# Quasitransient regimes of backward Raman amplification of intense x-ray pulses

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New powerful soft x-ray sources may be able to access intensities needed for backward Raman amplification (BRA) of x-ray pulses in plasmas. However, high plasma densities, needed to provide enough coupling between the pump and seed x-ray pulses, cause strong damping of the Langmuir wave that mediates energy transfer from the pump to the seed pulse. Such damping could reduce the coupling, thus making efficient BRA impossible. This work shows that efficient BRA can survive despite the Langmuir wave damping significantly exceeding the linear BRA growth rate. Moreover, the strong Langmuir wave damping can automatically suppress deleterious instabilities of BRA to the thermal noise. The class of "quasitransient" BRA regimes identified here shows that it may be feasible to observe x-ray BRA within available x-ray facilities.

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The anticipated high powers of new x-ray sources [1–3] may be sufficient to enable the ultrafast nonlinear manipulation of these pulses by means of stimulated Raman backscattering. The soft x-ray range appears to be the most promising for such experiments. The medium for Raman processing intense soft x-rays may be a thin condensed-matter sheet. The counterpropagating x-ray short seed pulse and longer pump pulse can be coupled within the sheet through the resonantly excited Langmuir wave. The optimal sheet thickness is half of the pump length. The sheet could be ionized and heated by the x-ray pulses themselves, or by an optical laser pulse when necessary, and inertially confined throughout the ultrafast nonlinear interaction. What is shown below is that such a plasma sheet can mediate Raman backscattering in a nonlinear compression regime even when the damping rate of the plasma wave exceeds significantly the linear Raman growth rate. The regimes identified here extend significantly the parameter domain in which the efficient Raman compression of ultraintense laser pulses may be achieved.

The necessary condition for the backward Raman amplification (BRA) of the seed x-ray pulse to occur within the sheet, i.e., within the half of the pump pulse duration $t_{\text{pump}}$, is

$$\Gamma_R t_{\text{pump}}/2 \gg 1, \tag{1}$$

where $\Gamma_R$ is the largest possible growth rate of the Raman instability of the seed pulse. For the transient Raman instability, this quantity is given by

$$\Gamma_R = \frac{e}{mc}\left(2\pi\frac{I\omega_e}{c\omega}\right)^{1/2}, \tag{2}$$

where $I$ is the pump intensity, $\omega$ is the pump frequency, $\omega_e$ is the plasma frequency, $\omega_e \ll \omega$, $c$ is the speed of light, $-e$ is the electron charge, and $m$ is the electron mass.

There may be, for instance, a possibility of noticeable BRA on the SASE-3 XFEL facility, anticipated to be capable of producing 100 GW pulses at wavelength $\lambda = 6.4$ nm (corresponding to the frequency $\omega = 3\times10^{17}$ s$^{-1}$) [2]. However, the assumption that the linear BRA regime is transient (i.e., that the Langmuir wave damping can be neglected within the pumped pulse duration) could be justified only for a small enough Langmuir wave damping, while the stronger damping would make BRA harder to detect.

The nonfocused output pulse diameter planned on the SASE-3 XFEL facility is 0.01 cm (corresponding to the intensity $I=10^{15}$ W/cm$^2$), and the output pulse planned duration is $t_{\text{pump}}=10^{-13}$ s. For the plasma concentration $n_e=10^{24}$ cm$^{-3}$, corresponding to the plasma frequency $\omega_e=5.6\times10^{16}$ s$^{-1}$, the linear transient BRA growth rate would be $\Gamma_R=10^{13}$ s$^{-1}$ so that $\Gamma_R t_{\text{pump}}/2=1/2$. Focusing XFEL output pulse to, say, a 0.001 cm diameter spot would produce the intensity $I=10^{17}$ W/cm$^2$, $\Gamma_R=10^{14}$ s$^{-1}$ and, respectively, $\Gamma_R t_{\text{pump}}/2=5$. This might already be sufficient for noticeable BRA, if the Langmuir wave damping indeed could be neglected. This is however not warranted, as clarified further, so that the damping impact on BRA must be taken into account.

In the presence of Langmuir wave damping $\nu$, a sufficiently short seed pulse amplified by a constant pump beam grows in the linear BRA regime proportional to $\exp[S(z,t)]$ with the exponent given at its large values, $S(z,t)\gg 1$, by the formula

$$S(z,t) \approx 2\Gamma_R\sqrt{\left(t-\frac{z}{c}\right)\frac{z}{c}} - \nu\left(t-\frac{z}{c}\right), \tag{3}$$

where $t$ is the time and $z$ is the distance in the direction of the seed pulse propagation ($0<z<ct$). This formula differs from the transient ($\nu=0$) linear BRA case (described, for instance, in [4]) just by the simple term $-\nu(t-z/c)$ describing the Langmuir wave damping behind the short seed pulse. Equation (3) can also be easily obtained from the basic equations presented below.

As seen from Eq. (3), the instability survives the Langmuir wave damping $\nu$ in the domain,

$$\frac{\tilde{\nu}^2}{1+\tilde{\nu}^2} < \frac{z}{ct} < 1, \quad \tilde{\nu} \equiv \frac{\nu}{2\Gamma_R}. \tag{4}$$

The amplified pulse maximum is located at

 

$$z = \frac{ct}{2}\left(1 + \frac{\tilde{\nu}}{\sqrt{1+\tilde{\nu}^2}}\right) \equiv z_M(t), \tag{5}$$

and

$$S(z_M,t) \approx \tilde{\Gamma}_R t, \quad \tilde{\Gamma}_R \equiv \frac{\Gamma_R}{\sqrt{1+\tilde{\nu}^2}+\tilde{\nu}}. \tag{6}$$

Thus, taking into account the Langmuir wave damping $\nu$, the necessary condition (1) for noticeable BRA should be replaced by

$$\tilde{\Gamma}_R t_{\text{pump}}/2 \gg 1. \tag{7}$$

For weak damping, $\tilde{\nu} \ll 1$, the linear BRA is transient, $\tilde{\Gamma}_R \approx \Gamma_R$, and the amplified pulse maximum moves with the speed $c/2$, i.e., twice slower than the seed so that the pulse quickly stretches.

