Saturation of Stimulated Raman Scattering by the Excitation of Strong Langmuir Turbulence

B. Bezzerides, D. F. DuBois, and Harvey A. Rose
Los Alamos National Laboratory, Los Alamos, New Mexico 87545
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Numerical studies of a model of stimulated Raman scattering (SRS) coupled to Langmuir turbulence show that once the collisional threshold is exceeded for parametric decay of the SRS daughter Langmuir wave the noncollisional dissipation resulting from Langmuir collapse can strongly saturate the Langmuir wave and the resulting SRS reflectivity. The model provides information on the total rate of hot electron generation. Scaling laws for the reflectivity as a function of laser and plasma parameters are found.

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Stimulated Raman scattering (SRS) of laser light from plasmas has been studied [1] since soon after the first observation and interpretation of the analogous effect of SRS from liquids and solids in the early 1960s. Since the early 1970s [2] many experiments have observed SRS from plasmas [3]. During the 1980s a number of important deductions were made concerning the comparison of theory and experiment [3–5].

In our opinion no theory to date accounts for all of these observations. The purpose of the present work is to investigate further the predictions of a nonlinear theory [6] which couples SRS to Langmuir turbulence. Our recent studies of the saturation of the Langmuir parametric decay instability (PDI) [7,8] lead to the premise that strong Langmuir turbulence effects can play a critical role in the saturation of SRS. This saturation by the decay cascade emanating from the daughter SRS Langmuir wave has been studied in the weak coupling limit where a finite set of linear Langmuir waves are excited [9]. Recently, Drake and Batha [10] hypothesized that the intensity of the daughter SRS Langmuir wave was saturated at the threshold for its subsequent decay into another Langmuir wave and an ion acoustic wave.

The studies to be reported here are based on numerical simulations of a one-dimensional reduced description of the coupling of backscatter SRS to Langmuir turbulence which is a variant of the equations used in Ref. [6]:

$$2i\omega_R(\partial_t+\nu_R-v_R\partial_x)A_R=-\omega_{p_0}^2(n/n_0)A_R+(\omega_{p_0}^2/8\pi e n_0)A_0\partial_x E^*\exp i(k_1x-\tfrac{3}{2}k_1^2v_e^2\omega_{p_0}^{-1}t)\,, \tag{1}$$

$$[2i\omega_{p_0}(\partial_t+\nu_e)-\omega_{p_0}^2(n/n_0)+3v_e^2\partial_x^2]E=\tfrac{1}{2}(\omega_{p_0}^2/c^2)(e/m)ik_1(A_0A_R^*)\exp i(k_1x-\tfrac{3}{2}k_1^2v_e^2\omega_{p_0}^{-1}t)\equiv S(x,t)\,, \tag{2}$$

$$[\partial_t^2+2\nu_i\partial_t-c_s^2\partial_x^2]n=(16\pi m_i)^{-1}\partial_x^2[|E|^2+c^{-2}(\omega_R^2|A_R|^2+\omega_0^2|A_0|^2)]\,. \tag{3}$$

Here we assume spatial variation only in the x direction and represent the total electromagnetic vector potential A^T (which is in a direction perpendicular to x) in the envelope form

$$A^T=\tfrac{1}{2}[A_0\exp i(k_0x-\omega_0t)+A_R\exp i(k_Rx-\omega_Rt)]+\text{c.c.}\,, \tag{4}$$

