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Raman gain against a background of non-thermal ion fluctuations in a plasma

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Abstract. A complex stimulated Raman scattering event against a background of non-thermal ion acoustic waves in an inhomogeneous plasma is described. We obtain analytic forms for the Raman gain due to a five-wave interaction consisting of conventional three-wave Raman scattering followed by the decay of the Raman Langmuir wave into a second Langmuir wave (or a second scattered light wave) and an ion acoustic wave. Very modest levels of ion waves produce a significant effect on Raman convective gain. A combination of plasma inhomogeneity and suprathermal ion fluctuations may offer a means for the control of Raman gain.

Stimulated Raman scattering (SRS) and stimulated Brillouin scattering (SBS) have long been a focus of attention in laser–plasma interaction physics (Liu et al. 1974). Both result in light being scattered – an outcome to be minimized in laser fusion – while the other waves involved in the decay (Langmuir waves in the case of SRS and ion acoustic waves in the case of SBS) themselves give rise to unwelcome side-effects. Hardly surprisingly in view of these consequences, both instabilities have been examined time and again and their principal characteristics established, although various questions on both SRS and SBS remain open even today. Of particular concern to indirect-drive laser fusion is the control of SRS from the gas-fill in hohlraum targets (Fernández et al. 1996; Kirkwood et al. 1996). Most of the studies carried out have dealt with the two instabilities as independent parametric decays, largely on account of their wide spectral separation. However, early experiments demonstrated an anticorrelation between the occurrence of SRS and SBS (Walsh et al. 1984) which has been confirmed subsequently (Labaune et al. 1997; Fuchs et al. 2000) from observations providing *direct* evidence of a spatial and temporal interaction between the ion acoustic waves (IAW) from SBS and the Langmuir waves from SRS. Langmuir waves began to grow only when the spectrum of IAWs dropped to thermal levels. Artificially enhancing the level of IAWs significantly reduced the density fluctuations associated with Langmuir waves. These results naturally focus attention on those plasma–wave interactions that couple Langmuir waves and ion acoustic waves. These are the *Langmuir decay instability* (LDI), in which a primary Langmuir wave decays into a daughter Langmuir wave and an ion acoustic wave (DuBois and Goldman 1965), and the *electromagnetic decay instability* (EDI), where the daughter Langmuir wave is replaced by a second scattered electromagnetic wave (Shukla et al. 1983; Baker 1996). The link between SRS and the presence or otherwise of ion waves has been firmly established. Secondary counter-propagating

Langmuir waves needed for LDI have been identified directly using Thomson scattering (Labaune et al. 1997). Experiments with hohlraum targets have shown that SRS reflectivity rises with increased ion-wave damping (Fernández et al. 1996; Kirkwood et al. 1996).

These results prompted us to re-examine the question of Raman growth in the presence of IAWs. We have not attempted to describe the undoubtedly complex events that occur in the experiments. Instead, we explore a model of Raman gain in which both primary and secondary decays take place in an *inhomogeneous plasma*, so that the parametric resonances are local in nature. In our model, we assume that the two distinct decay mechanisms SRS–LDI and SRS–EDI occur with interaction regions that coincide. We refer mainly to SRS–LDI to illustrate our argument. The coincident SRS–LDI event can be described by a local five-wave process. Our model is further characterized by *prescribing* an enhanced level of IAWs. These may be present as a consequence of SBS, possibly enhanced by seeding, or pumped, for example by the optical mixing of crossed laser beams, or produced by other unidentified sources, as observed in recent experiments (Labaune et al. 1995). The secondary decay process is not then a parametric *instability* but is, in isolation, stable, giving rise only to a nonlinear frequency shift due to the suprathermal level of the IAWs. Hence we refer to these processes as SRS–LD and SRS–ED.

