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PARAMETRIC INSTABILITIES IN PICOSECOND TIME SCALES

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INTRODUCTION

The coupling of intense laser light with plasmas is a rich field of plasma physics, with many applications. Among these are inertial confinement fusion (ICF), x-ray lasers, particle acceleration, and x-ray sources. Parametric instabilities have been studied for many years because of their importance to ICF; with laser pulses with duration of approximately a nanosecond, and laser intensities in the range $10^{14} - 10^{15} \text{ W/cm}^2$ these instabilities are of crucial concern because of a number of detrimental effects. Although the laser pulse duration of interest for these studies are relatively long, it has been evident in the past years that to reach an understanding of these instabilities requires their characterization and analysis in picosecond time scales. At the laser intensities of interest, the growth rate for stimulated Brillouin scattering (SBS) is of the order of picoseconds, and of an order of magnitude shorter for stimulated Raman scattering (SRS).

In this paper we will discuss SBS and SRS in the context of their evolution in picosecond time scales. We will describe the fundamental concepts associated with their growth and saturation, and recent work on the nonlinear treatment required for the modeling of these instabilities at high laser intensities. Following this, we will review two types of

experiments. The first is an experiment performed with laser pulses with a duration of 600 ps, which clearly demonstrates the short-time nature of many observed features. The second experiment describes the potentials of using picosecond laser pulses as a "pump" for the study of parametric instabilities. It is hoped that this review will motivate future studies, both experimental and theoretical, that will help the understanding of growth and saturation of instabilities and the interplay between them.

SCATTERING INSTABILITIES IN UNDER-DENSE PLASMAS

The coupling between laser light and a plasma can be strongly influenced by the excitation of waves by a variety of instabilities [1, 2]. These instabilities can be most simply represented as the resonant coupling of the incident laser light into two other waves. Neglecting magnetic fields, these daughter waves are simply high frequency electron plasma waves (Langmuir), low frequency ion acoustic waves, and scattered electromagnetic waves. Two important instabilities are SRS and SBS. In SRS, the laser light can decay into a scattered light wave and an electron plasma wave. In the SBS instability, the two daughter waves are a scattered light wave and a low frequency ion acoustic wave. Other important instability are the two plasmon decay instability in which the laser light drives a pair of electron plasma waves, and the decay into an electron plasma wave plus an ion acoustic wave (or a purely growing ion fluctuation). In addition to the above processes, the laser light can produce density modulations which lead to either self-focusing or filamentation.

Stimulated Raman scattering

The Raman instability (SRS) represents the resonant decay of an incident photon into a scattered photon plus an electron plasma wave [1, 2, 3]. The frequency and wave number matching conditions are

$$\omega_o = \omega_s + \omega_{epw} \quad (1)$$

$$\mathbf{k}_o = \mathbf{k}_s + \mathbf{k}_{epw} \quad (2)$$

where $\omega_o(\omega_s)$ and $\mathbf{k}_o(\mathbf{k}_s)$ are the frequency and wave number of the incident (scattered) light wave, and $\omega_{epw}, (\mathbf{k}_{epw})$ is the frequency (wave number) of the electron plasma wave. Since $\omega_s \geq \omega_{epw}$, this process requires that $\omega_o > 2\omega_{epw}$; i.e. $n_e < 1/4n_c$. The wavenumber $\mathbf{k}$ varies from $\mathbf{k}_s \sim \mathbf{k}_o$ for $n_e \sim n_c/4$, to $\mathbf{k}_s \sim 2\mathbf{k}_o$ for $n_e \ll n_c/4$. Similarly, the frequency of the backscattered Raman lighth varies from $\omega_s \sim 1/2\omega_o$ at $n_e \sim n_c/4$, to $\omega_s \sim \omega_o$ at $n_e \ll n_c/4$. The frequency of the backscattered light is red shifted by the Bohm-Gross frequency [1]

$$\omega_{BG} = \omega_p \sqrt{1 + \frac{12\omega_o^2}{\omega_p^2} \frac{kT_e}{mc^2}} \quad (3)$$

where ω_p is the plasma frequency

$$\omega_p = \sqrt{\frac{4\pi n_e^2}{m_e}}, \quad (4)$$

T_e is the electron temperature, n_e is the electron density, and n_c is the critical electron density.

To illustrate the physics of the Raman instability, consider a light wave with electric field amplitude E_o and wavenumber $\mathbf{k}_o$ propagating through a plasma. Let the density be

rippled along k_o by the density fluctuation δn associated with an electron plasma wave. Since the electrons are oscillating with velocity $v_{os} = eE_o/m\omega_o$, a transverse current $\delta J = -ev_{os}\delta n$ is generated. With proper matching of the wavenumbers and frequencies, this transverse current generates a scattered light wave with amplitude δE . In turn, the scattered light wave beats with the incident light wave to produce a variation in the wave pressure: $\nabla(E_o/8\pi) = \nabla(E_o \cdot \delta E/4\pi)$. The variation in the wave pressure pushes plasma from regions of high pressure to regions of low pressure, generating density perturbations which can reinforce the initial fluctuation. The presence of this feedback loop leads to the Raman instability.

Maximum growth occurs when both the plasma wave and the down-shifted light wave are resonant. Taking $\omega = \omega_{epw} + i\gamma$, ignoring damping, and considering backscatter gives

$$\gamma = \frac{k v_{os}}{4} \sqrt{\frac{\omega_{epw}}{\omega_o - \omega_{epw}}} \quad (5)$$

where k is determined by the frequency and wavenumber matching conditions

$$k = k_o + \frac{\omega_o}{c} \sqrt{1 - \left(\frac{\omega_{epw}}{\omega_o}\right)^2} \quad (6)$$

For example, the SRS growth rate for $0.35 \mu\text{m}$ light irradiating a plasma with a density of $0.2 n_c$ at an intensity of 10^{15}W/cm^2 is $\gamma \simeq 1.6 \times 10^{13} \text{sec}^{-1}$, with a resulting growth time of approximately 0.07 ps.

