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Review

# Key plasma parameters for resonant backward Raman amplification in plasma

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**Abstract.** Backward Raman amplification and compression in plasma enables pulse compression to intensities not available using material gratings. In order to achieve the highest intensities and efficiencies in the compression effect, in a manner robust both to noise and other competing plasma effects, both resonance effects and detuning effects are exploited. Here we offer a simplified guide to how some of the key plasma parameters and laser parameters should be picked in order to achieve robust and efficient amplification.

## 1 Introduction

It is hoped, by using resonant backward Raman amplification in plasmas, to achieve the next generation of laser intensities [1]. This method contemplates a counter-propagating geometry, where the energy is stored in a long pump pulse, or possibly a train of pump pulses. The energy of this pump is then depleted by a short, Raman-downshifted, counter-propagating seed pulse. Under Raman amplification, the seed pulse grows to intensities far higher than the pump pulse, effectively capturing the pump pulse energy in a pulse of shorter duration, or, effectively, achieving what might be called Raman compression. The limiting effects in plasma are the nonlinear effects associated with the nearly relativistic electron velocities in the wave fields, but these limits arise at intensities far beyond those tolerated by the material gratings needed for the current state-of-the-art technique of chirped pulse amplification (CPA) [2].

Material gratings apparently cannot tolerate laser pulses so intense that the electron quiver energy reaches the material ionization energy. In contrast, the limiting quiver energy for Raman amplification sets in only when the electrons reach relativistic energies, which is typically larger than material ionization energies by a factor of $10^5$ to $10^6$. Thus, in principle, plasma can mediate intensities larger than material grating by a factor $10^5$ to $10^6$.

For a given electron quiver energy, the laser pulse intensity is inversely proportional to the square of the laser wavelength, so, in principle, at shorter wavelengths higher intensities could be reached both for Raman amplification techniques in plasma as well as for CPA techniques. The shortest laser wavelength that could, in practice, be

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handled by material gratings using CPA is about 1/3 micron. This sets the largest tolerable intensity on the gratings to about $10^{12}$ to $10^{13}$ W/cm$^2$.

To give a feel for the opportunities in plasma, at a wavelength of 1/3 micron, unfocused laser intensities of $10^{18}$ W/cm$^2$ might be reached through backward Raman amplification. If produced with apertures of just 1 cm$^2$, the laser pulse would reach powers of an exawatt ($10^{18}$ W); if produced with apertures of $30 \times 30$ cm$^2$, zettawatt powers ($10^{21}$ W) could be reached. Moreover, in contrast to CPA, wavelengths much shorter than 1/3 micron can be accommodated in plasma, so that even much higher powers, in principle, could be reached.

Note that resonant backward Raman amplification seeded by noise is considered to be one of the fastest and often the most disruptive of the laser instabilities in plasma. It is, for example, a concern in delivering energy to the highly compressed pellet of plasma in inertial confinement fusion. In the application considered here, however, the same fast process, can be used constructively, allowing a well-prepared counter-propagating laser seed to consume the energy in a powerful pump beam faster than any of the competing parasitic processes.

Significant theoretical effort has been exerted in the past several years to optimize the efficiency and to make the effect robust to non-ideal effects, such as plasma noise or imperfect laser pulses. At the same time, many of the key ideas have been implemented in experiments, where some of the important uncertainties were retired and some of the basic elements of the nonlinear regime were demonstrated [3–14]. The effects that were experimentally demonstrated include the theoretical predictions of advantageously detuning the 3-wave resonance by chirping the pump together with employing gradients in the plasma density [15,16].

Here we provide a simplified guide to some of the key plasma parameters and key laser parameters that improve the chances of achieving robust amplification. We try to provide dimensional parameters, in practical units, to facilitate the implementation of the theoretical ideas. The guide is by no means complete, but hopefully it is useful.