where ω_0,k_0 and ω_R,k_R are the incident and scattered light frequencies and central wave numbers, respectively. In this paper we chose ω_0 and ω_R to be related by the SRS frequency matching condition, $\omega_0=\omega_R+(\omega_{p_0}^2+3k_1^2v_e^2)^{1/2}$, where $k_1=k_0-k_R$ is the wave number of the most unstable daughter Langmuir wave at the reference density, n_0. The equation for the pump wave, A_0, is obtained from (1) by the formal substitutions $A_R\to A_0$, $\omega_R\to\omega_0$, $E^*\to E$, $\nu_R\to\nu_0$, $v_R\to v_0$, and $\omega_{p_0}\to-\omega_{p_0}$. The incident and scattered light group velocities, v_0 and v_R, are taken to be in opposite directions. Here E is the slowly varying envelope of the electrostatic field where the complete field is written as $\frac{1}{2}E\exp(-i\omega_{p_0}t)+\text{c.c.}$, and $\omega_{p_0}^2=4\pi e^2n_0/m$. The ion density (or low frequency part of the electron density) is given as n_0+n where n_0 is the averaged or reference background density and n the fluctuation. The collisional damping operators, ν_R,ν_0, on the light are given by standard formulas. The damping operators ν_e and ν_i are local in k space where $\nu_e(k)$ is the sum of collisional damping—$\nu_e(k=0)$—which is independent of k, and Langmuir Landau damping whose dissipation is properly interpreted as energy going into hot electrons [7,8]; $\nu_i(k)$ is the usual ion wave Landau damping, $\omega_i=kc_s$, and c_s the sound speed, $c_s=(ZT_e/m_i)^{1/2}$ in the limit $ZT_e/T_i\gg1$, and $v_e^2=T_e/m_e$.

It is well known [11] that when the intensity of the SRS daughter Langmuir wave, W_{k_1}, exceeds the damping threshold value—$E_T^2/4\pi n_0T_e$ which is equal to $16[\nu_e(k_2)/\omega_p](\nu_i/\omega_i)$—that this mode can decay into another Langmuir wave and an ion acoustic with wave vectors $\mathbf{k}_2$ and $\mathbf{k}_2'$ where $\mathbf{k}_1=\mathbf{k}_2+\mathbf{k}_2'$. This threshold is effective, provided $3k_1\Delta kv_e^2<\nu_e(k_1)\omega_{p_0}$, even when W_{k_1} is a superposition of modes, $W_{k_1}=|E_1|^2(4\pi n_0T_e)^{-1}$, where E_1 is the envelope relative to k_1:

$$E_1(x,t)=\exp(-ik_1x)\sum_{k\epsilon\Delta k}E_k\exp(ikx)L_s^{-1/2}\,; \qquad k_1-\Delta k/2<k<k_1+\Delta k/2\,. \tag{5}$$

The decay cascade provides a pathway for the saturation of SRS via the dissipation associated with Langmuir turbulence. This threshold is generally much lower than the Langmuir wave breaking threshold [10]. Typically we find [12] very low reflectivities are required to exceed the PDI threshold if I_0 exceeds the homogeneous absolute damping threshold for SRS.

As a beginning step toward understanding we have numerically solved (1)-(3) for an initially homogeneous plasma with a finite SRS-active region of length L. This model can be thought of as representing a finite slab of plasma or a finite region of intense pump field as would occur in a laser hot spot with or without self-focusing [13]. (Such hot spots resulting from a random phase plate, for example, have a longitudinal correlation length $\sim f^2\lambda_0$ in terms of the f number of the lens.) We use a spectral method based on a simulation cell of length $L_S > L$ with incoming boundary conditions for the light and effective outgoing boundary conditions for the sound and Langmuir waves in the active region. The collisional damping length for the Langmuir waves confines them to the active region. We have arbitrarily killed the stimulated Brillouin scattering (SBS) coupling by spatial enveloping and dropping the $|A_0|^2$ contribution to the ponderomotive force in (3). Our experience [6] leads us to believe that with respect to SBS coupling, our results represent an upper bound on the SRS reflectivity.

In this Letter we will emphasize the long time behavior, independent of the details of the onset of the laser pulse, after which the system appears, for fixed laser intensity, to approach a statistically stationary state independent of initial noise levels. The early transient behavior depends in detail on these levels and on the detailed rise of the laser pulse. Typically there is a strong initial spike in the reflected intensity during which wave breaking saturation effects, not included in our model, may be important. Once the ions have time to move the reflectivity settles down to much lower levels. We find that the asymptotic turbulent state is independent of initial conditions for a wide range of initial conditions. We assume that this state is a valid prediction of the model even for initial conditions leading to strong transients.