The model is as follows. Frequency matching requires, for SRS, that $\omega_0 = \omega_1 + \omega_{L1}$ and, for LD, that $\omega_{L1} = \omega_{L2} + \omega_s$, where ω_0 is the laser frequency, ω_1 is the scattered-light frequency, ω_{L1} and ω_{L2} are the primary and secondary Langmuir frequencies, and ω_s is the IAW frequency. In the case of ED, ω_{L2} is replaced by ω_2, the secondary scattered-light wave frequency. We consider SRS backscattering, since this has the largest growth rate $\gamma = \gamma_R$, where

$$\gamma_R = \frac{k_{L1}\omega_p v_0}{4(\omega_1\omega_{L1})^{1/2}},$$

ω_p is the electron plasma frequency, v_0 is the electron quiver velocity in the laser electric field, and k_{L1} is the Raman Langmuir wavenumber. The spatially dependent wavenumbers satisfy their respective dispersion relations,

$$k_{0,1,2}^2 = \frac{\omega_{0,1,2}^2 - \omega_p^2}{c^2}, \qquad k_{L1,L2}^2 = \frac{\omega_{L1,L2}^2 - \omega_p^2}{3v_e^2}, \qquad k_s^2 = \frac{\omega_s^2}{(c_s + u)^2}.$$

Here v_e is the electron thermal velocity, c_s is the ion acoustic speed and u is the expansion velocity of the plasma. The wavenumber mismatch factors for SRS and LD respectively are

$$K_R(x) = k_0(x) + k_1(x) - k_{L1}(x), \qquad K_L(x) = k_s(x) - k_{L1}(x) - k_{L2}(x).$$

The resonances are assumed to be coincident at the origin, where wavenumber matching is exact for both primary and secondary processes: $K_R(0) = K_L(0) = 0$. Resonance widths are determined mainly by the spatial dependence of the Langmuir wavenumbers. Figure 1 shows the geometry of the five-wave SRS–LD interaction: the laser pump, IAW and primary Langmuir wave enter the resonance from the left, while the scattered light and secondary Langmuir wave enter from the right. The geometry is such that $k_{L1} \approx k_{L2} \approx \frac{1}{2}k_s$.

We assume that the IAW, while consistent with a low level of density fluctuations, is sufficiently above 'noise' that its amplitude remains undepleted across the gain region. We also neglect depletion of the laser pump. The problem then amounts to

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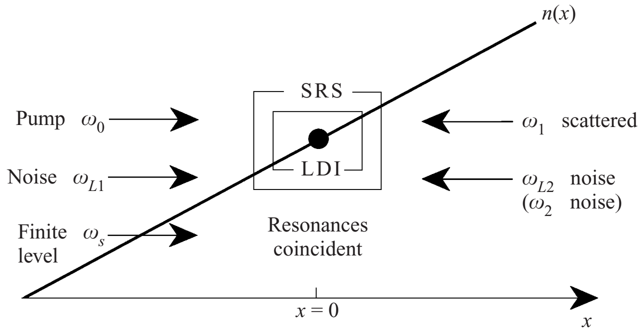

Figure 1. The five-wave geometry of the SRS–LD interaction.

a study of the evolution of the Raman-scattered light wave, and primary and secondary Langmuir waves each growing from noise in the presence of the laser pump and IAW. This interaction is distinct from a coupling of two *instability* mechanisms in which the Langmuir wave generated by SRS acts as the pump for a secondary instability. In our configuration, the IAW is present from the outset, and acts solely as a moving 'grating', reflecting each Langmuir wave into the other. The coupled equations for an inhomogeneous plasma in the WKB approximation are

$$\dot{\psi}_1 - V_1\psi_1' + \nu_1\psi_1 = \gamma_R\psi_{L1}\exp\left(i\int K_R\,dx\right), \tag{1a}$$

$$\dot{\psi}_{L1} + V_{L1}\psi_{L1}' + \nu_{L1}\psi_{L1} = \gamma_R\psi_1\exp\left(-i\int K_R\,dx\right) + \Delta_N\psi_{L2}\exp\left(-i\int K_L\,dx\right), \tag{1b}$$

$$\dot{\psi}_{L2} - V_{L2}\psi_{L2}' + \nu_{L2}\psi_{L2} = -\Delta_N\psi_{L1}\exp\left(i\int K_L\,dx\right), \tag{1c}$$

