast" time derivative of the pump amplitude, one obtains the following universal equations containing just one parameter Q:

$$a_{1\zeta} = -b_1 f_1 \ , \tag{16}$$

$$f_{1\zeta} = a_1 b_1^*, \tag{17}$$

$$b_{1\tau} = a_1 f_1^* - \imath Q b_{1\zeta\zeta} + \imath |b_1|^2 b_1, \tag{18}$$

$$Q = \frac{(k_a + k_b)^2 c^2 \omega_b c_b'}{4\omega_e \omega_a (c_a + c_b)}. \tag{19}$$

The parameter Q characterizes the group velocity dispersion of amplified pulse and depends only on the ratio of the plasma to laser frequency $q \equiv \omega_e/\omega_b$. In strongly under-critical plasmas, where $q \ll 1$, one has $Q = q/2$; in nearly critical plasmas, where $q \to 1$, one has $Q = 0.5/\sqrt{1-q^2} \gg 1$.

In strongly undercritical plasmas, which is of major interest here, the amplified pulse intensity I, fluence w and effective duration Δt_b can then be expressed in these variables as

$$I = \frac{G|b_1|^2 \omega_e}{4\lambda_b} \left(\frac{I_0^2}{2I_{\rm br}^2}\right)^{1/3} , \tag{20}$$

$$w = \frac{G\tau}{\lambda_b} \left(\frac{I_0}{4I_{\rm br}}\right)^{1/3} , \tag{21}$$

$$\Delta t_b = \frac{w}{\max_\zeta I} = \frac{4\tau}{\omega_e \max_\zeta |b_1|^2} \left(\frac{I_{\rm br}}{2I_0}\right)^{1/3} , \tag{22}$$

where

$$G = m_e^2 c^4/e^2 = 0.3\,\mathrm{J/cm}, \quad \lambda_{\rm b} = 2\pi/\mathrm{k_b}\,, \tag{23}$$

and

$$I_{\rm br} = n_e m_e c^3 q/16 \tag{24}$$

is the threshold pump intensity for resonant Langmuir wave breaking. The formula for fluence here assumes nearly complete pump depletion.

Eqs. (16)- (18) will be solved now for small input Gaussian seed pulses of the form

$$b_1(\zeta, 0) = \frac{b_{1_0}}{\sqrt{D\pi}} \exp\left[-\frac{(\zeta-\zeta_0)^2}{D}\right]$$

with $b_{1_0} = 0.05$, $D = 1$ and $\zeta_0 = 10$. No auxiliary chirping of the seed pulse is needed here, though it may be useful in less undercritical plasmas[37].

IV. DISPERSIONLESS REN REGIME

First, consider extremely undercritical plasmas where the group velocity dispersion can be neglected, so that the approximation $Q = 0$ is good enough.

Fig. 1 shows the rescaled amplified pulse amplitude $|b_1|$ as a function of the delay time ζ at several amplification times τ for $Q = 0$. The amplified pulse may have its maximum amplitude from either the first spike or from later spikes. This maximum is depicted in Fig. 2. The initial nearly linear part of the curve in Fig. 2 corresponds to the classical π-pulse regime. In what we call the REN regime, the leading spike growth saturates, while the second spike grows, reaching even higher intensity. Then the second spike growth saturates, while the third spike grows, reaching even higher intensity yet. The spikes do not filament and remain distinguishable for a while. As seen from the Fig. 2, the top amplified pulse amplitude can be nearly double the largest leading spike amplitude, so that output intensity can be nearly 4 times the leading spike theoretical limit.

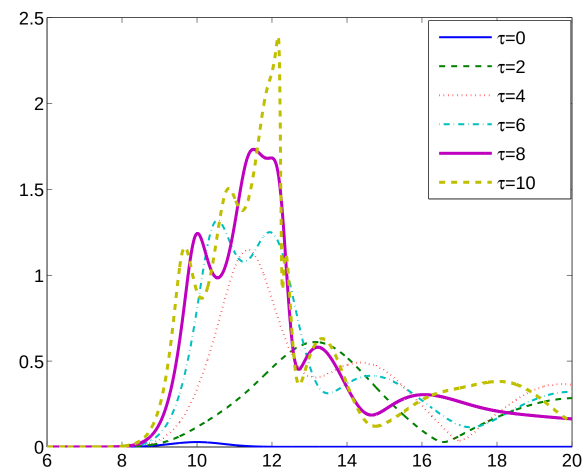

FIG. 1. The amplified pulse amplitude $|b_1|$ vs. the delay time ζ at several values of the amplification time τ in extremely undercritical plasma.