Damping of both daughter waves introduces a pump threshold intensity for driving the instability. Damping can arise by both collisional and collisionless mechanisms and by inhomogeneity in the plasma. If we assume collisional damping of both the plasma wave and scattered light wave in an analogous fashion to inverse bremsstrahlung, net growth requires that

$$\gamma > \frac{1}{2} \left(\frac{\omega_{epw}}{\omega_o - \omega_{epw}} \right) \nu_{ei} \quad (7)$$

where ν_{ei} is the electron-ion collision frequency. Collisional stabilization can be particularly effective in high Z plasmas irradiated by short wavelength laser light. For example, consider $0.35 \mu\text{m}$ light incident on an Au plasma with an electron temperature of 2 keV, $Z=50$, and density of $2 \times 10^{21} \text{cm}^{-3} (\sim n_c)$. Then $\nu_{ei} \sim 2.4 \times 10^{13} \text{sec}^{-1}$ and the collisional threshold intensity would be about $5 \times 10^{14} \text{W/cm}^2$.

A density gradient limits the region over which three waves can resonantly interact. Propagation of wave energy out of this region then represents an effective dissipation which must be overcome by the incident pump. To illustrate the gradient threshold, consider Raman backscatter at a density $n_e \ll n_c$ in a plasma with a linear density profile with scale length, $L = [d/dx(\ln n)]^{-1}$. The Rosenbluth criterion for $e^{2\pi}$ convective growth is

$$\frac{\gamma^2}{|K v_1 v_2|} > 1 \quad (8)$$

where γ is the growth rate in a uniform plasma, $v_{1,2}$ denotes the group velocity of the two daughter waves and K is the gradient of the mismatch in the wavenumbers. Inserting the growth rate, we obtain

$$\left(\frac{v_{os}}{c}\right)^2 > \frac{2}{k_o L} \left(\frac{c}{v_{gs}}\right) \quad (9)$$

v_{gs} is the group velocity of the scattered light wave. For a plasma with $L \sim 10^4 \lambda_o$ irradiated with $0.35 \mu\text{m}$ laser light, this threshold intensity is only $3 \times 10^{14} \text{W/cm}^2$.

Stimulated Brillouin Scattering

The Brillouin instability represents the decay of an incident laser light wave into a scattered light wave plus an ion acoustic wave [1, 2, 3]. As can be readily seen from the frequency matching conditions, this instability is operative throughout the plasma with $n < n_c$. In addition, by the Manley-Rowe relations, nearly all the energy of the incident light can be transferred to the scattered light wave since the ion acoustic frequency is much less than the laser frequency. Hence, this instability can significantly degrade the absorption and/or modify its location.

The feedback leading to this instability is analogous to that already discussed for the Raman instability except now the density fluctuation is that associated with an ion acoustic wave. Maximum growth occurs when the ion acoustic wave and down-shifted light wave are both resonant. As an example, we consider backscatter ($k = 2k_o$) and neglect damping. For $\gamma \ll kc_s$ (weak coupling limit)

$$\omega = kc_s + i\gamma \quad (10)$$

$$\gamma = \frac{1}{2\sqrt{2}} \frac{k_o v_{os} \omega_{pi}}{\sqrt{\omega_o k_o c_s}}. \quad (11)$$

For $0.35 \mu\text{m}$ light with an intensity of 10^{14}W/cm^2 in a plasma with a density of $0.5n_c$ and an electron temperature of 2 keV, $\gamma \sim 2.5 \times 10^{12}$ and the growth time is about 0.4 ps. There is an instability threshold due to damping of the waves, so then net growth of Brillouin backscatter requires

$$\frac{v_{os}}{v_e} > 4 \frac{\nu_{ei}}{\omega_o} \frac{\nu_i}{\omega_{ia}} \quad (12)$$

where v_e is the electron thermal velocity and ν_{ei} is the ion-electron collision frequency and ν_i is the ion damping rate. This intensity threshold is typically quite low.

Inhomogeneity in either density, expansion velocity, or temperature of the plasma introduces a threshold for the instability. Rosenbluth criterion for $e^{2\pi}$ convective growth with a gradient in expansion velocity gives for the backscatter instability

$$\frac{v_{os}}{v_e} > \frac{16}{k_o L_\nu} \frac{n_{cr}}{n} \quad (13)$$

$$L_\nu = \left(\frac{1}{c_s} \frac{\delta v_e}{\delta x} \right)^{-1} \quad (14)$$

If we consider a plasma with a density of $0.5n_c$, an electron temperature of 2keV, and $L_\nu = 10^3 \lambda_o$, the gradient threshold intensity is then about 10^{14}W/cm^2 .

SPATIAL AND TEMPORAL EVOLUTION OF PARAMETRIC INSTABILITIES IN THE LINEAR APPROXIMATION

In this section, we will restrict ourselves to static homogeneous plasmas and will consider 1D interacting plasma slab model in which $x = L$ and $x = 0$ denote the points where the laser light, propagating from right to left, enters and goes out respectively. Such a 1D modelization is considered to describe properly the backscattering instabilities (beam to beam reflection).