In the opposite case of strong damping, $\tilde{\nu} \gg 1$, the linear BRA instability survives just in the narrow domain trailing the seed, $0 < t - z/c < t\tilde{\nu}^{-2}$ (assuming that the seed is shorter than even this narrow domain). Then, $t - z_M/c \approx t\tilde{\nu}^{-2}/4$ and the maximal growth rate is substantially smaller, $\tilde{\Gamma}_R \approx \Gamma_R^2/\nu \ll \Gamma_R$.

Note that in strong damping regimes, due to the pumped pulse being much shorter than the amplification distance, the above formulas can easily be generalized for varying damping $\nu$ (as long as the variation is negligible within the short pumped pulse duration). Namely, $\nu$ in formulas (3)–(6) should then just mean the Langmuir wave damping in the pumped pulse space-time location. This damping may vary significantly along the pumped pulse trajectory. As seen from Eq. (6), the linear BRA survives even for the Langmuir wave damping growing along the pumped pulse trajectory not faster than $t$. For $\nu(t,z=ct) \propto t$, the exponential linear BRA disappears. Yet, the pumped pulse amplitude still grows $\propto t$, as clarified below. The pumped pulse duration decreases in this case $\propto 1/t$. Such a distance-dependent damping can imitate nonlinear BRA saturation and pulse compression, while staying in linear BRA regimes. The effect might be also relevant to optical pulse BRA regimes, providing a possible explanation of relatively low efficiency observed so far in the experiments [5–7].

The actual damping of Langmuir wave is caused primarily by electron-ion collisions and resonant Cherenkov-type electron-wave interaction (causing so-called Landau damping),

$$\nu = \nu_{\text{cln}} + \nu_{\text{Lnd}}. \tag{8}$$

The collisional damping rate of the Langmuir wave is approximately a quarter of the electron-ion collision rate,

$$\nu_{\text{cln}} \approx \nu_{\text{ei}}/4. \tag{9}$$

For a nearly ideal classical plasma with singly charged ions, the rate of electron-ion collisions $\nu_{\text{ei}}$ can be written as

$$\nu_{\text{ei}} \approx \frac{4}{3}\sqrt{\frac{2\pi}{m}}\frac{\Lambda n_e e^4}{T_e^{3/2}}, \tag{10}$$

where $n_e$ is the electron concentration and $\Lambda$ is the Coulomb logarithm.

Landau damping is proportional to the slope of the distribution of resonant electrons moving with the phase velocity of Langmuir wave,

$$v_{\text{ph}} \approx cq_L/2, \quad q_L \equiv \omega_e/\omega \quad (\ll 1). \tag{11}$$

To avoid excessive Landau damping, it is necessary to keep the phase velocity of the Langmuir wave much larger than the thermal electron velocity $\sqrt{T_e/m}$. This is possible for sufficiently low electron plasma temperatures $T_e$ such that

$$T_e \ll mc^2 q_L^2/4 \equiv T_M. \tag{12}$$

The linear Landau damping rate for Maxwellian electron distribution is

$$\Gamma_{\text{Lnd}} = \frac{\omega_e\sqrt{\pi}}{(2q_T)^{3/2}}\exp\left(-\frac{1}{2q_T}\right), \quad q_T \equiv \frac{T_e}{T_M} \ll 1. \tag{13}$$

The actual Landau damping is sensitive to the distribution of resonant electrons and might differ significantly from Eq. (13). Deviations from Eq. (13) may occur even when the bulk electrons are Maxwellian, since resonant electrons both are more strongly affected by nonlinear interactions and are more slowly relaxed to equilibrium than are the bulk electrons. Under certain conditions, Maxwellian tails of the electron distribution might not even have time enough to form, if plasma is also heated on the ultrafast time scale of the interaction. When tails do form, Landau damping still may be much smaller than Eq. (13) since the resonant electron distribution gets smoothed within just a few plasma-wave periods in the most favorable BRA regimes where the wave amplitudes are about the wavebreaking threshold. Therefore, acceptable values of the small parameter $q_T$ might, in fact, be as large as 1/4–1/3. (Smaller values of $q_T$ are also acceptable, as long as the collisional damping not too strong, but intense x-ray pulses tend to heat plasma to higher temperatures $T_e$, closer to $T_M$, through the inverse bremsstrahlung and the Langmuir wave damping.)