V. THE EFFECT OF GROUP VELOCITY DISPERSION

For less extreme, though still strongly undercritical plasmas, the group velocity dispersion can become important in the REN regime. This is in contrast to the π-pulse regime for which the group velocity dispersion is negligible in strongly undercritical plasmas[13]. The importance of even rather small group velocity dispersion in the REN regime is illustrated in Figs. 3 and 4 which show the dispersion effect at small $Q \approx q/2$.

Fig. 4 shows the maximum pulse amplitude $\max_\zeta |b_1|$ as a function of the amplification time τ. Note that the

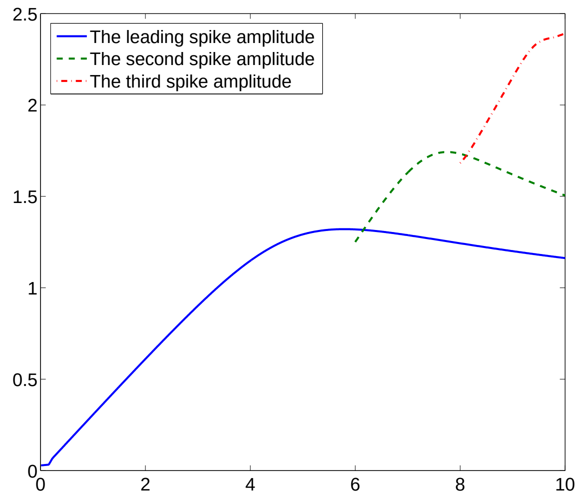

FIG. 2. The maximal amplitude of the amplified pulse $\max_\zeta |b_1|$ as a function of the amplification time τ in extremely undercritical plasma.

π-pulse regime corresponds to the joint straight part of the curves located approximately at times $\tau < 2$. Here, there is indeed no Q-dependence, indicating the negligibility of the group velocity dispersion. However, in the REN regime ($\tau > 3$), the Q-dependence becomes increasingly prominent. Larger Q corresponds to smaller pulse amplitudes, because group velocity dispersion tends to stretch the pulses. It also tends to delay the onset of the longitudinal filamentation instability, thus enabling yet larger output fluences if not intensities.

It can be seen from Figs. 3 and 4 that significant additional growth of the amplified pulse intensity and fluence can occur not only beyond the classical π-pulse regime, but even after the leading spike saturation. In extremely undercritical plasmas, where $Q \lesssim 0.01$, subsequent to the leading spike saturation, the amplified pulse intensity and fluence can increase further by a factor of about 3. In denser, but still strongly undercritical plasmas with $Q \sim 0.02 - 0.03$, the amplified pulse growth subsequent to the leading spike saturation can be about 2-fold in intensity and about 4-fold in fluence.

For example, for $\lambda_b = 1/4\,\mu$m and $I_0 = I_{\rm br}/2$, and $Q = 0.025$ (corresponding to $\omega_e/\omega_b = 0.05$), the fluence achievable in the REN regime is 120 kJ/cm^2. Here, the plasma concentration is $n_e = 4.5 \times 10^{19}\,{\rm cm}^{-3}$ and the input pump intensity is $I_0 = 1.7 \times 10^{14}\,{\rm W/cm}^2$; the pump duration is 0.7 ns, the amplified pulse output duration is 94 fs and the intensity is $1.2 \times 10^{18}\,{\rm W/cm}^2$.

Note that these intensities are more than 10 times larger than intensities reached in the recent numerical simulations[51]. One reason why the REN regime was not reached in these simulations might be because of instabilities arising from numerical noise. A number of instabilities arise from noise, whether real noise or numerical