where V_i and ν_i ($i = 1, L1, L2$) denote group velocities and damping rates, $\Delta_N = \frac{1}{4}\omega_p N$, and $N = \delta n_i/n_0$ is the normalized IAW density fluctuation, with n_0 the plasma density. Dots and primes denote time and space derivatives respectively. The scattered wave is subject to collisional damping, $\nu_1 = \omega_p^2\nu_{ei}/2\omega_1^2$, and the Langmuir waves are subject to collisional and Landau damping, $\nu_{L1,L2} = \gamma_{LD} + \frac{1}{2}\nu_{ei}$, where ν_{ei} is the electron–ion collision frequency and γ_{LD} the Landau-damping rate.

We look first at the effect of damping combined with a finite level of IAWs in the homogeneous limit. Taking $K_R = K_L = 0$ and assuming quantities to vary as $\exp(\gamma t)$, the cubic equation resulting from (1) gives the threshold and growth rate. In the absence of the secondary decay ($\Delta_N = 0$), we recover the peak SRS homogeneous plasma growth rate $\gamma = \gamma_R$, which applies well above the convective instability threshold set by $\gamma_R^2 = \nu_1\nu_{L1}$. Including the LD, a finite level of IAWs increases the convective threshold to

$$\gamma_R^2 = \nu_1\left(\nu_{L1} + \frac{\Delta_N^2}{\nu_{L2}}\right), \tag{2}$$

which shows the contrasting effect of the damping of the two Langmuir waves. Two distinct regimes are apparent, determined by the ratio $\rho_0 = \Delta_N/\nu_{L1}$ ($\nu_{L1} \approx \nu_{L2}$ for LD). For weakly damped Langmuir waves, the threshold increase could be substantial, since ρ_0 can be large for even low levels of IAWs. Indeed, if damping is weak so that $\gamma_{LD} \ll \nu_{ei} \ll \Delta_N$, the threshold is set only by the level of IAWs,

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$\gamma_R = \omega_p \Delta_N / \omega_1$. Clearly, whenever $\rho_0 > 1$, the influence of the IAWs on SRS is significant.

It is the effect of the finite level of IAWs through the nonlinear frequency shift Δ_N in inhomogeneous plasma that is the focus of this work, so we confine attention to the undamped equations. Considering the temporal behaviour $(\partial/\partial x \to 0)$, we obtain from (1) the homogeneous growth rate $\gamma = (\gamma_R^2 - \Delta_N^2)^{1/2}$. Whenever the LD frequency shift $\Delta_N \geqslant \gamma_R$, Raman growth is substantially suppressed. In the absence of SRS, the LD is stable, giving rise only to a frequency shift Δ_N.

Henceforth, we consider an inhomogeneous plasma and analyse the stationary convective process $(\partial/\partial t \to 0)$. The mismatch factors, taken to be linear, are $K_{R,L}(x) = K'_{R,L}(0)x$. The scale lengths implicit in this model are the SRS coupling length $(V_1 V_{L1})^{1/2}/\gamma_R$, the LD coupling length $(V_{L1} V_{L2})^{1/2}/\Delta_N$ and the resonance widths set by the mismatch factors: $(K'_R)^{-1/2}$ for SRS and $(K'_L)^{-1/2}$ for LD. The SRS resonance is approximately $\sqrt{2}$ larger than the LD resonance, since these are determined mainly by the Langmuir waves. In the absence of LD, we retrieve the well-known Rosenbluth gain factor (Liu et al. 1974): $G_R = 2\pi\gamma_R^2/V_1 V_{L1}|K'_R|$. Formally, we can define a negative 'gain' factor for the LD: $G_N = 2\pi\Delta_N^2/V_{L1}V_{L2}|K'_L|$.