Formulas (9) and (10) for collisional damping of Langmuir wave can be presented in the form

$$\nu_{\text{cln}} \approx \frac{2\sqrt{2}\Lambda r_e \omega^3}{3\sqrt{\pi} q_T^{3/2} c\omega_e}, \quad r_e \equiv \frac{e^2}{mc^2} \approx 2.818 \text{ fm}. \tag{14}$$

For the above $\lambda = 6.4$ nm ($\omega = 3\times 10^{17}$ s$^{-1}$) and $n_e = 10^{24}$ cm$^{-3}$, ($\omega_e = 5.6\times 10^{16}$ s$^{-1}$, $T_M = 4.6$ keV), and for $q_T = 1/4$ ($T_e = 1.15$ keV), it follows $\nu_{\text{cln}} = 6.4\times 10^{13}$ s$^{-1}$. Assuming that the Landau damping of Langmuir wave can be neglected due to the nonlinear suppression, modified growth rate (6), for $\Gamma_R = 10^{14}$ s$^{-1}$, would be $\tilde{\Gamma}_R = 7.3\times 10^{13}$ s$^{-1}$. Then, $\tilde{\Gamma}_R t_{\text{pump}}/2 = 3.7$ which is sufficient for the linear BRA to appear.

However, to enter the nonlinear stage of BRA, where the pumped pulse is compressed, the value of parameter $\tilde{\Gamma}_R t_{\text{pump}}/2$ should be larger. To make possible even the most advanced nonlinear compression, it would be desirable to focus the XFEL output pulse to $D=1$ $\mu$m diameter spot. This would produce the intensity $I=10^{19}$ W/cm$^2$, $\Gamma_R=10^{15}$ s$^{-1}$ and, respectively, $\Gamma_R t_{\text{pump}}/2=50$. If the Landau damping of the Langmuir wave were negligible and other parameters were the same as above, this would imply the transient BRA compression regime similar to that described in [4]. Note that, in this example, the focused pump diffraction time $D^2/(\lambda c)\sim 5\times 10^{-13}$ s is still larger than the pump duration, so that diffraction can still be neglected. The pump intensity is still below the Langmuir wave breaking threshold $I_{\text{thr}}=n_e mc^3 q_L/16=3\times 10^{19}$ W/cm$^2$, so that complete pump depletion is possible [4]. Yet, the pump intensity is already so high that the plasma heating via the inverse bremsstrahlung of pump energy may be important. Under conditions when the electron cooling by thermal conduction and by collisions with ions is negligible, the pump inverse bremsstrahlung heats the electron plasma by

$$\delta T_e = \frac{\nu_{ib} t_h I}{C_e n_e c}, \tag{15}$$

where $t_h$ is the time of heating, $C_e=3/2$ is the specific heat per electron, and $\nu_{ib}$ is the rate of inverse bremsstrahlung [8],

$$\nu_{ib} = \nu_{ei}\omega_e^2/\omega^2. \tag{16}$$

The largest heating occurs at the edge where the pump enters the plasma. There, $t_h=t_{\text{pump}}$ and $\delta T_e=1.2$ keV, so that no extra plasma heating is needed. Note that the laser pump energy per plasma electron can be evaluated as

$$T_P = \frac{2I}{n_e c}. \tag{17}$$

For the above parameters, $T_P=4$ keV, so that the energy loss due to the plasma heating is relatively small. In this transient nonlinear BRA regime, the amplified pulse can be compressed to the duration of plasma-wave period, i.e., $10^{-16}$ s, so that the nonfocused BRA output intensity can reach $4\times 10^{21}$ W/cm$^2$.

Consider now how much of the Langmuir wave damping could really be tolerated in such nonlinear BRA regimes. This is important due to the following factors:

(1) the suppression of the Landau damping might not be sufficient to neglect the damping altogether (note that linear Landau damping for a Maxwellian electron distribution at the above parameters, $\nu_{\text{Lnd}}=4\times 10^{16}$ s$^{-1}$, is close to the plasma frequency);

(2) the collisional damping might be larger in the presence of multicharged plasma ions;

(3) the collisional damping might be larger, if extra plasma heating is not provided, at smaller pump intensities, or at distances further from pump entrance in plasma; and

(4) even under conditions when the linear, or even moderate nonlinear BRA stage is strongly affected by the Langmuir wave damping, the damping might become less important during the advance nonlinear BRA stage, when the pumped pulse is strongly compressed.

To evaluate the tolerable damping, the evolution of the space-time envelopes of the pump and pumped laser pulses and resonant Langmuir wave experiencing some damping can be described by the following equations:

$$a_t - ca_z = -\Gamma_R bf,$$

$$b_t + cb_z = \Gamma_R af^*,$$

$$f_t + \nu f = \Gamma_R ab^*. \tag{18}$$

Here $a$ and $b$ are the dimensionless space-time envelopes of the pump and the pumped pulse electron quiver velocities normalized so that the power densities $I_a$ and $I_b$ are $I_b/|b|^2 \approx I_a/|a|^2 = \pi c(m_e c^2/e)^2 a_0^2/\lambda^2 = 2.736\times 10^{24} a_0^2/\lambda^2[\text{nm}]$ W/cm$^2$, and $f$ is the normalized Langmuir wave envelope.

In new variables,

$$\zeta = \Gamma_R(t - z/c), \quad \tau = \Gamma_R z/c, \tag{19}$$

Equations (18) take the form

$$2a_\zeta - a_\tau = -bf,$$

$$b_\tau = af^*,$$

$$f_\zeta + 2\tilde{\nu} f = ab^*. \tag{20}$$

Under conditions when the pumped pulse length is much shorter than the amplification distance, the term $a_\tau$ is much smaller than $a_\zeta$. Note that without the terms $a_\tau$ and also $\tilde{\nu}f$, Eqs. (20) have self-similar solutions of the form

$$b = \tau B(\eta), \quad a = A(\eta), \quad f = F(\eta), \quad \eta \equiv \tau\zeta. \tag{21}$$

which describe the pulse amplification and contraction in the well-known $\pi$-pulse regime [9]. In this regime, the ratio $f_\zeta/f \propto \tau$ grows along the pumped pulse path, so that even a growing Langmuir wave damping coefficient, $\tilde{\nu}$, could be tolerated, so long as it grows no faster than $\tau$.