## 2 Basic equations for resonant Raman amplification

Backward Raman amplification is, in general, a process that transfers energy from a pump laser pulse (of frequency $\omega_a$) to a counter-propagating seed laser pulse of lower frequency $\omega_b < \omega_a$ through an electron excitation of the medium. In a much narrower sense, it is the resonant 3-wave decay process wherein a pump photon decays into a counter-propagating seed photon and Langmuir plasmon. The narrower sense is the one used here, where Raman amplification describes the process mediated by a plasma for high laser intensity applications. The plasma medium is essential to our use here, because, as opposed to other Raman media such as gases or fibers, the plasma is a medium already broken down ionized, and so can withstand extreme intensities.

The resonance conditions (or the quantum energy and momentum conservation laws) are

$$\omega_a = \omega_b + \omega_e, \tag{1}$$

$$k_a + k_b = k_e, \tag{2}$$

where $\omega_e$ is the Langmuir wave frequency, $k_{a,b,e} = 2\pi/\lambda_{a,b,e}$ are respective wavenumbers, and $\lambda_{a,b,e}$ are respective wavelengths in plasma. Pump and seed laser frequencies $\omega_a$ and $\omega_b$ are expressed in terms of respective laser wavelengths in vacuum $\lambda_{a0}$ and $\lambda_{b0}$ as follows

$$\omega_{a,b} = \frac{2\pi c}{\lambda_{a0,b0}} \approx \frac{1.9 \times 10^{11}}{\lambda_{a0,b0}/\text{cm}}\,\text{sec}^{-1},$$

where $c$ is speed of light in vacuum.

The Langmuir wave frequency $\omega_e$ is close to the electron plasma frequency $\omega_0$,

$$\omega_e \approx \omega_0 = \sqrt{\frac{4\pi n_e e^2}{m_e}} = 5.6 \times 10^4 \sqrt{n_e\,\mathrm{cm}^3}\,\mathrm{sec}^{-1},$$

where $n_e$ is the plasma electron concentration, $e$ is the electron charge and $m_e$ is the electron mass. Laser pulses can propagate in plasmas only of plasma frequencies smaller than laser frequency. The critical plasma concentration $n_{cr}$ for the seed laser pulse propagation is such that

$$\frac{4\pi e^2 n_{cr}}{m_e} = \omega_b^2.$$

It follows that

$$n_{cr} = \frac{\pi m_e c^2}{e^2 \lambda_{b0}^2} \approx \frac{1.1 \times 10^{13}}{\lambda_{b0}^2/\mathrm{cm}^2}\,\mathrm{cm}^{-3}.$$

In strongly under-critical plasmas, $n_e \ll n_{cr}$, where

$$\omega_0 = \omega_b \sqrt{\frac{n_e}{n_{cr}}} \ll \omega_b,$$

one has $\lambda_{a,b} \approx \lambda_{a0,b0}$ and

$$k_a \approx k_b \approx k_e/2 \approx \omega_a/c.$$

## 3 Langmuir wave breaking threshold

In the above resonant 3-wave decay process, the Langmuir wave acquires the fraction $\omega_e/\omega_a$ of energy lost by the laser pump. This acquisition occurs, at a given spatial location, within the short time equal to the duration of the amplified pulse as it crosses that location in the plasma. Assuming that no damping occurs within such a short time, the Langmuir wave holds the acquired energy within this time. To the extent that the Langmuir wave holds this energy, it then facilitates the transfer of pump energy to the backward propagating seed.

The energy density of Langmuir wave that can be sustained by the plasma is limited, however, by the phenomena of wave breaking. The Langmuir wave breaking occurs when the amplitude of the electron quiver velocity in the wave reaches the wave phase velocity

$$v_{\mathrm{ph}} = \frac{\omega_e}{k_e} \approx \frac{c\omega_e}{2\omega_a}.$$

The Langmuir wave energy density at wavebreaking is

$$W_{\mathrm{ebr}} = n_e m_e v_{\mathrm{ph}}^2/2.$$

The laser pump energy density at wavebreaking is

$$W_{\mathrm{br}} = W_{\mathrm{ebr}} \frac{\omega_a}{2\omega_e}.$$

The factor 2 in the denominator here is due to the fact that the pump slice losing energy within a given plasma layer is twice thicker than this layer.