Our simulations show that the reflectivity R, which is equal to $(\omega_R/\omega_0)^2(v_R/v_0)|A_R/A_0|^2$, is an intermittent function of time with nearly 100% amplitude modulations. In Fig. 1 we show the time averaged spatial profiles of scattered wave intensity, the Langmuir intensity $|E|^2$, and the density fluctuation. Note that the averaged density is no longer uniform because the time averaged ponderomotive forces of the high frequency waves expels electrons and ions from the confined region of interaction. This induced density inhomogeneity produces, as expected, a finite frequency spread in the power spectrum of the backscattered light. Note that for convenience we have chosen our reference density n_0 to be the average density so that $\int_0^L dx\, n(x) = 0$.

The density depressions caused by localized hot spots in a laser beam, perhaps including self-focusing effects, can locally perturb the density gradient profile to allow absolute instability of SRS which otherwise might not occur [13]. The hot spot length rather than the detuning length in the unperturbed gradient [10] then controls the strength of SRS.

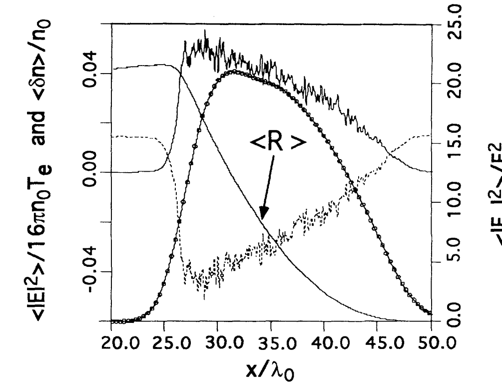

FIG. 1. For the parameters: $\lambda_0 = 0.35\ \mu$m, and $L/\lambda_0 = 22.5$, $I_0/I_T = 64$ where $I_T = 2\times10^{14}$ W/cm^2, is the absolute instability threshold intensity. The plasma parameters are $n/n_c = 0.2$, $T_e = 1$ keV, $T_e/T_i = 2.4$, $Z = 4$, $A = 9$ for which $k_1\lambda_D = 0.125$. Time averaged (over many laser cycles) spatial profiles of the excited Langmuir energy, $\langle|E(x)|^2\rangle/16\pi n_0 T_e$ (solid curve), the fractional density modulation, $\langle\delta n(x)\rangle/n_0$ (dashed curve), and the local reflectivity marked R (in arbitrary units). The laser enters from the left. The maximum reflectivity was about 1.2%. The solid line with dots is the time averaged profile of the square of envelope of the daughter SRS Langmuir waves, $\langle|E_1(x)|^2\rangle$, in units of the PDI threshold intensity E_T^2.

Simulations clearly show that driven caviton collapse occurs for the parameter regime of Fig. 1 where $k_1\lambda_D \simeq 0.12$ with a collapse duration of a fraction of a ps. At lower densities, e.g., $n/n_c = 0.1$ where $k_1\lambda_D \simeq 0.22$ for the parameters of Fig. 1, we find that collapse is suppressed. We find that the collapse is nucleated [7] at spatial scales associated with the density modulations at $k \simeq 2k_1$ arising from the PDI. For the higher values of $k_1\lambda_D$ there is not a sufficient inertial range between nucleation scales and dissipation scales for collapse to develop [14]. In Fig. 2 we plot the time averaged reflectivity $\langle R\rangle$ versus the laser intensity I for several parameter regimes. The Langmuir turbulence saturates $\langle R\rangle$ at low values. Also shown is the fraction of laser energy, $D/\langle R\rangle$, which is absorbed by noncollisional Langmuir dissipation where $DI = 2\sum[\nu_e(k) - \nu_e(0)]|E_k|^2$. This is a measure of the energy going into hot electrons via collapse. In the simulations reported here, except for the weakest value of I, more than 50% of the Langmuir energy goes into hot electrons.