The largest tolerable damping, growing like $\tau$, can be written as $2\tilde{\nu}=q_\nu\tau$ (i.e., $\nu=q_\nu\Gamma_R^2 z/c$). To evaluate the largest tolerable damping gradient, $q_\nu$, note that Eqs. (20), with $2\tilde{\nu}=q_\nu\tau$ (but without the term $a_\tau$), still have self-similar solutions of the form of Eq. (21). These solutions can be determined from the ordinary differential equations,

$$2A_\eta = -BF, \quad B + \eta B_\eta = AF^*, \quad F_\eta + q_\nu F = AB^* \tag{22}$$

with the following initial conditions (or, in the physical problem, boundary conditions) at $\eta \to +0$:

$$A(+0) = \cos(\epsilon/\sqrt{2}), \quad F(+0) = \sqrt{2}\sin(\epsilon/\sqrt{2}),$$

$$\epsilon \equiv \Gamma_R \int b(0,t)dt. \tag{23}$$

The resulting normalized pump intensity $A^2$ is shown in Fig. 1 for several values of parameters $\epsilon$ and $q_\nu$. As seen from the figure, a higher damping gradient of the Langmuir wave (large $q_\nu$) inhibits pump depletion by smaller seeds (small $\epsilon$),

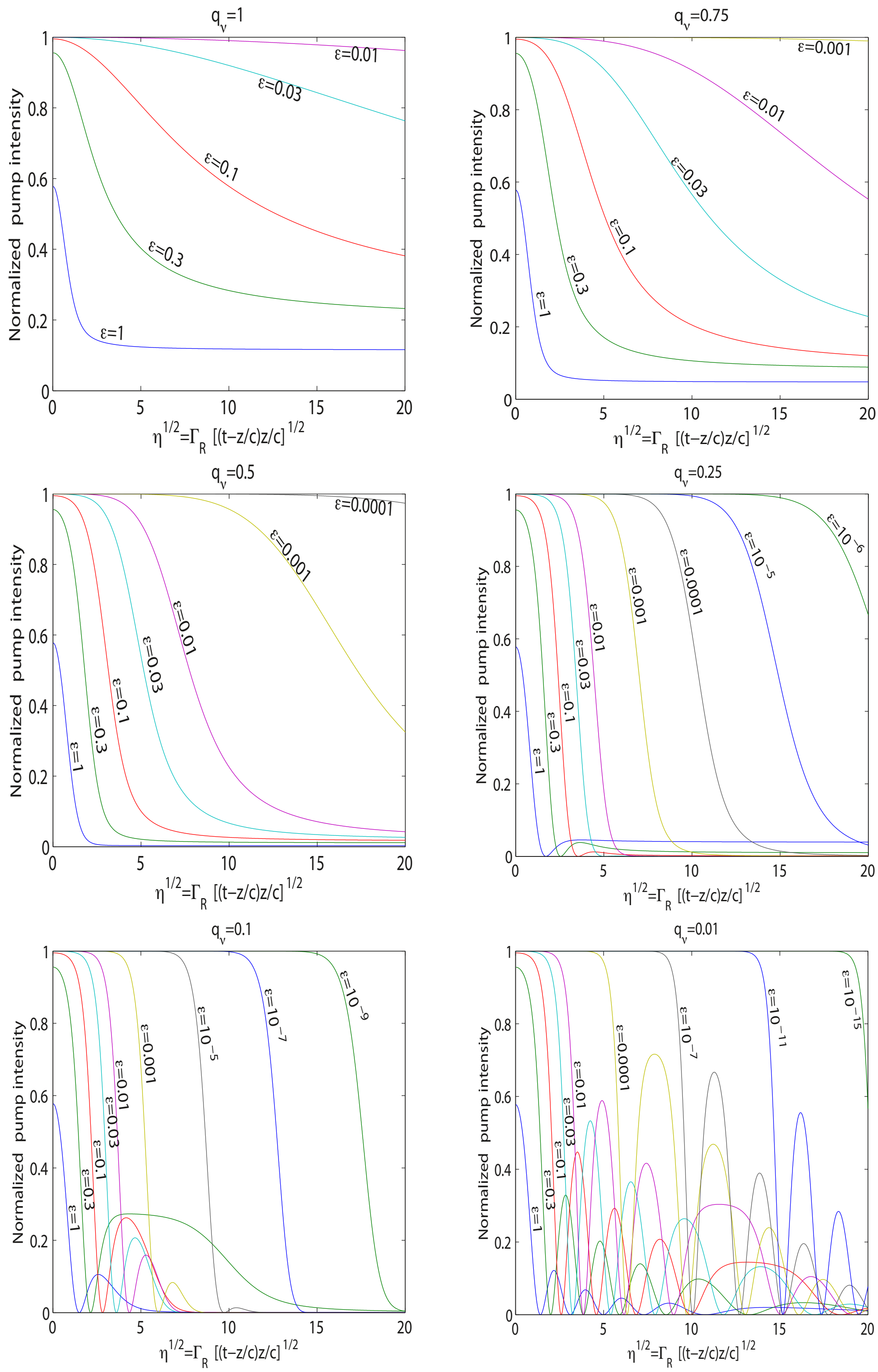


FIG. 1. (Color online) Normalized pump intensity $A^2$ in self-similar BRA regimes, for several values of the integrated seed amplitude $\epsilon$ and Langmuir wave damping gradient $q_\nu$.