Thus, the pump intensity at the Langmuir wave breaking threshold can be put as

$$I_{\mathrm{br}} = cW_{\mathrm{br}} \approx \frac{n_e m_e c^3 \omega_e}{16\,\omega_a} = I_M \left(\frac{n_e}{n_{cr}}\right)^{3/2}, \tag{3}$$

where

$$I_M \approx \frac{n_{cr} m_e c^3}{16} \approx \frac{\pi m_e^2 c^5}{16 e^2 \lambda_a^2} \approx \frac{1.7 \times 10^9}{\lambda_a^2/\mathrm{cm}^2} \frac{\mathrm{W}}{\mathrm{cm}^2} \,. \tag{4}$$

For pump laser pulses of intensities exceeding the wave-breaking threshold, $I_0 > I_{\mathrm{br}}$, and considering only the standard resonant 3-wave decay process, only a fraction $I_{\mathrm{br}}/I_0$ of the pump energy would be depleted [1]. Thus, our primary attention is given to backward Raman amplification regimes where $I_0 < I_{\mathrm{br}}$, so that there is no Langmuir wave breaking. In these regimes, the hydrodynamic or fluid description is applicable.

## 4 Linear regime for backward Raman decay instability

The linear regime of resonant Raman amplification is the regime in which the pump depletion is negligible. The linear growth rate for the Stokes ($\omega_b = \omega_a - \omega_e$) wave amplitude during the standard 3-wave backward Raman decay instability of a monochromatic pump wave is

$$\gamma_0 = \sqrt{\frac{\omega_a \omega_e \overline{v_a^2}}{2c^2}}.$$

The averaged over the wavelength square of the electron quiver velocity in this wave, $\overline{v_a^2}$, is related to the pump intensity $I_0$ as

$$\overline{v_a^2} = \frac{4\pi e^2 I_0}{m_e^2 \omega_a^2 c} = \frac{c^2 I_0}{16 I_M}.$$

It follows that

$$\gamma_0 = \sqrt{\frac{\omega_a \omega_e I_0}{32\, I_M}} \,. \tag{5}$$

In regimes in which wavebreaking is absent, namely $I_0 < I_{\mathrm{br}} = I_M(\omega_e/\omega_a)^3$, it follows from Eq. (5) that $\gamma_0 \ll \omega_e$. The bandwidth of the resonance does not overlap the anti-Stokes wave at $\omega_b^+ = \omega_a + \omega_e$. The instability can therefore be described by the standard 3-wave backward Raman decay. When a long pump pulse with a sharp front encounters a short, small counter-propagating seed pulse at the Stokes resonant frequency $\omega_b = \omega_a - \omega_e$, the Stokes wave amplitude grows behind the seed proportionally to

$$\exp\left[2\gamma_0 \sqrt{(t - z/c)z/c}\right].$$

Here $t$ is the time passed after the seed (propagating in the positive direction $z$ with the speed $c$) meets the pump front at $z = 0$. The Stokes pulse maximum, located at $z = ct/2$, propagates with the speed $c/2$ and grows with the rate $\gamma_0$. Since the front propagates at $c$ while the maximum propagates only at $c/2$, the linear regime displays a broadening of the seed pulse.

The linear regime is valid so long as the pump does not undergo significant depletion, and hence may be regarded as approximately constant.

## 5 Pump depletion regime

The linear amplification stage ends when the pump depletion becomes significant. This occurs first near the maximum of the amplified pulse. Then the rear layers of amplified pulse (no longer reached by the pump) are shaded and stop growing, while ahead of the maximum (where the pump depletion is still small) the amplification persists as in the linear instability regime. As a result, the pulse reshapes in such a way that its maximum relocates closer to the leading original small seed. That original seed propagates nearly with the speed $c$, essentially unchanged by the interaction. Thus, the amplified pulse maximum propagates with a super-luminous velocity. As the maximum grows. it thus contracts and becomes entirely localized near the region $z \simeq ct$.