Our goal is to determine the gain from the behaviour of the (steady-state undamped) solutions to (1) over the complex plane subject to the boundary conditions $|\psi_{L1}|_{x=-\infty} = \delta_{L1}$, $|\psi_1|_{x=+\infty} = \delta_1$, $|\psi_{L2}|_{x=+\infty} = \delta_{L2}$ under the additional assumption that the phases of sources entering the resonance from opposite sides are statistically independent. The Stokes analysis used to find asymptotic solutions is presented elsewhere (Barr et al. 2000), and leads to expressions for the amplitudes of each wave emerging from the resonant region:

$$\begin{aligned}|\psi_{L1}^+|^2 = {} & \delta_{L1}^2 \exp(G_R - G_N) \\ & + \delta_{L2}^2 \frac{V_{L2}}{V_{L1}} \exp(G_R)\,[1 - \exp(-G_N)] \\ & + \delta_1^2 \frac{V_1}{V_{L1}} [\exp(G_R) - 1],\end{aligned} \tag{3a}$$

$$\begin{aligned}|\psi_1^-|^2 = {} & \delta_{L1}^2 \frac{V_{L1}}{V_1} \exp(-G_N)\,[\exp(G_R) - 1] \\ & + \delta_{L2}^2 \frac{V_{L2}}{V_1} [1 - \exp(-G_N)][\exp(G_R) - 1] \\ & + \delta_1^2 \exp(G_R)\end{aligned} \tag{3b}$$

$$|\psi_{L2}^-|^2 = \delta_{L1}^2 \frac{V_{L1}}{V_{L2}} [1 - \exp(-G_N)] + \delta_{L2}^2 \exp(-G_N), \tag{3c}$$

where the superscripts $\pm$ refer to that side of the resonance from which the wave emerges. This is the main result of this paper. The contributions from each source to net gain in the three daughter waves as a function of the SRS gain G_R and the LD negative gain G_N are clearly seen from (3). First, the gains in the primary Langmuir wave and the scattered wave that derive from the *primary* Langmuir wave source δ_{L1} are suppressed by G_N. Secondly, the LD coupling to SRS due to the IAWs allows gain in the Raman products to derive from the *secondary* Langmuir wave source δ_{L2}. This is enhanced by the finite level of IAWs, and emerges to replace the primary Langmuir wave as the main source whenever $G_N \gg 1$. Finally, gain from the scattered-wave source δ_1 is unchanged.

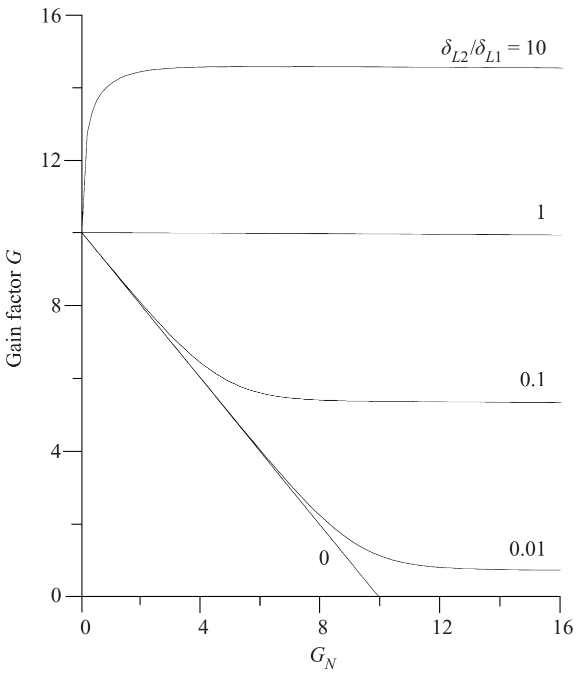

Figure 2. Net gain of the Langmuir wave G versus G_N for a range of source ratios δ_{L2}/δ_{L1}. A Raman gain $G_R = 10$ is assumed.