while allowing pump depletion for larger seed energy. This can be an advantageous effect, since the pump propagation through the plasma is then more robust to premature backscattering off small unintentional waves, such as seeded by thermal fluctuations. An appropriate Langmuir wave damping can also advantageously suppress secondary spikes in the pumped pulse, while even improving the leading spike amplification by preventing the energy flowing back to the pump from the trailing part of the leading spike. This allows the single amplified spike to consume nearly all the pump

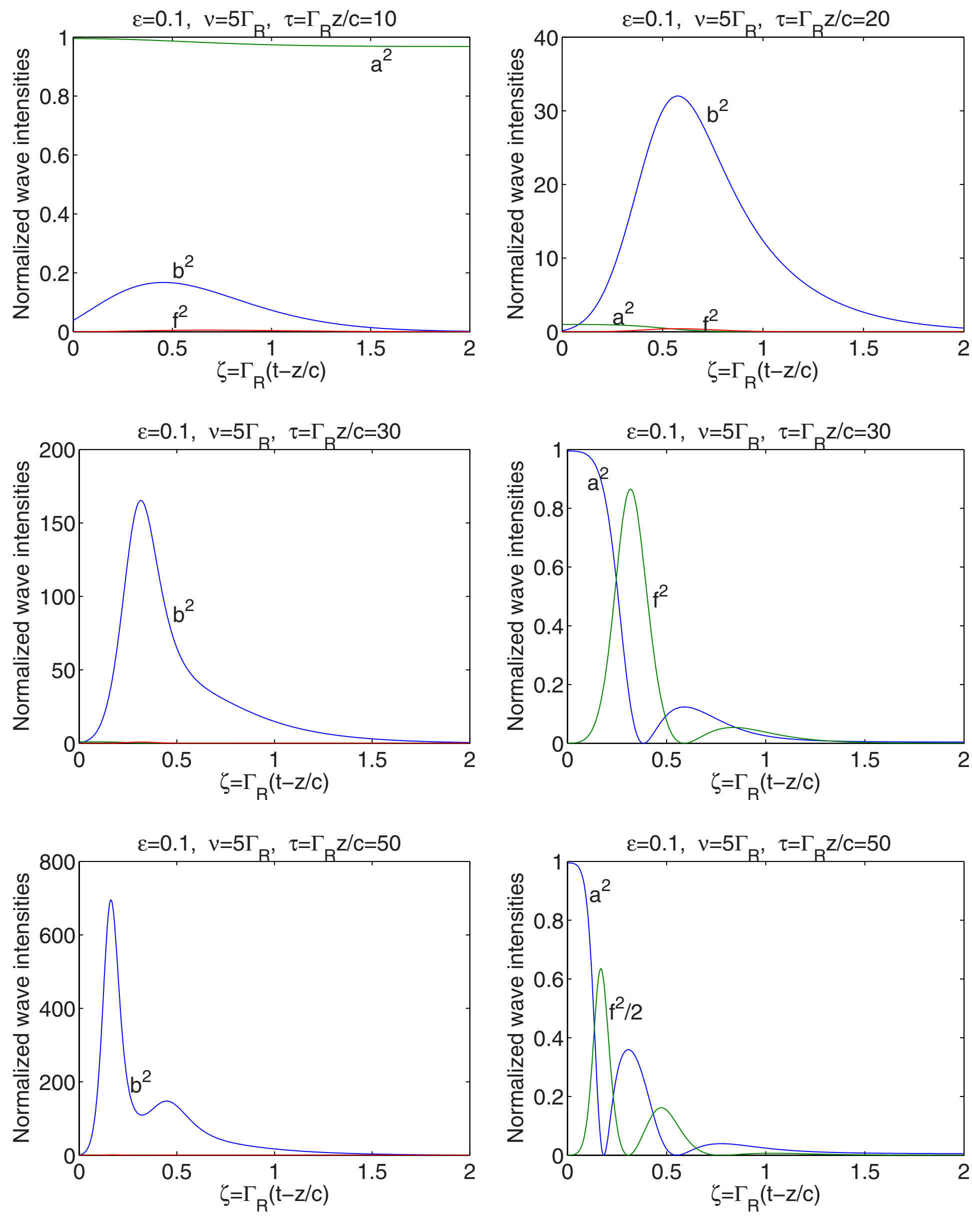


FIG. 2. (Color online) Evolution of normalized wave intensities $b^2$, $a^2$ and $f^2$ in quasitransient BRA regime with the integrated seed amplitude $\epsilon$=0.1 and Langmuir wave damping $\nu$=5$\Gamma_R$.

energy. (Note that the pump energy spent for Langmuir wave excitation constitutes just a small fraction, $\omega_e/\omega \ll 1$, of the energy consumed by the pumped pulse.)

Most importantly, these results appear not to be very sensitive to details of the Langmuir wave damping variation with the amplification distance. What really matters is the damping value at the final stage of amplification when the most of the pump energy is consumed by the amplified pulse (this is because the distance needed for pumped pulse amplification grows proportional the pulse energy). The maximum tolerable damping value at the final BRA stage can be expressed in the terms of the parameter $q_\nu$ as

$$\nu_f \sim q_\nu \Gamma_R^2 t_{\text{pump}}/2.$$

For $q_\nu$=0.5, $\Gamma_R$=$10^{15}$ s$^{-1}$ and $\Gamma_R t_{\text{pump}}/2$=50, it follows that $\nu_f \sim 2.5\times 10^{16}$ s$^{-1}$, so that Langmuir wave damping as large as about half of the plasma frequency could be tolerated.