In this regime, the growth is no longer exponential; it is much slower than exponential. Because it is slower than exponential, the number of exponentiations occurring between the original small seed and the amplified pulse maximum remains approximately the same, roughly equal to $\Lambda_0$,

$$\Lambda_0 \approx 2\gamma_0 \sqrt{(t - z_M/c)t}\,, \tag{6}$$

where $z_M$ is the spatial location at which the seed reaches its maximum intensity. It follows from Eq. (6) that

$$z_M \approx c\left(t - \frac{\Lambda_0^2}{4\gamma_0^2 t}\right).$$

The duration of the amplified spike decreases inversely proportional to the time $t$, while the amplitude increases proportional to $t$. The spike evolves, in fact, in a self-similar way. The spike intensity in the energy containing region around the maximum can be roughly approximated by the formula

$$I_b \approx 2I_0 \left\{ \frac{2\gamma_0 t}{\Lambda_0 \cosh\left[2\gamma_0^2 t(z - z_M)/\Lambda_0 c\right]} \right\}^2.$$

The fluence of this leading spike is

$$w_{b1} \approx 8I_0 t/\Lambda_0,$$

which constitutes $4/\Lambda_0$ fraction of the encountered pump fluence $w_a = 2I_0 t$.

This formula takes into account just the first term of the expansion in the small parameter $1/\Lambda_0 \ll 1$. A more accurate calculation [1] gives

$$\frac{w_{b1}}{w_a} \approx \frac{4}{\Lambda_0 + 2}.$$

The rest of the pump energy goes to the next spikes. The duration of the leading spike is

$$\Delta t_b = \frac{w_{b1}}{\max_z I_b} \approx \frac{\Lambda_0}{\gamma_0^2 t}\,. \tag{7}$$

All these formulas are applicable for the regime in the pump depletion regime only, namely for $\gamma_0 t > \Lambda_0$, when $\gamma_0 \Delta t_b < 1$. The duration of the original seed pulse should not exceed $\Delta t_b$, so that the entire original seed remains ahead of the amplified spike. This gives the condition that the bandwidth of the original seed pulse should be no smaller than the bandwidth of the anticipated output pulse.

## 6 Relativistic electron nonlinearity

The amplification in the pump depletion regime is limited by the relativistic electron nonlinearity of the leading spike. The relativistic increase of the electron mass in the spike reduces local plasma frequency, thereby producing the frequency shift

$$\delta\omega_{nl} = \frac{\omega_e^2 \overline{v_b^2}}{4\omega_b c^2},$$

where

$$\overline{v_b^2} = \frac{4\pi e^2 I_b}{m_e^2 \omega_b^2 c} \approx \frac{c^2 I_b}{16 I_M}.$$

is the square of the electron quiver velocity in the amplified pulse, averaged over a wavelength. Using the above formulas, the nonlinear frequency shift at the top of amplified pulse can be put in the form

$$\delta\omega_{nl} \approx \omega_e^3 t^2 \left(\frac{I_0}{16 I_M \Lambda_0}\right)^2,$$

with corresponding cumulative nonlinear phase shift

$$\delta\phi_{nl} = \int dt \delta\omega_{nl} \approx \frac{\omega_e^3 t^3}{3} \left(\frac{I_0}{16 I_M \Lambda_0}\right)^2.$$

The phase in the 3-wave interaction governs whether energy flows from the pump to the seed pulse, or whether the energy flows in the opposite direction. Thus, should the phase shift exceed $\delta \sim 1$, it could result in a significant slowing down (or even reversal) of the longitudinal amplification. A phase shift of that magnitude might also cause transverse filamentation of the amplified pulse. Thus, the amplification time should be kept short enough to avoid such a large phase change, namely below the maximum amplification time $t_M$

$$t < t_M = \frac{1}{\omega_e} \left[3\delta \left(\frac{16 I_M \Lambda_0}{I_0}\right)^2\right]^{1/3}.$$

The maximum amplification time, in turn, determines the shortest achievable duration of the leading amplified spike

$$\Delta t_b > \Delta t_{bm} = \frac{4}{\omega_a} \left(\frac{2\Lambda_0 I_M}{3\delta I_0}\right)^{1/3}, \quad (8)$$

as well as the largest achievable fluence,

$$w_b < w_{bM} = \frac{64 I_M}{\omega_e} \left(\frac{3\delta I_0}{2\Lambda_0 I_M}\right)^{1/3}, \quad (9)$$

and the largest achievable intensity,

$$I_b < I_{bM} = \frac{w_{bM}}{\Delta t_{bm}}. \quad (10)$$

Note that, for given $\omega_a$, $I_0$, $\Lambda_0$ and $\delta$, the shortest achievable duration of the leading amplified spike does not depend on the plasma concentration, whereas the largest achievable fluence and intensity are inversely proportional to $\omega_e \propto \sqrt{n_e}$. This applies just as long as the plasma is dense enough to avoid the Langmuir wave breaking, i.e. for $\omega_e > \omega_a (I_0/I_M)^{1/3}$.