As a diagnostic of the envelope of the "daughter" part of the Langmuir field which couples directly to the SRS we have calculated $E_1(x,t)$ defined in (5) where the band of modes centered about k_1 is determined as follows: The reflectivity is recalculated using only those modes in Δk

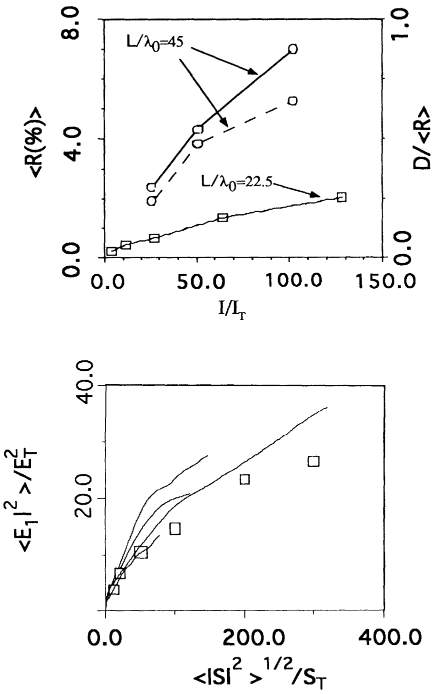

FIG. 2. Top panel: Time averaged reflectivity R (%) (solid curves) vs laser intensity I in units of the absolute instability threshold intensity I_T for two values of the length L of the SRS active region. Except for L all parameters are as in Fig. 1. Also shown is the ratio, D/R, of the energy going into hot electrons by noncollisional dissipation to the reflected laser energy. Bottom panel: The local relationship between the modulus squared of the envelope of the daughter Langmuir waves, $\langle |E_1|^2\rangle$, in units of the PDI threshold E_T^2 and the local source [the right-hand side of Eq. (2)], in units of the linear threshold source amplitude, $S_T = 2\omega_{p_0} v_e(k=0)E_T$. Top curve: $L/\lambda_0 = 45$, $I/I_T = 26$. Next lowest curve: same as in Fig. 1. Third curve: $L/\lambda_0 = 22.5$, $I/I_T = 130$, and ν_{ei} is half of that in Fig. 1. Bottom curve: $L/\lambda_0 = 22.5$, $I/I_T = 128$. Unless otherwise stated the parameters are the same as in Fig. 1. Boxes: determined from homogeneous turbulence simulations.

which are increased from 0 to its final value, Δk_1, for which the reflectivity is, say, 90% of the reflectivity calculated with all of the modes. It is found that $3k_1 \Delta k_1 v_e^2 < \nu_e(k_1)\omega_{p_0}$. [Note that it is physically meaningless to look at a single Fourier mode, say $E(k_1)$, whose amplitude depends on the length L_S. It is necessary to take L_S larger than the important correlation length, $(\Delta k)^{-1}$, to get physically meaningful results.] A typical Fourier energy spectrum for a fairly strongly driven regime is shown in Fig. 3 where the bandwidth Δk_1 turns out to be essentially the FWHM of the peak centered at k_1. The remaining modes represent the turbulence resulting from cascades and collapse which does not couple directly to SRS. More weakly driven simulations show [12] spectral

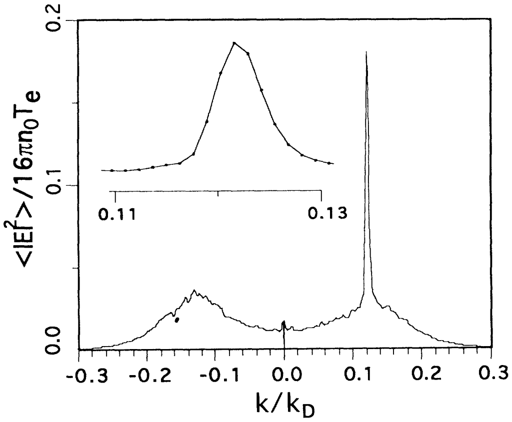

FIG. 3. Time averaged modal energy spectrum, $\langle |E(k)|^2\rangle / 16\pi n_0 T_e$, of Langmuir fluctuations vs k/k_D for the parameters of Fig. 1. The inset is an enlargement of the peak near $k_1\lambda_D = 0.125$. More weakly driven spectra of this type show well defined cascade features.

characteristics of modified weak turbulence cascades [7,8].