These conclusions are supported by numerical solution of Eq. (20) (without the term $a_\tau$) with initial and boundary conditions

$$a(+0,\tau) = \cos(\epsilon/\sqrt{2}), \quad f(+0,\tau) = \sqrt{2}\,\sin(\epsilon/\sqrt{2}),$$

$$b(\zeta,0) = 0. \tag{24}$$

Normalized wave intensities $b^2$, $a^2$, and $f^2$ in quasitransient BRA regimes are shown in Figs. 2–4 for several values of the amplification distance, integrated seed amplitude $\epsilon$ and Langmuir wave damping $\nu$. Here and further, constant damping rate $\nu$ is considered (rather than a damping rate that grows with $\tau$).

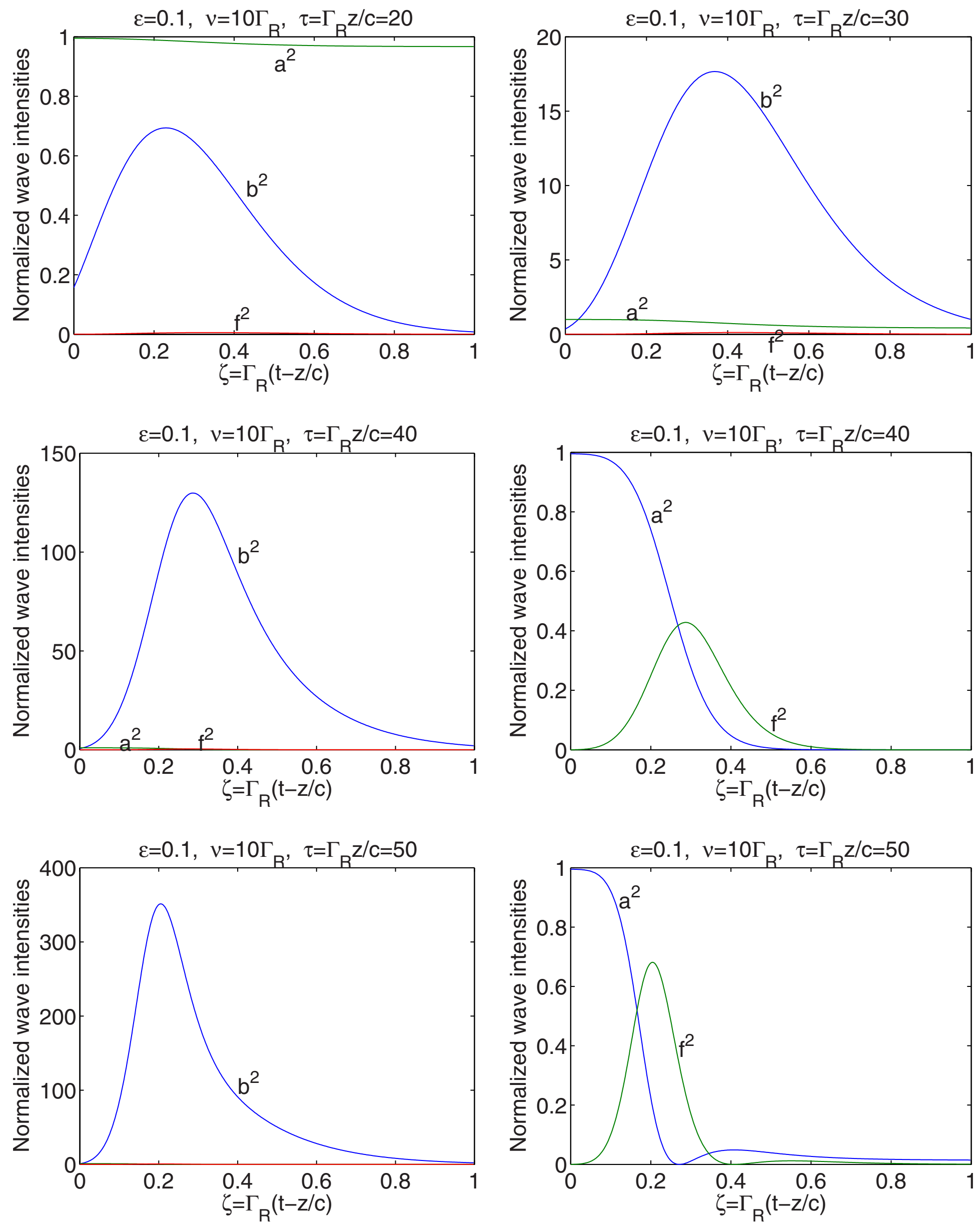


FIG. 3. (Color online) Evolution of normalized wave intensities $b^2$, $a^2$ and $f^2$ in quasitransient BRA regime with the integrated seed amplitude $\epsilon=0.1$ and Langmuir wave damping $\nu=10\Gamma_R$.

These figures show that significant Langmuir wave damping indeed can be tolerated in quasitransient BRA regimes, with nearly total pump depletion even when the Langmuir wave damping rate significantly exceeds the largest linear Raman growth rate. The largest tolerable damping is an increasing function of the amplification distance and integrated seed pulse amplitude. When the damping exceeds this limit, the pump depletion decreases. Yet even in such nearly linear BRA regimes, significant amplification and compression of the pumped pulse can occur.