Standard Three Wave Coupling

In the standard decay regime, the coupled mode equations take the form

$$\left(\frac{\partial}{\partial t} + V_0 \frac{\partial}{\partial x} + \nu_0\right)a_0 + \gamma_0 a_1 a_2 = S_0 \quad (15)$$

$$\left(\frac{\partial}{\partial t} + V_1 \frac{\partial}{\partial x} + \nu_1\right)a_1 - \gamma_0 a_0 a_2^* = S_1 \quad (16)$$

$$\left(\frac{\partial}{\partial t} + V_2 \frac{\partial}{\partial x} + \nu_2\right)a_2 - \gamma_0 a_0 a_1^* = S_2 \quad (17)$$

where a_0 , a_1 and a_2 stand for the amplitudes of the incoming laser light, the transverse and the longitudinal daughter waves respectively. These amplitudes have been normalized to have $a_0(L) = 1$ at the incoming point $x = L$ and at the intensity peak of the laser irradiation; γ_0 is the linear homogeneous growth rate of the parametric instability. The quantities V_α and ν_α denote the group velocity and the linear damping of wave α . In this section we will restrict ourselves to backscattering instabilities only, so that one has $V_0 < 0$, $V_1 > 0$ and $V_2 < 0$. The source terms S_α on the right-hand side of Eqs. (15), (16), (17) account for the thermal noise emission of wave α and are therefore stochastic functions in space and time. The statistical properties of the source functions S_α will be described in the next Subsection.

Equations (15,16,17) are correct only in the so-called standard decay limit where the envelope approximation is valid for each of the interacting waves. In the case of SRS the validity condition for this approximation can be written in physical units as

$$I_{14}\lambda^2 \ll 3.64 \cdot 10^5 \frac{(n/n_c)^{1/2}[1 - (n/n_c)^{1/2}]}{[\sqrt{1 - n/n_c} + \sqrt{1 - 2(n/n_c)^{1/2}}]^2} \quad (18)$$

and in the case of SRS it reads

$$I_{14}\lambda^2 \ll 1.5 \frac{n_c}{n} \sqrt{1 - n/n_c} (Z/A)^{1/2} [T_e + \frac{3T_i}{Z}(1 + K_2^2\lambda_D^2)]^{3/2} \sqrt{1 + K_2^2\lambda_D^2} \quad (19)$$

where T_e and T_i are the electron and ion temperature in units of keV; the quantity $K_2^2\lambda_D^2$ is given by $K_2^2\lambda_D^2 = 7.83 \cdot 10^{-3} T_e (n_c/n - 1)$. In both expressions (18, 19) $I_{14}\lambda^2$ denotes the laser flux in units of $10^{14} \text{ W cm}^{-2} \mu\text{m}^2$.

Statistical Properties of the Source Terms S_α

The source terms S_α on the right-hand side of Eqs. (15,16,17) are stochastic functions in space and time which have to reproduce properly the equilibrium statistical properties of fields a_α within the spectral domain of interest for the instability. These stochastic processes will have the form of white noises with the following statistical properties

$$\langle S_\alpha(x, t) \rangle = 0 \quad (20)$$

$$\langle S_\alpha(x, t) S_\alpha^*(x', t') \rangle = (2\pi)^2 \Sigma_\alpha^{1D} \delta(x - x') \delta(t - t') \quad (21)$$

where $\langle \rangle$ denotes statistical average. The space-time variables appearing in Eqs.(20, 21) are the slowly varying quantities corresponding to the envelope approximation. The 1D spectral density Σ_α^{1D} is determined from the condition that the equilibrium fluctuations

corresponding to Eqs. (15,16,17) in the limit $\gamma_0 \rightarrow 0$ are identical to the thermal equilibrium values. One obtains

$$S_{\alpha}^{1D} = \frac{\nu_{\alpha}}{(2\pi)^4} (T_{\alpha}^{eff}/\omega_{\alpha} N_0) K_{\alpha}^2 \Delta\Omega_{\alpha} \quad (22)$$

where $\mathbf{K}_{\alpha}$ is the wave vector of wave α associated with the usual three wave resonance conditions and $\Delta\Omega_{\alpha}$ is the solid angle in $\mathbf{K}_{\alpha}$ -space corresponding to the domain of interest for wave α . The effective temperature T_{α}^{eff} is given [4], for the transverse (TW) and electron plasma wave (EPW), by $T_{TW}^{eff} = T_{EPW}^{eff} = T_e$ and, for the ion sound wave (ISW), by $T_{ISW}^{eff} = T_e(1 + ZQ)/[1 + (ZT_e/T_i)Q]$, with $Q = (v_{the}/v_{thi}) \exp[-(ZT_e + 3T_i)/2T_i]$ where T_s and $v_{th,s} = (T_s/m_s)^{1/2}$ denote the temperature (in units of energy) and the thermal velocity of species s , respectively. The normalizing constant N_0 is the density of quanta for the pump wave

$$N_0 = \frac{n_{cr} m_e c^2}{4 \omega_{\alpha}} \left(\frac{v_{os}}{c} \right)^2 \quad (23)$$

where v_{os} is the electron quiver velocity $v_{os} = q_e E_0 / m_e \omega_{\alpha}$, E_0 denoting the maximum amplitude of the pump wave electric field at its entry point $x = L$. In the case of comparison with 1D particle in cell (PIC) simulations, the whole factor $T_{\alpha}^{eff} K_{\alpha}^2 \Delta\Omega_{\alpha}$ should be determined to match the equilibrium characteristics of the numerical scheme.