The greatest output pulse fluences and intensities are achieved in the plasma of lowest possible density still in the hydrodynamic limit, namely the wavebreaking density (when $\omega_e = \omega_a (I_0/I_M)^{1/3}$). Maximized over the plasma density, the maximum fluence can then be put as

$$\max_{n_e} w_{bM} = \frac{64 I_M}{\omega_a} \left( \frac{3\delta}{2\Lambda_0} \right)^{1/3} , \tag{11}$$

and the maximum intensity can then be put as

$$\max_{n_e} I_{bM} = 16 \left( I_M^2 I_0 \right)^{1/3} \left( \frac{3\delta}{2\Lambda_0} \right)^{2/3} . \tag{12}$$

Note that the maximal fluence here does not depend on the pump intensity $I_0$. The factor $64 I_M/\omega_a$ in the formula for fluence can also be presented in the form

$$\frac{64 I_M}{\omega_a} = \frac{2 m_e^2 c^4}{e^2 \lambda_a} \approx \frac{0.6}{\lambda_\mathrm{a}/\mathrm{cm}} \ \mathrm{J/cm}^2 . \tag{13}$$

## 7 Pump backscattering by noise

Within the amplification time $t_M$, there could occur a large number of the linear instability exponentiations,

$$\Lambda_M = \gamma_0 t_M = \sqrt{\frac{\omega_a}{\omega_e}} \left( 3\delta\, \Lambda_0^2 \right)^{1/3} \left( \frac{2 I_M}{I_0} \right)^{1/6} ,$$

particularly in strongly under-critical plasmas, where $\omega_a \gg \omega_e$. If this number exceeds the number of exponentiations needed for the plasma noise to grow and deplete the pump, $\Lambda_M > \Lambda_{noise}$, the pump pulse might not reach the seed pulse at all because of the premature Raman backscattering by noise.

To prevent the premature backscattering and to ensure efficient amplification up to the maximum intensities, it is necessary to exclude the noise that grows by $\Lambda_{noise}$ exponentiations, while allowing the useful seed that grows by the $\Lambda_0 < \Lambda_{noise}$ exponentiations that are needed for entering the pump depletion regime. This can be achieved by an appropriate detuning the Raman resonance.

The frequency detuning $\delta\omega_{dt}$ in the moving location of the usefully amplified pulse may be allowed to grow with the amplification time $t$; it just should remain smaller than the growing frequency bandwidth of the pulse,

$$\Delta\omega_b = 1/\Delta t_b \approx \gamma_0^2 t/\Lambda_0 .$$

For a parasitic pulse, the opposite condition should be satisfied, namely, the parasitic pulse bandwidth should become smaller than the resonance detuning before the pump is parasitically depleted. Thus, the pump backscattering by noise can be suppressed without suppression of the useful amplification by the resonance detuning satisfying conditions

$$\frac{1}{\Lambda_0} > \frac{\delta\omega_{dt}}{\gamma_0^2 t} > \frac{1}{\Lambda_{noise}} . \tag{14}$$

The detuning can be produced, for instance, by the pump pulse frequency chirp, or by the plasma density gradient. By combining different kinds of the detuning, it is possible to arrange in different ways for the suppression of the parasitic scattering.

As a result of the resonance detuning, the amplified pulse acquires the frequency shift equal to the averaged frequency detuning,

$$\delta\omega_b = \overline{\delta\omega_{dt}}.$$

For a linear detuning, we have

$$\delta\omega_{dt} \propto t,$$

so that

$$\delta\omega_b = \delta\omega_{dt}/2.$$

In particular, when the detuning is produced by the pump frequency linear chirp, the amplified pulse acquires the same chirp as the pump has, because the encountered pump length is twice larger than the amplification distance.