More specifically, Fig. 2 shows the pulse evolution in quasitransient regimes with Langmuir wave damping moderately large compared to the largest linear Raman growth rate, $\nu=5\Gamma_R$. In such regimes, even an initially small enough amplified pulse, $\epsilon=0.1$, becomes, at an advanced nonlinear BRA stage, not very sensitive to the damping and resembling the leading amplified spike in purely transient BRA regimes. At the latest nonlinear BRA stage even the second amplified spike appears. The advanced amplified pulse intensity $b^2$ greatly exceeds the pump intensity $a^2$ so that the latter is barely seen at the $b^2$ plot for $\tau=\Gamma_R z/c=30$ (or 50). Therefore, the pump and Langmuir wave intensities, $a^2$ and $f^2$, are shown separately, for the same parameters, in the expanded scale plots.

A stronger damping, $\nu=10\Gamma_R$, significantly prolongs the linear stage of amplification of the pulse of the same initial integrated amplitude, $\epsilon=0.1$, and reduces the amplified pulse intensity even at an advanced nonlinear stage. This is clearly seen from comparison of plots presented in Figs. 2 and 3. On each plot, the major parameters are indicated. The amplified pulse intensity reduction becomes less pronounced, but still remains even at the final BRA stage, as seen from plots Fig.

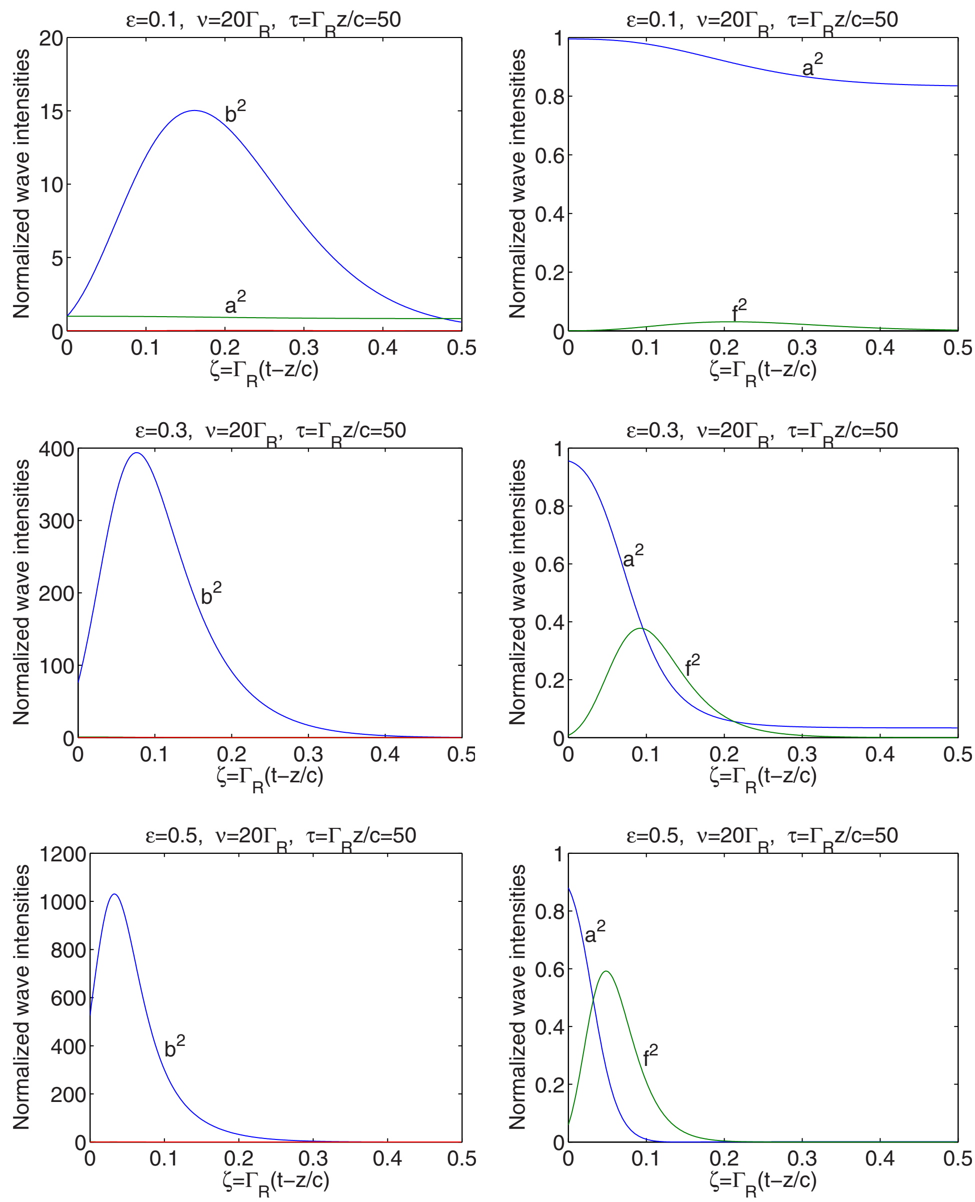


FIG. 4. (Color online) Final normalized wave intensities $b^2$, $a^2$ and $f^2$ in quasitransient BRA regime with the integrated seed amplitude $\epsilon=0.1;0.3;0.5$ and Langmuir wave damping $\nu=20\Gamma_R$.

2 and 3 at $\tau=\Gamma_R z/c=50$. Due to the amplified pulse intensity reduction, the pump intensity $a^2$ remains slightly visible at the $b^2$ plot for $\tau=\Gamma_R z/c=40$ (and is even marked therein). To show the pump and Langmuir wave intensities, $a^2$ and $f^2$, properly, the separate expanded scale plot is presented for $\tau=\Gamma_R z/c=40$ (and similarly for $\tau=\Gamma_R z/c=50$).