When the energy flux of each Langmuir wave source entering the resonance from opposite sides is equal, i.e. $T \equiv \delta_{L2}^2 V_{L2}/\delta_{L1}^2 V_{L1} = 1$, the terms involving G_N cancel. Net gain reverts to the levels expected in the absence of secondary decay. Thus the condition $T = 1$ must be violated to suppress or enhance the usual Raman gain. Since, for LD, $V_{L1} \approx V_{L2}$ and if the two Langmuir waves emerge from equal thermal sources $\delta_{L1} \approx \delta_{L2}$, as at first sight seems most likely, then indeed $T \approx 1$. However, should one or other of the Langmuir wave sources be enhanced, there is potential for either *decreased* or *increased* net gain, depending on whether or not $\delta_{L1} > \delta_{L2}$. Since the counter-propagating waves enter the resonance from opposite sides, there may well be good reasons for these sources to be both non-thermal and distinct, given the environment created by the intense laser driver. For instance, there may be regions where Raman gain is low or even sub-threshold, contiguous with localized high-intensity speckles, which provide enhanced source levels of the Raman Langmuir wave; then $\delta_{L2} \ll \delta_{L1}$, and Raman gain is strongly suppressed. Consider the gain due to $\delta_{L1,L2}$ only, and assume that the scattered-wave source is small compared with electrostatic sources (set $\delta_1 = 0$). Figure 2 shows the net gain of the Raman Langmuir wave, $G = \ln|\psi_{L1}^+/\delta_{L1}|^2$, versus G_N, assuming a Raman gain $G_R = 10$ for various ratios δ_{L2}/δ_{L1}. When $G_N \gg 1$, this approaches $G = G_R + \ln T$, so that total suppression occurs if $\delta_{L2}/\delta_{L1} < \exp(-\frac{1}{2}G_R)$. Also in this limit, the ratio of the Langmuir wave amplitudes becomes $|\psi_{L2}^-/\psi_{L1}^+|^2 = \delta_{L1}^2/\delta_{L2}^2 \exp(-G_R) = \exp(-G)$, providing a direct measure of the net gain G without requiring individual knowledge of the source amplitudes or of G_R and G_N. For SRS–ED decay, expressions similar to (3) can be obtained, with $L2 \rightarrow 2$. Again assuming that the electrostatic source is dominant, we find that the net gain in both the Raman Langmuir and scattered waves is unconditionally quenched by the presence of the IAW; the gain for the Raman Langmuir wave is $G = G_R - G_N$.

To illustrate the above effects, we use two parameter regimes – one typical of indirect-drive ignition experiments with hohlraum targets (Fernández et al. 1996; Kirkwood et al. 1996), the other corresponding to the plasma source in the experiments of Labaune et al. (Labaune et al. 1997, 1998; Depierreux et al. 2000) In the first, strong Landau damping would be expected to be dominant, while in the second cooler source, Landau damping is weak. Collisional damping is assumed weak for both cases. Typical parameters for ignition experiments are laser intensity $I = 2 \times 10^{15}$ W cm^{-2}, wavelength $\lambda = 0.35$ μm, density $n_0/n_c = 0.1$ and temperature $T = 3$ keV (Fernández et al. 1997). Then $k_{L1}\lambda_D = 0.38$ and a Landau damping rate $\gamma_L \approx 0.075\omega_p$ implies a strongly damped regime. A measure of the secondary decay is $\rho_0 = \Delta_N/\nu_{L1} \approx 3.3N$. Thus $\rho_0 \ll 1$ for all but unreasonably large values of N, and hence IAWs would have little effect on SRS. However, SRS has been shown to depend on the IAW damping rate and hence on the IAW amplitude (Fernández et al. 1997). In our model, this requires Landau damping sufficiently weak that $\rho_0 > 1$. In fact, Langmuir-wave damping is believed to be much weaker than linear Landau damping would imply (Afeyan et al. 1998).