Space-Time Behavior in the Linear Regime

In the limit of small amplitude daughter waves, i.e. for $|a_1| \ll 1$ and $|a_2| \ll 1$, one can neglect the nonlinear pump depletion and solve the *linearized* system of Eqs. (15,16,17) where a_0 is now a *given* function of space and time. In the case of a rarefied ($n/n_c \ll 1$) and optically thin ($\nu_0 L/|V_0| \ll 1$) interacting plasma one can also neglect the linear dispersion and damping of the incident pump wave so that the pulse shape remains constant during a propagation through the interacting plasma. For the sake of simplicity we model the incident pulse of peak intensity I_{max} and characteristic duration t_p by a *square pulse* of intensity I_0 and duration t_p . The intensity I_0 is computed from the condition that the energy of the square pulse be identical to the incident energy. Accordingly, a_0 is equal to unity within the pulse interval $\max[0, L - |V_0|t] \leq x \leq \min[L, L - |V_0|(t - t_p)]$ and to zero elsewhere. This simple modeling is justified by numerical solutions to Eqs.(1) which show that varying the pulse shape does not lead to significantly different values of the time integrated reflectivity [5]. The solution of the system (16)-(17) with a square pulse as a pump wave can be carried out by means of the Green's function method after specifying the initial and boundary conditions [6]. For the initial stage of the pulse front propagation through the plasma, i.e. for $t < L/|V_0|$, we can assume vanishing initial conditions for the plasma fluctuations. Instead, we can impose the boundary conditions at the front of the pulse corresponding to the thermal equilibrium, because the laser pulse encounters a thermal plasma at its front. In the subsequent period of time $t > L/|V_0|$, one has to solve the system (16)-(17) taking for $a_{\alpha}(x, t)$ at $t = L/|V_0|$ the value reached at the end of the initial period and imposing thermal equilibrium for a_1 at $x = 0$. One can neglect the additional boundary condition corresponding to thermal equilibrium for a_2 at $x = L$ since, in the domain of parameters of concern in a picosecond time scale laser-plasma interaction, it would not lead to significantly different results from those obtained without it except within a negligibly thin boundary layer located around $x = L$. Following the procedure described above one obtains the fluctuations $\langle |a_{\alpha}|^2 \rangle$ as complicated integrals over space and time which can be calculated either asymptotically in the limit $\gamma_0 \rightarrow \infty$ (large gain factor limit) or perturbatively in the limit $\gamma_0 \rightarrow 0$ (small gain factor limit). From these two limits it is possible to design [6] tractable and uniform approximations to the fluctuations $\langle |a_{\alpha}|^2 \rangle$, i.e. expressions which are correct for both small and large values

of the gain factor. These expressions can be conveniently written in terms of the usual Thomson scattering reflectivity $R^{TS}(x)$ given by

$$R^{TS}(x) = (2\pi)^2 \left(\frac{\gamma_1^{AD}}{\nu_1} + \frac{\gamma_2^{AD}}{\nu_2} \right) \frac{\gamma_0^2 x}{2V_1^2} \quad (24)$$

It is interesting to discuss the solutions to Eqs.(16-17) as functions of the normalized damping, ν_2/γ_0 of daughter wave 2. In the strongly damped regime defined by the inequality $\nu_2/\gamma_0 \gg (1 + V_1/|V_0|)^{-1/2}$, one obtains the following expressions

$$\langle |a_1(x, t)|^2 \rangle = \langle |a_1|^2 \rangle_{eq} + \quad (25)$$

$$R^{TS}(l_{eff}) \frac{1}{\sqrt{1 + \pi G(l_{eff})}} \left[\frac{\exp G(l_{eff}) - 1}{G(l_{eff})} \right]$$

$$\langle |a_2(x, t)|^2 \rangle = \langle |a_2|^2 \rangle_{eq} + \quad (26)$$

$$R^{TS}(V_1 \frac{1 - \exp[-\frac{2\nu_2}{|V_0|}(x - L + |V_0|t)]}{2\nu_2}) \frac{\exp G(l_{eff})}{\sqrt{1 + \pi G(l_{eff})}}$$

where $\langle |a_\alpha|^2 \rangle_{eq}$ denotes the equilibrium fluctuations of wave α corresponding to the limit $\gamma_0 \rightarrow 0$. The gain factor $G(x)$ and the effective length l_{eff} are given by

$$G(x) = \frac{2\gamma_0^2 x}{\nu_2 V_1} \quad (27)$$

$$l_{eff} = \text{Min}[x, \frac{V_1(x - L + |V_0|t)}{|V_0| + V_1}]. \quad (28)$$

From the expression of the gain factor $G(l_{eff})$ appearing in Eqs.(26,27), one can see that the interaction is characterized by two different phases. The first phase takes place in the time interval $(L-x)/|V_0| < t < L/|V_0| + x/V_1$ and corresponds to exponentially growing in time fluctuations (and logarithmic corrections). The wave amplitude instantaneous growth rate $\gamma \equiv \gamma_0^2/[\nu_2(1 + V_1/|V_0|)]$ may be interpreted as corresponding to a spatial amplification in the frame moving with the laser front. One has $G(l_{eff}) = 2K X(t)$ where $K \equiv \gamma_0^2/[\nu_2(V_1 + |V_0|)]$ denotes the spatial amplification growth rate in the moving frame, and where $X(t) = x - L + |V_0|t$ is the instantaneous distance between the pulse front and the point x . The second phase takes place in the time interval $L/|V_0| + x/V_1 < t < t_p + (L-x)/|V_0|$ and corresponds to the spatial amplification in the laboratory frame. During this phase the gain factor $G(l_{eff})$ remains equal to the usual spatial amplification factor $G(l_{eff}) = (2\gamma_0^2 x)/(\nu_2 V_1)$ and the fluctuations $\langle |a_\alpha|^2 \rangle$ do not longer evolve in time. In the complementary, weakly damped regime defined by the inequality $\nu_2/\gamma_0 \ll (1 + V_1/|V_0|)^{-1/2}$, one obtains the expressions

$$\langle |a_1(x, t)|^2 \rangle = \langle |a_1|^2 \rangle_{eq} + \quad (29)$$

$$R^{TS}(l_{eff}) \frac{\exp \Gamma(\tau_{eff})}{\sqrt{[1 + q_1^2 G(\lambda_{eff})][1 + r_1^2 G(\lambda_{eff})]}} \left[\frac{\exp G(\lambda_{eff}) - 1}{G(\lambda_{eff})} \right]$$