To gain some insight into the validity of a local description of the Langmuir turbulence we have plotted in Fig. 2 (bottom panel) the time averaged modulus $\langle |E_1(x)|^2\rangle / E_T^2$ versus the rms of the right-hand side of (2), which we call $\langle |S(x)^2|\rangle^{1/2}$, both quantities are obtained locally from a global simulation. In Fig. 2 these curves are shown for a range of values of I, L, and ν_{ei} over which there is a wide range of the absolute magnitudes of $\langle |E_1|^2\rangle$ and $\langle |S|\rangle$. The fact that all of these curves lie remarkably close to one another suggests that there is a nearly universal behavior of these local quantities when scaled in terms of the PDI threshold field, E_T, and the source strength S_T which would lead to E_T in the *linear* response of E; i.e., $S_T = 2\omega_{p_0} E_T \nu_e(k_1)$. We have also plotted the ratio $\langle |E_1|^2\rangle / E_T^2$ obtained from a homogeneous, source driven simulation of the Langmuir turbulence as a function of $|S|/S_T$. A fair fit to the data in Fig. 2 can be obtained for this ratio with the formula $(\langle |S|\rangle / S_T)^{2/\eta}$ with $3.2 < \eta < 3.5$. A rough, mean-field-theory, scaling formula for $\langle R\rangle$ can then be obtained as follows: Ignore fluctuations about the mean, assume $|(\partial_t + \nu_R)A_R| \ll |v_R \partial_x A_R|$; neglect pump depletion; write $|E| = \langle |E_1(x)|^2\rangle^{1/2}$ using the fitting formula; substitute for $|S|$ using (2). This produces a nonlinear ordinary differential equation for $|A_R|$ which can be integrated to yield the scaling formula $(\langle R\rangle/\nu_i)^{\eta-1} \propto \nu_e^{\eta-3} I^2 L^{2\eta}$. We have verified by simulations some predictions of this formula that $\langle R\rangle \propto \nu_i$, for $\nu_i > 0.02$, and increases less strongly than a linear function of I (Fig. 2) and increases with L with a power law between L^2 and L^3. For L fixed (say by the hot spot length) this scaling with ν_e is much weaker than the $\langle R\rangle \propto \nu_e L^2$ proposed in [10]. Remarkably when $L = 2(v_e/\omega_p)L_n$, the detuning length, our scaling gives $\langle R\rangle \propto v_i v_e^3 I^{2/(\eta-1)}$; the result in [10] is the $\eta \rightarrow \infty$ limit of this. The numerical values of $\langle R\rangle$ predicted by the complete form of this formula can differ from the simulation values by factors as large as 3 to 5. We attribute this discrepancy to nonlocal and fluctuation effects and phase-dependent effects on the light propagation. The local analysis breaks down in the simulation region to the right where A_R is weak and an exponentially varying function of x. Aspects of this in the weak regime have been discussed in [9]. The local gain length in this region is comparable to the estimated correlation length of the local turbulence. To the extent that the local analysis is valid, as suggested by Fig. 2, the turbulence properties, including the dissipation into hot electrons, depend only on the strength of $|S|$. Since $S \propto IR^{1/2}$, the scaling law for R predicts $|S|^{\eta-1} \propto (IL)^{\eta}$. Thus equal values of IL give roughly equal turbulence levels; and the value of I for which collapse effects dominate scales as L^{-1}. A more detailed analysis of all these points and others will be given in a longer paper [12].

Clearly there are important physical effects ignored or only crudely treated in this modeling. We regard this as only a first step in understanding the saturation of SRS by strong Langmuir turbulence which may provide a paradigm for further studies. Our results show that once the amplitude of the SRS daughter Langmuir wave exceeds its parametric decay damping threshold the strong dissipation resulting from Langmuir collapse can saturate the SRS reflectivity, R. This has led to simple mean-field scaling laws for R and provides a basis for and an improvement on the conjecture of Drake and Batha [10].

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