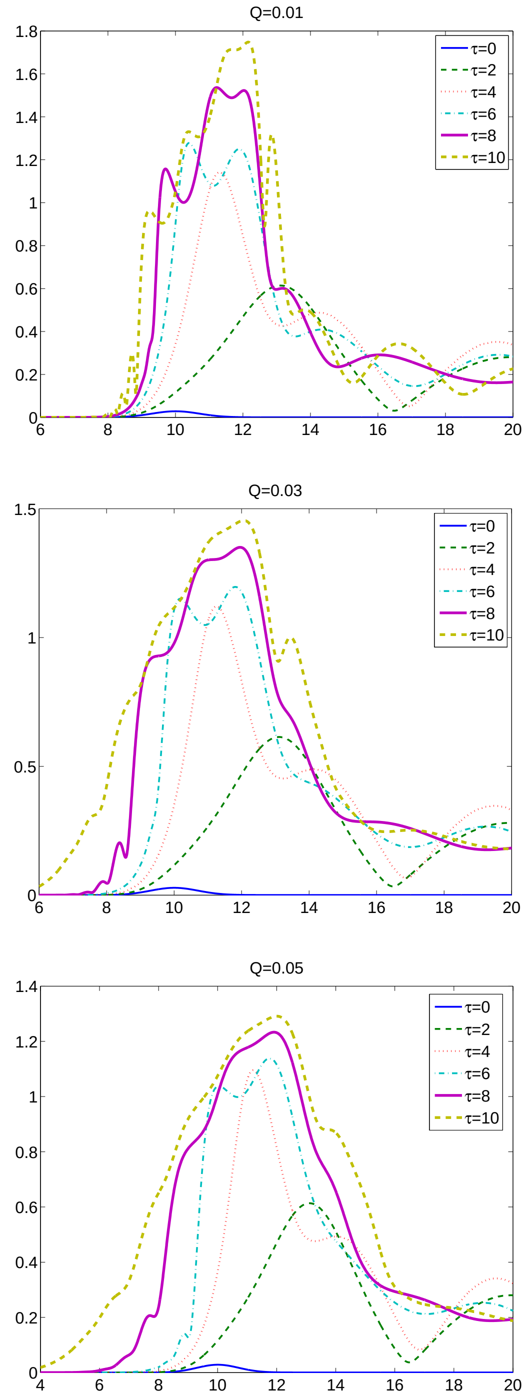

FIG. 3. The amplified pulse amplitude $|b_1|$ vs. the delay time ζ at several values of the amplification time τ and the dispersion parameter Q.

noise. The numeral noise in particle-in-cell codes might even exceed real plasma noise. In any event, the instabilities, whatever the origin, might be suppressed, say, by applying selective detuning techniques[3,4,17–19]. Since these techniques were not employed in [51], the REN regime could be unreachable. In simulations of much larger

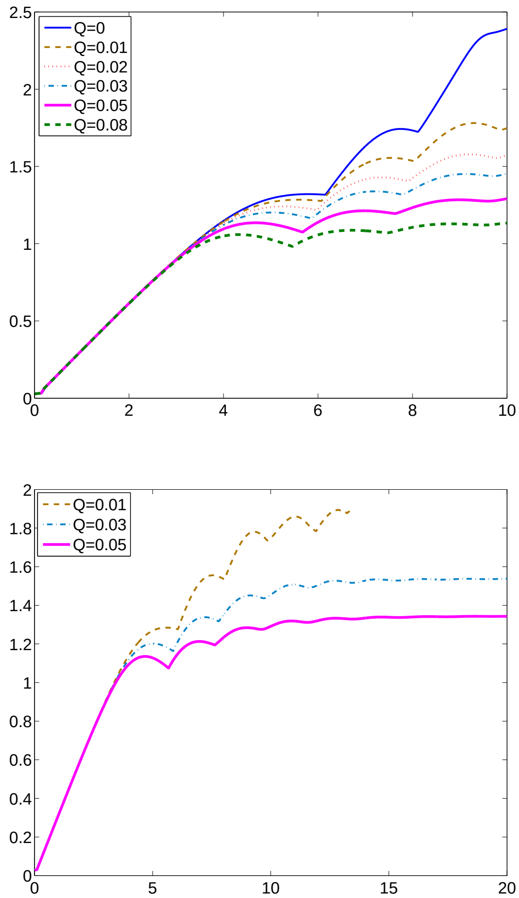

FIG. 4. The maximal amplitude of the amplified pulse $\max_\zeta |b_1|$ as a function of the amplification time τ for several values of the dispersion parameter Q.

pump intensities than discussed here, $I_0 \approx 30 I_{\rm br} \gg I_{\rm br}$, somewhat larger output intensities, like 4×10^{17} W/cm^2 , were reported[33]. Similarly, possibly because of numerical poise, the REN regime was apparently not reached also in those simulations.

Note that the ability to compress laser pulses from ns to 100 fs duration may allow direct BRA of currently available powerful 1/4 micron wavelength ns laser pulses to ultrahigh powers. The regimes found here can further enhance multi-step BRA schemes[52,53], as well as possible combinations of such schemes with other currently considered methods of producing ultra-high laser intensities, like[54–61].

VI. SUMMARY

In summary, an amplification regime is identified here wherein output pulse intensities and fluences substantially surpass the previous theoretical limit for strongly undercritical plasmas. The new intensity and fluence limits are produced by the initially sub-dominant spikes of the amplified wavetrain, which were not previously thought to be important for achieving the largest output pulses. In addition, the amplified pulse regular group velocity dispersion, in spite of being small in strongly undercritical plasmas, is shown nevertheless to be capable of delaying the pulse filamentation, thus allowing further pulse amplification to even larger output fluences.

VII. ACKNOWLEDGMENTS

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