$$\langle |a_2(x, t)|^2 \rangle = \langle |a_2|^2 \rangle_{eq} + R^{TS}\left(\frac{V_1(x - L + |V_0|t)}{|V_0|}\right) \left[\frac{\exp G(\lambda_{eff}) - 1}{G(\lambda_{eff})} \right] \frac{\exp \Gamma(\tau_{eff})}{\sqrt{[1 + q_2^2 G(\lambda_{eff})][1 + r_2^2 G(\lambda_{eff})]}} \quad (30)$$

where the gain factor $G(x)$ and the damping factor $\Gamma(t)$ are given by

$$G(x) = \frac{2\gamma_0 x}{V_1} \quad (31)$$

$$\Gamma(t) = 2\nu_2 t \quad (32)$$

and where the effective length λ_{eff} and time τ_{eff} are given by

$$\lambda_{eff} = 2\sqrt{V_1 \text{Min}[x, \frac{V_1(x-L+|V_0|)}{2(|V_0|+V_1)}] \text{Max}[(t-\frac{x}{V_1}-\frac{L}{|V_0|}), \frac{1}{2}(t+\frac{x}{|V_0|})]} \quad (33)$$

$$\tau_{eff} = \text{Max}[(t-\frac{x}{V_1}-\frac{L}{|V_0|}), \frac{1}{2}(t+\frac{x}{|V_0|})]. \quad (34)$$

The quantities q_α and r_α are defined by

$$\begin{aligned} q_1 &\equiv \sqrt{\frac{\pi}{2}} \lambda_{eff} / \text{Min}[x, \frac{V_1(x-L+|V_0|)}{2(|V_0|+V_1)}] \\ q_2 &\equiv \sqrt{\frac{\pi}{2}} (t + \frac{x-L}{|V_0|}) / \text{Max}[(t-\frac{x}{V_1}-\frac{L}{|V_0|}), \frac{1}{2}(t+\frac{x}{|V_0|})] \\ r_1 = r_2 &\equiv \sqrt{\frac{\pi \text{Max}[0, t-(\frac{x}{V_1}+\frac{L}{|V_0|})x-|V_0|]}{2 \text{Min}[2(t+\frac{x}{|V_0|}), t+(\frac{x}{V_1}+\frac{L}{|V_0|})x-|V_0|]}} \end{aligned}$$

The validity domain of the expressions (29), (30) is actually restricted by the inequality $t < t^*$, where $t^* \sim \text{Min}[L/|V_2|, (\gamma_0^2 L)/(\nu_2^2 V_1)]$ is the characteristic time at which the fluctuations behave as predicted by the usual time asymptotic analysis (namely, they reach steady state in the convective regime or grow in time in the absolute regime). It is important to realize that it takes the time t^* to distinguish between absolute and convective evolution. In the case of short pulse experiments in preformed plasma, the pulse duration t_p is typically much shorter than t^* so that the concepts of absolute or convective instabilities are irrelevant for the proper description of the scattering instabilities [7, 8]. In this case, the *whole interaction* takes place in the transient stage of the evolution which is correctly given by the Eqs. (29), (30). From the expression of the gain factor $G(\lambda_{eff})$ appearing in Eqs. (29), (30), one finds that the transient stage ($t < t^*$) of the interaction is characterized by two different regimes. The first phase takes place in the time interval $(L-x)/|V_0| < t < L/|V_0| + (x/V_1)(2+V_1/|V_0|)$ and corresponds to exponentially growing in time fluctuations (and logarithmic corrections). The wave amplitude instantaneous growth rate $\gamma \equiv \gamma_0/(1+V_1/|V_0|)^{1/2}$ can again be interpreted as corresponding to a spatial amplification in the frame moving with the pulse front. One indeed finds that $G(\lambda_{eff}) = 2KX(t)$, where the spatial amplification growth rate K in the moving frame is now computed in the weakly damped limit, namely $K = \gamma_0/\sqrt{(V_1+|V_0|)(V_2+|V_0|)} \simeq \gamma_0/|V_0|\sqrt{1+V_1/|V_0|}$. One should stress that neglecting the pulse propagation amounts to taking $|V_0| = +\infty$, which would lead to overestimate of the instantaneous growth rate by the factor $\sqrt{2}$ in the underdense limit $|V_0| \sim |V_1| \sim c$. The second phase of the evolution takes place in the subsequent time interval $L/|V_0| + (x/V_1)(2+V_1/|V_0|) < t < t_p + (L-x)/|V_0|$ and is associated with fluctuations behaving like $\exp G_{tot}$ (plus logarithmic corrections), where the total gain factor $G_{tot} \equiv G(\lambda_{eff}) + \Gamma(\tau_{eff})$ is given by

$$G_{tot} = 4\gamma_0 \sqrt{\frac{x}{V_1}(t - \frac{x}{V_1} - \frac{L}{|V_0|})} - 2\nu_2(t - \frac{x}{V_1} - \frac{L}{|V_0|}) \quad (35)$$

which corresponds to a time behavior which is not of an exponential type [6, 7, 9]. The effective wave amplitude instantaneous growth rate $\gamma_{eff}(t)$ defined by $\gamma_{eff}(t) \equiv \frac{1}{|a_\alpha|} \frac{\partial |a_\alpha|}{\partial t}$ is a decreasing function of time which reads

$$u_H(t) = \omega \sqrt{t - x/V_1 - L/|V_0|} \quad (36)$$

At a given point x , the finite value of the pulse duration t_p limits the extension of the interaction period to $t < t_p + (L - x)/|V_0|$. All the results we have discussed above describe properly the space-time evolution of the instability provided t is taken *inside* the interaction period $(L - x)/|V_0| < t < t_p + (L - x)/|V_0|$ (and, in the weakly damped regime, provided the inequality $t_p < t^*$ is fulfilled). In the opposite case $t > t_p + (L - x)/|V_0|$, the point x is located in the wake of the pulse and the fluctuations at this point correspond to freely propagating waves 1 and 2 which have been left behind the pulse and which suffer convective and dissipative losses, respectively.