## 8 Example of resonant backward Raman amplification parameters

To illustrate how the issues described and the formulas given in the preceding sections can be used to design a practical experiment, we give as an example the resonant backward Raman amplification and compression of a pump laser with wavelength $\lambda_a = 1/3\,\mu\text{m} = 1/3\ 10^{-4}\,\text{cm}$. For this wavelength, the pump frequency is

$$\omega_a = 2\pi c/\lambda_a \approx 5.65 \times 10^{15}\,\text{sec}^{-1},$$

so that the critical plasma concentration is

$$n_{cr} \approx 10^{22}\,\text{cm}^{-3},$$

and the wavebreaking pump intensity at the critical plasma concentration is

$$I_M \approx 1.5 \times 10^{18}\,\text{W/cm}^2.$$

For a pump of intensity $I_0 = 10^{15}\,\text{W/cm}^2$, the amplification can proceed without wavebreaking in plasmas of densities

$$n_e > n_{ebr} = n_{cr}(I_0/I_M)^{2/3} \approx 7.5 \times 10^{19}\,\text{cm}^{-3}.$$

The largest output intensity can be produced at the smallest of the these concentrations, namely $n_e = n_{ebr}$. At this density, the Langmuir wave frequency, approximately equal to the electron plasma frequency, is

$$\omega_e \approx 4.9 \times 10^{14}\,\text{sec}^{-1}.$$

Suppose an initial seed intensity and duration such that the number of exponentiations during the linear amplification stage is $\Lambda_0 = 6$. Suppose also that $\delta = 1$ is the tolerable nonlinear phase detuning of the 3-wave resonance. The output intensity achieved at the wavebreaking density is then

$$I_{bM} \approx 8.5 \times 10^{17}\,\text{W/cm}^2\,, \tag{15}$$

with an output duration of

$$\Delta t_{bm} \approx 1.3 \times 10^{-14}\,\text{sec}\,. \tag{16}$$

The initial seed duration $\Delta t_0$ should be about the output pulse duration $\Delta t_{bm}$. The initial seed intensity can be put as

$$I_{b0} \approx I_M \frac{256\pi\Lambda_0 \exp(-2\Lambda_0)}{\omega_a \omega_e \Delta t_0^2} \sim 10^{13}\,\mathrm{W/cm^2}.$$

The amplification time is

$$t_M \approx 8.2 \times 10^{-12}\,\mathrm{sec}.$$

The pump duration is $2t_M$. The thickness of the plasma layer is

$$L = ct_M \approx 2.5\,\mathrm{mm}.$$

However, to achieve the output intensity, given by Eq. (15) and the output duration, given by Eq. (16), it is necessary to overcome parasitic processes. For example, for the pump to reach the seed, we must ensure that it can traverse distance $L$ without incurring significant backscattering by noise, which grows exponentially during the linear stage of the backward Raman instability with the growth rate of

$$\gamma_0 \approx 7.5 \times 10^{12}\,\mathrm{sec}^{-1},$$

so, within the amplification time $t_M$, or plasma length $L$, the linear instability could make

$$\Lambda_M = \gamma_0 t_M \approx 62$$

exponentiations. Clearly, unless the parasitic backscattering were suppressed, any realistic noise level would prematurely deplete the pump.

As an example of noise, consider, for instance, the equilibrium thermal noise of plasma. The energy density of the resonant Langmuir wave noise in the plasma of electron temperature $T_e = 100\,\mathrm{eV}$ can be evaluated as

$$W_{eT} \sim \frac{T_e \gamma_0^2}{2\pi\lambda_a c^2} \sim 5 \times 10^{-9}\,\frac{\mathrm{J}}{\mathrm{cm}^3}.$$

The resonant Langmuir wave energy density needed for the pump depletion is

$$W_{ed} \approx \frac{2I_0\omega_e}{c\,\omega_a} \approx 6 \times 10^3\,\frac{\mathrm{J}}{\mathrm{cm}^3}.$$