Even stronger damping, $\nu=20\Gamma_R$, delays the linear stage of amplification of a pulse with the same initial integrated amplitude, $\epsilon=0.1$, so much so that the pump depletion remains small even at the final BRA stage $\tau=\Gamma_R z/c=50$, as seen from the respective plots in the Fig. 4. Due to the great reduction of the amplified pulse intensity, the pump intensity $a^2$ is clearly seen even at the $b^2$ plot for $\epsilon=0.1$. However, to clearly show the small depletion of the pump intensity and small Langmuir wave intensity, the expanded scale plot is presented separately for $\epsilon=0.1$.

Even at such stronger damping, $\nu=20\Gamma_R$, the BRA efficiency and pulse compression can be improved significantly by using stronger seed pulses. The improvement is apparent in Fig. 4 for a larger initial integrated amplitude, $\epsilon=0.3$. The amplified pulse intensity $b^2$ greatly exceeds the pump intensity $a^2$, so that the latter is again barely seen at the $b^2$ plot for $\epsilon=0.3$. Therefore, the pump and Langmuir wave intensities, $a^2$ and $f^2$, are shown separately, for the same parameters, in the expanded scale plot. This plot shows, in particular, that the initial seed integrated amplitude $\epsilon=0.3$ produces nearly complete pump depletion at the final BRA stage.

For an even larger initial seed integrated amplitude, $\epsilon=0.5$, the pulse compression at the final BRA stage is essentially restored, and the amplified pulse becomes about 1000 times more intense than the input pump. The expanded scale plot for the final pump and Langmuir wave intensities, $a^2$ and $f^2$ at $\epsilon=0.5$, shows, in particular, that the pump depletion reaches 100%.

For the parameters considered above for the advanced nonlinear regime, the resonant electrons are not thermalized within the output pumped pulse duration (though the Maxwellization of the resonant electrons has time enough to occur within the pump duration). This permits the Langmuir waves to smooth the electron distribution, and thus to reduce the Langmuir wave Landau damping below that given by Eq. (13). The reduction by just a half already enables efficient quasitransient BRA regimes.

While these results are general enough to apply to any Raman backscatter regimes, including optical ones, the quasitransient regimes are of the most crucial importance specifically for x-ray BRA. This is because the x-ray BRA needs plasma densities as large as those of condensed matter in order to provide enough coupling between the pump and seed laser pulses, thus making strong damping of Langmuir waves inevitable. The above theory extends the parameter range for efficient BRA. It makes more feasible approaching the theoretical short-wavelength limit for efficient x-ray BRA found in [10], as well as exploiting relic lattice effects to enhance the dispersion of the x-ray group velocity and exceed the theoretical limit for the output pulse fluence in homogeneous plasma, as proposed in [11].

The major results of this work are:

(i) it is shown that efficient backward Raman amplification (BRA) is possible even for the Langmuir wave damping significantly exceeding the linear Raman growth rate, and the largest tolerable damping is determined quantitatively;

(ii) moreover, the strong Langmuir wave damping is shown to be capable of automatically suppressing deleterious instabilities of BRA to the thermal noise; and

(iii) a class of BRA regimes is identified, called herein "quasitransient" regimes, in which the observation of x-ray BRA may be feasible within already available technologies, in particular, on the currently built SASE-3 XFEL facility.

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[1] R. Tatchyn *et al.*, Nucl. Instrum. Methods A **375**, 274 (1996); www-ssrl.slac.stanford.edu/lcls
[2] V. Ayvazyan *et al.*, Eur. Phys. J. D **37**, 297 (2006); SLAC-PUB-12114; Technical Design TESLA XFEL 2002–09 Report (DESY 2002–167), Hamburg, Germany, 2002.
[3] W. Barletta *et al.*, R. Nuovo Cim. Soc. Ital. Fis. **29**, 1 (2006); www.elettra.trieste.it/FERMI
[4] V. M. Malkin, G. Shvets, and N. J. Fisch, Phys. Rev. Lett. **82**, 4448 (1999).
[5] W. Cheng, Y. Avitzour, Y. Ping, S. Suckewer, N. J. Fisch, M. S. Hur, and J. S. Wurtele, Phys. Rev. Lett. **94**, 045003 (2005).
[6] J. Ren *et al.*, Nat. Phys. **3**, 732 (2007).
[7] J. Ren *et al.*, Phys. Plasmas **15**, 056702 (2008).
[8] I. P. Shkarofsky, T. W. Johnston, and M. P. Bachynsky, *The Particle Kinetics of Plasmas* (Addison-Wesley, Reading, MA, 1966), p. 258.
[9] G. L. Lamb, Jr., Phys. Lett. A **29**, 507 (1969); Rev. Mod. Phys. **43**, 99 (1971); *Elements of Soliton Theory* (Wiley, New York, 1980); V. E. Zakharov, JETP Lett. **32**, 589 (1980); S. V. Manakov, *ibid.* **35**, 237 (1982); V. A. Gorbunov, V. B. Ivanov, S. B. Papernyi, and V. R. Startsev Izv. Akad. Nauk SSSR, Ser. Fiz. **48**, 1580 (1984) [Bull. Acad. Sci. USSR, Phys. Ser. (Engl. Transl.) **48**, 120 (1984)]; J. Coste and C. Montes, Phys. Rev. A **34**, 3940 (1986).
[10] V. M. Malkin, N. J. Fisch, and J. S. Wurtele, Phys. Rev. E **75**, 026404 (2007).
[11] V. M. Malkin and N. J. Fisch, Phys. Rev. Lett. **99**, 205001 (2007).