NONLINEAR EVOLUTION OF STIMULATED RAMAN SCATTERING

In our review of scattering instabilities on the short time scale we discuss physical processes which are not affected by the hydrodynamic evolution of laser produced plasmas. Our simplified theoretical models treat plasma as a homogeneous medium with the fixed density and temperatures, but include several nonlinear terms describing complicated wave evolution. It has become apparent that the proper understanding of experimental results must involve nonlinear physics even on the scale of several picoseconds. Fast growing SRS can easily reach nonlinear regime within this time period. The primary nonlinear effects related to SRS correspond to coupling between Langmuir and ion acoustic waves, which results in parametric decay cascade and/or Langmuir collapse. The broad spectra of Langmuir and ion acoustic fluctuations produced during this nonlinear SRS evolution can have in turn dramatic effect on other, slowly growing scattering processes including Brillouin, forward Raman and non resonant interactions involving anti-Stokes components of electromagnetic waves. Review of recent theoretical results on the saturation of SRS and the effect of SRS produced non thermal fluctuation levels on the secondary scattering processes will be given in the next two subsections.

Theoretical Model

Since 1987 [9, 10, 11] the nonlinear evolution of SRS has been studied within the context of the model, which includes the system of Zakharov and electromagnetic wave equations. The Zakharov equations describe coupling between Langmuir and sound waves. They include physics related to modulational instability, parametric decay instability (PDI), PDI cascade, and the effects related to strong Langmuir turbulence and collapse. The ponderomotive force of electromagnetic waves drives the SRS resonant Langmuir wave at the relatively large wave number, making a modulational instability the unlikely mechanism of nonlinear SRS saturation [11]. An important nonlinear coupling described by this model [9, 10, 11] corresponds to interaction between stimulated Brillouin scattering and SRS. SBS produces sound waves, which can affect nonlinear evolution of Langmuir waves by destroying SRS matching conditions or by triggering the localization of Langmuir fields. The actual scenario of SRS/SBS coupling depends on the linear growth of both instabilities and initial noise levels of plasma fluctuations (cf. previous section). We have identified the range of plasma parameters, relevant for many current short wavelength laser-plasma interaction experiments, where SBS does not grow in time and has negligible effect on the evolution of SRS at early times. One can achieve similar conditions for even wider range of parameters but for the short laser pulses, when slowly growing SBS has no time to reach large levels of reflectivity during the laser pulse duration. By eliminating at first strong SBS we are able to make progress in our understanding of nonlinear SRS behavior. The

correct theory of SRS saturation is still uncertain and the existing models always lead to much higher values of reflectivity as compared to experimental observations. Next section will address some of the questions related to the saturation of SRS in the context of 2D modeling of the instability. Unfortunately analysis of the two dimensional models of SRS is still impractical from the theoretical and numerical point of view. Our conclusions are based on the numerical solutions to the system of Zakharov and electromagnetic wave equations

$$i\left(\frac{\partial E}{\partial t} + \nu_L \cdot E\right) + \frac{3}{2} \frac{v_e^2}{\omega_p} \frac{\partial^2 E}{\partial x^2} - \frac{\omega_p}{2} N E \quad (37)$$

$$= \frac{\pi e n_0^2}{\omega_p} \frac{\partial}{\partial x} (\Psi_0 \Psi_1^*) + S_L(x, t),$$

$$\frac{\partial^2 N}{\partial t^2} + 2\nu_A \cdot \frac{\partial N}{\partial t} - c_s^2 \frac{\partial^2 N}{\partial x^2} \quad (38)$$

$$= \frac{\partial^2}{\partial x^2} \left(\frac{|E|^2}{16\pi n_0^4 m_i} + \frac{Z m_e}{4 m_i} |\Psi_0|^2 \right) + S_A(x, t),$$

$$i\left(\frac{\partial \Psi_0}{\partial t} + \nu_0 \Psi_0\right) + \frac{c^2}{2\omega_0} \frac{\partial^2 \Psi_0}{\partial x^2} - \frac{\omega_p^2 - \omega_0^2}{2\omega_0} \Psi_0 \quad (39)$$

$$= \frac{\omega_p^2}{2\omega_0} N \Psi_0 - \frac{e}{4\omega_0 m_e} \left(\Psi_1 \frac{\partial E}{\partial x} \right),$$

$$i\left(\frac{\partial \Psi_1}{\partial t} + \nu_1 \Psi_1\right) + \frac{c^2}{2\omega_1} \frac{\partial^2 \Psi_1}{\partial x^2} - \frac{\omega_p^2 - \omega_1^2}{2\omega_1} \Psi_1 \quad (40)$$

$$= \frac{\omega_p^2}{2\omega_1} N \Psi_1 - \frac{e}{4\omega_1 m_e} \Psi_0 \frac{\partial E^*}{\partial x},$$