Hence, the number of the Langmuir wave amplitude exponentiations from the noise to the pump depletion level is

$$\Lambda_{noise} \approx \frac{1}{2}\log\left(\frac{W_{ed}}{W_{eT}}\right) \approx 14.$$

To suppress the parasitic backscattering of pump by this level of noise, while not suppressing the useful amplification, as in Eq. (14), we apply a linear frequency detuning of the 3-wave resonance, which allows the linear instability to make, say, only 10 exponentiations,

$$\delta\omega_{dt} = \frac{\gamma_0^2 t}{10}.$$

The maximum frequency detuning, reached at the end of the useful backward Raman amplification, namely at $t = t_M$, is

$$\delta\omega_{dtM} = \frac{\gamma_0^2 t_M}{10} = \gamma_0 \frac{\Lambda_M}{10} \approx 4.6 \times 10^{13}\,\mathrm{sec}^{-1}.$$

This represents less than 1% of the laser pump frequency, which could be accomplished by chirping the pump laser. Alternatively, it represents a 10% change in the plasma frequency, which could be accomplished by an electron density variation of 20%. Significantly, without this detuning, the output pulse would reach only about a fraction $14/62$ of its potential energy and a fraction $(14/62)^2$ of its potential intensity.

## 9 Summary and discussion

What we presented here is a practical guide to using the key formulas for resonant Raman amplification and compression of intense laser pulses in plasma. By way of an illustrative example, we showed that to obtain the highest intensities, it is not nearly enough at all just to create conditions for amplification; it is crucial, simultaneously, to create conditions that suppress unwanted effects. We showed how one of the important unwanted effects that requires suppression is premature backscattering of the pump by plasma noise. The illustrative example showed how, in fact, a factor of as much as 16 would be lost in output intensity were the detuning methods not employed.

It is important to point out that both the unwanted effects considered in this brief guide, and others not considered here, are gone into in much greater depth in the archival literature. These include parasitic effects by plasma noise on Raman scattering of the pump and amplified pulses [1,15–19], pulse scattering by plasma density inhomogeneities [20], and the generation of super-luminous precursors of the amplified pulse [21]. Advantageous compression regimes are also made possible by selecting appropriate seed parameters, which include both the seed duration and intensity [21,22] as well as the the seed chirp [23].

The filamentation and the detuning of amplified pulses due to the relativistic electron nonlinearity are also treated in much greater depth [24–26]. The regime in density-temperature space of robust operating regimes has also been delineated in greater depth [27]. Other issues gone into greater depth include pulse depletion and plasma heating through inverse bremsstrahlung [28–33].

In addition to overcoming issues associated with competing deleterious effects, such as through resonance detuning, there are also important suggestions on how to improve the efficiency or make the effect more robust through multi-stage procedures [34,35]. These suggestions generally entail more than one plasma amplification stage, with the different stages featuring different plasma densities. The joint optimization of using resonant Raman compression together with CPA compression to produce the pump beam was also considered [36].

It is also important to point out that we omitted for consideration here certain alternative regimes or mechanisms of laser compression through counter-propagating waves in plasma. These alternative mechanisms include amplification based on Compton backscattering in very rarefied plasma, where wave trapping effects are important and where collective electron fields are ignorable [37]. The Compton backscattering suggestion was important, in part, in that it in fact stimulated the consideration of the resonant Raman backscattering considered here, where, rather than being negligible, collective electron effects dominate [1]. Similarly not dealt with here are interesting suggestions relating to the deep wavebreaking regime, which also occurs in relatively rarefied plasma and also for which a kinetic treatment is necessary [38]. Also, promising suggestions for amplification using Brillouin backscattering effects [39–41] are similarly not considered here. We also did not touch here on effects associated with the propagation of lasers in finite-width channels [42–44].

However, while this guide has necessarily been brief, and while much has not been considered here, it is hoped that the systematic presentation of a few of the most important effects will be a useful guide in experiments aimed at producing the resonant Raman backscattering and compression effect. It is further hoped that this effect will fulfill its full promise, namely to facilitate the next generation of laser intensities.

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