In contrast, the source used by Labaune et al. (1995, 1997, 1998) consisted of exploded plastic foils, for which we take $I = 9\times10^{13}$ W cm^{-2}, $\lambda = 1$ μm, $n_0/n_c = 0.1$, $T = 0.5$ keV and scale length 1 mm. Then $k_{L1}\lambda_D = 0.15$, $\gamma_L \approx 2 \times 10^{-8}\omega_p$ and $\rho_0 \approx 10^7 N$. Since $\rho_0 \gg 1$ for very low levels of IAWs, the coupling to the secondary Langmuir wave will be important in this case. The convective threshold is then set by the IAW level, $\gamma_R = \omega_p\Delta_N/\omega_1$, which implies $v_0/c = (\omega_p^3/k_{L1}^2c^2\omega_1)^{1/2}N$, so that $v_0/c \approx 0.14N$. Since $v_0/c \approx 0.008$, this requires $N = 6\%$ at the convective threshold. Were the intensity to be enhanced within 'speckles' of localized hot spots (Tikhonchuk et al. 1996), even higher IAW levels would be required to bring SRS to the convective threshold. Such levels are substantially above what is observed (Labaune et al. 1998; Depierreux et al. 2000). We can therefore conclude that these experiments are well above the damping threshold.

Suppose that damping can be neglected and growth is convectively saturated owing to plasma inhomogeneity. For the above conditions, $G_R = 1.25$, so that observable gain requires some form of local enhancement in intensity, as in 'speckles'. Also, $G_N \approx 5.5\times10^4N^2$, so that an IAW density fluctuation of only $N \approx 0.4\%$ gives $G_N = 1$. For SRS suppression, we require, in addition, that $\delta_{L1} > \delta_{L2}$. Labaune et al. measured the ratio of secondary to primary Langmuir wave amplitudes of the density fluctuations at about 6%, from which, assuming $G_N \gg 1$, we predict a net gain $G = 2\ln|\psi_{L1}^+/\psi_{L2}^-| \approx 5.6$. This compares with the measured gain (amplitude squared) of $G \approx 7$ (Labaune et al. 1998). Assume a local enhancement of intensity within a speckle of a factor 5, so that $G_R = 6.25$. Taking $N = 1\%$ and $\delta_{L2}/\delta_{L1} = 0.1$ gives a reduced gain of $G = 2$, while for $N = 2\%$ and $\delta_{L2}/\delta_{L1} = 0.04$, there is complete quenching. It appears that modest levels of ion fluctuations, typically around 1%, combined with a differential in the Langmuir-wave source levels, is all that is needed for the SRS signal to be reduced significantly.

The strong evidence for SRS suppression while SBS is active is a compelling argument for identifying the SBS IAW with the IAW of the SRS–LD process. Phase matching dictates that this is so only for densities approaching the quarter-critical density, although observations show SRS inhibition over a wide range of densities. However, exploding-foil experiments with evolving inhomogeneous density and flow-velocity profiles constitute an environment in which SBS ion waves can adjust their wavenumbers by propagation to match those of the SRS–LD process. Our model does not in fact depend on the ion fluctuation satisfying the usual IAW dispersion relation – any finite level of non-resonant ion waves (ω_s, k_s) would suffice, and indeed this has been observed with spectra that allow both SRS–LD and SRS–ED decays to arise at virtually any density (Labaune et al. 1995). We emphasize too that the Rosenbluth gain G_R, the IAW level N and possibly the source levels themselves are each dependent on laser intensity which itself is likely to have a complex spatial pattern due, for example, to a speckle distribution. If the IAW level is generated by SBS, the negative gain G_N will rise exponentially, while G_R increases only linearly, with laser intensity. The theory developed here may offer an explanation in part for features observed in the interplay between SRS and SBS in laser-produced plasmas. For instance, in Labaune et al. (1997) and Fuchs et al. (2000), the SRS Langmuir wave was diminished by the presence of SBS and further suppressed by seeding the IAWs. Other seeding experiments are suggested by our results. Another observation connecting IAWs with the SRS signal was its observed dependence on the damping of the IAW (Fernandez et al. 1996; Kirkwood et al. 1996). Clearly, the IAW level is sensitive to its damping rate, so that, as our model shows, this affects the net Raman gain.

We have determined the net gain of Raman-generated waves due to five-wave SRS–LD and SRS–ED interactions when a finite level of ion waves is present in an inhomogeneous plasma. Modest ion fluctuation levels have a significant effect on gain, which is sensitively dependent on the source levels from which the waves amplify. Inhomogeneous plasma combined with a suprathermal level of IAWs may afford a means for the control of SRS.

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