where $v_e = (T_e/m_e)^{1/2}$ and $c_s = (ZT_e/m_i)^{1/2}$. Equations (37) - (40), provide improved approximation to the description of scattering instabilities as compared to Eqs. (15) - (17). They include several nonlinear terms and describe variations of the fields not in terms of slowly varying amplitudes a_0, a_1, a_2 but in terms of normalized ion density $N(x, t)$, time enveloped electrostatic field amplitude E and electron quiver velocities Ψ_0, Ψ_1 in the electric fields of electromagnetic pump and SRS backscattered wave, respectively. The complete electrostatic field reads $E = (1/2)(E(x, t) \exp(-i\omega_p t) + c.c.)$, and the electron quiver velocity in the electric fields of electromagnetic waves has the following form $V_e = (1/2) \sum_{r=0,1} (\Psi_r \exp(-i\omega_r t) + c.c.)$, where $\omega_1 = -\omega_p + \omega_0$. By enveloping E, Ψ_0, Ψ_1 with respect to high frequency variations we keep in Eqs. (37)-(40) terms varying on the scale of ion motion. There is, however no restriction on the spatial variations of the physical quantities. As before ν_0, ν_1 describe damping coefficients of the pump and scattered light, respectively, and ν_L, ν_A stand for linear damping operators of Langmuir fields and density disturbances, respectively. The latter contain Landau damping and effect of electron-ion collisions for electrostatic fields. The source terms S_L, S_A are added in order to sustain the level of fluctuations of SRS active modes consistent with the thermal equilibrium values (cf. section *Statistical Properties of the Source Terms S_α*) in the absence of the pump. The Zakharov equations are part of Eqs. (37)-(38). They are driven by the ponderomotive forces proportional to $\Psi_0 \Psi_1^*$ (37) and $|\Psi_0|^2$ (38). The first one, which occurs on the right hand side of Eq. (37) drives SRS, and the second one is responsible for SRS interaction. The validity conditions for Zakharov equations, which are always satisfied in our simulations, are given by $|E|^2/4\pi n_0^2 T_e < 1$ and $|N| < 1$. We have analyzed the wide range of plasma parameters with the common characteristic that the length of the system L is always much longer than the damping length of Langmuir waves V_L/ν_L . It is one of the important properties of short wavelength laser produced

Figure 1: (a) Typical time history of the SRS reflectivity. Different types of behavior correspond to : (1) linear growth, (2) first saturation due to the PDI, (3) transient nonlinear regime characterized by the development of the PDI cascade and its disruption, (4) final saturation. (b) SRS reflectivity and enhanced Brillouin reflectivity R_B from the numerical solutions to Eqs. (37)-(40) for the following parameters: $n_0/n_c = 0.2$, $T_e = 1.5 \text{ keV}$, $ZT_e/T_i = 8$, $I = 10^{14} \text{ W/cm}^2$, $L = 60 \mu\text{m}$.

plasmas, where collisional damping on Langmuir waves plays an important role.

SRS Reflectivity

It has been shown in Ref. [12] that the main features of the one dimensional Zakharov - electromagnetic wave equations model (Eqs. (37) - (40)) can be divided into transient processes and the final saturation regime, which is independent of the initial conditions (cf. Fig. 1(a)).

During the transient evolution SRS grows at first linearly (phase (1) in Fig. 1) and saturates due to the onset of parametric decay instability of resonantly driven by SRS Langmuir wave (phase (2) in Fig. 1(a)). SRS never reaches the linear asymptotic regimes related to absolute or convective instabilities because the nonlinear effects due to PDI occur earlier. The linear growth rate is given by the gain factor (36) from the transient regime of the instability. The first nonlinear saturation corresponds to a strong flash of backscattered radiation (phase (2) in Fig. 1(a)). Next in the transient nonlinear regime ((3) in Fig. 1(a)) we observe development of the PDI cascade and its disruption due to different spatio-temporal evolution of SRS and PDI. In the asymptotic regime ((4) in Fig. 1(a)) SRS reflectivity saturates at the level which is well approximated by the scaling formula

$$R = \frac{8}{27} \gamma_R^2 L^3 \nu_A^{k_A} \frac{(k_L v_e)^2 \omega_p}{\omega_0^2 \omega_A V_0 V_1^2}. \quad (41)$$

In Eq. (41) L is the plasma length, $\nu_A^{k_A}$ represents the Landau damping coefficient of the ion acoustic wave (ω_A, k_A) produced during PDI of the SRS driven Langmuir wave

(ω_L, k_L) ; V_0 and V_1 stand for the group velocities of the pump (ω_0, k_0) and backscattered (ω_1, k_1) waves, respectively. The homogenous SRS growth rate coefficient is given by $\gamma_{0R} = (1/4)v_0 k_L (\omega_p^2 / \omega_1 \omega_L)^{1/2}$, where v_0 is the electron quiver velocity amplitude in the electric field of the pump wave. The waves participating in SRS interaction satisfy the usual resonant matching conditions $\omega_0 = \omega_L + \omega_1$, and $k_0 = k_L + k_1$. The secondary process of the PDI corresponds to the resonant conditions: $\omega_L = \omega_L^{(1)} + \omega_A$ and $k_L = k_L^{(1)} + k_A$, where $(\omega_L^{(1)}, k_L^{(1)})$ characterizes the first component of the Langmuir wave cascade driven by PDI. The matching conditions for PDI lead to the following solutions for the wave numbers involved in the first stage of the cascade: $k_L^{(1)} = -k_L + \Delta k$, $k_A^{(1)} = 2k_L - \Delta k$, where $\Delta k = (2/3)k_D \sqrt{Z m_e / m_i}$.

The scaling expression for the reflectivity (41) has been derived without pump depletion effects. Therefore it is only valid for small values of R . It also gives reasonable predictions for the plasma density $n_0^e/n_c \leq 0.23$. Note that expression (41) does not depend on the damping coefficient, $\nu_L^{k_L}$ of Langmuir waves. The particular mechanism of Langmuir wave dissipation, i.e. for example Landau and collisional damping, or increased dissipation due to collapse processes (cf. Ref. [13]), does not directly affect SRS reflectivity but controls the spectrum of plasma waves and hot electron production. Equation (41) can be rewritten in more explicit form

$$R = 6.6 \cdot 10^{-7} I_{14} \lambda^2 \left(\frac{L}{\lambda} \right)^3 T_e^2 \left(\frac{\nu_A^{k_A}}{\omega_A} \right) \Xi(n_0^e/n_c), \quad (42)$$

where electron temperature T_e is expressed in keV , and Ξ is a smoothly varying function of the electron density n_0^e/n_c ,

$$\Xi = \frac{(n_0^e/n_c) [1 - (n_0^e/n_c)^{1/2}] [\sqrt{1 - n_0^e/n_c} + \sqrt{1 - 2(n_0^e/n_c)^{1/2}}]^4}{\sqrt{1 - n_0^e/n_c} [1 - 2(n_0^e/n_c)^{1/2}]} \quad (43)$$

The expressions (41), (42) for the SRS reflectivity have been derived from scaling arguments [12], which includes an estimate of the Langmuir wave amplitude E_L participating in the SRS interaction,

$$E_L = E_{th} \sqrt{\frac{2}{3}} A, \quad (44)$$

where $A = \gamma_{0R}^2 L / \nu_L^{k_L} |V_1|$ is the SRS amplification coefficient and E_{th} is the threshold of PDI

$$\frac{E_{th}}{\sqrt{4\pi n_0^e T_e}} = \frac{16 \nu_L^{k_L} \nu_A^{k_A}}{\omega_p \omega_A}. \quad (45)$$

From expressions (44), (45) we can derive validity condition for the Zakharov equations in terms of relevant for SRS parameters

$$1.22 \cdot 10^{-3} I_{14} \lambda^2 \left(\frac{\nu_A^{k_A}}{\omega_A} \right) \left(\frac{L}{\lambda} \right) \Xi_E(n_0^e/n_c) < 1, \quad (46)$$

where Ξ_E is the function of electron density,

$$\Xi_E = \frac{[\sqrt{1 - n_0^e/n_c} + \sqrt{1 - 2(n_0^e/n_c)^{1/2}}]^2}{\sqrt{[1 - 2(n_0^e/n_c)^{1/2}]}} \quad (47)$$

The limit (46) has been derived from the estimate of the electrostatic field amplitudes in the asymptotic saturation regime, where the fluctuations have already reached lower

levels as compared to transient regimes ((2) and (3) from Fig.1(a)). Therefore (46) provides only necessary condition for the validity of our numerical calculations and it does not exclude possibility of a strong enhancement of electric fields ($|E|^2/4\pi n_0^e T_e > 1$) in the transient regimes. Our model does not include nonlinear kinetic effects and cannot properly describe evolution of strong electrostatic fields. Finally we have considered plasma parameters, where SRS cannot compete with SRS, i.e. approximately the gain of SRS $G_B = \gamma_{0B}^2 L / V_0 \nu_A^{2k_0}$ is less than unity. In fact this limitation is stronger than necessary, in particular for short pulse interaction experiments when the long time behavior of SRS is irrelevant.

Laser Light Scattering from SRS Produced Fluctuations

The nonlinear evolution of Langmuir waves, which is responsible for the saturation of SRS, results in a broad spectra of the electrostatic fluctuations. The PDI cascade, localization of Langmuir waves and spatial evolution of the instabilities produce enhanced fluctuation levels which in turn can scatter the electromagnetic pump via Brillouin, forward Raman, and/or blue-shifted Raman interactions. These processes can be particularly important for the short pulse laser plasma interaction experiments, where fast growing SRS and subsequent PDI of Langmuir waves can enhance ion fluctuations by several orders of magnitude on the time scale when, for example, ordinary SRS has no chance to grow. We will concentrate in this review on the description of the enhancement of the Brillouin scattering by the backward SRS (cf. Ref. [14]), which can produce the most pronounced and unique effects in experiments.

Figure 1(b) illustrates the time evolution of SRS reflectivity together with the enhanced Brillouin reflectivity R_B . For the conditions of this run ($n_0/n_e = 0.2$, $T_e = 1.5keV$, $ZT_e/T_i = 8$, $I = 10^{14}W/cm^2$, $L = 60\mu m$, $\lambda = 0.5\mu m$) we do not observe a temporal growth of the ordinary stimulated Brillouin scattering. The sudden onset of the backscattered Brillouin signal takes place during the intermediate nonlinear evolution of SRS (phase (3) from Fig. 1(a)) and it is related to the broad ion acoustic spectrum of fluctuations produced at that time. It coincides with the development of high components of the PDI cascade and with the formation of continuum in the fluctuation spectra. The level of Brillouin scattering depends on the enhancement of the ion acoustic spectrum. In particular on the level of fluctuations produced at the wave number $k = k_A^B = 2k_0$, which participates in the Brillouin interaction. The extend of the enhanced region in the ion wave spectrum

scales as the overall number of steps in the PDI cascade. It is proportional to the ratio E_L^2/E_{th}^2 , i.e. for the saturated SRS the number of steps in the PDI cascade is simply $\frac{2}{3}A$ (cf. Eq. 44). Using this estimate, we may derive simple strong enhancement criteria for Brillouin scattering, in terms of the backward SRS amplification coefficient

$$\frac{2}{3}A > |k_A^B - k_A| / 2\Delta k, \quad (48)$$

where $2\Delta k$ is the separation between ion modes produced by PDI. Figure 2(a) shows enhanced spectra corresponding to the late moment in time of the run illustrated in Fig. 1(b). The enhanced part of the spectrum is dominated by broad, but distinct peaks corresponding to modes produced by PDI. The value of $k = k_A^B = 0.22k_D$ is located very close to the fourth component of the PDI cascade which is a mode propagating in the opposite direction to the pump. Electromagnetic radiation which is scattered due to interaction with this mode would produce primarily the blue shifted component observed in the frequency spectrum (cf. Fig. 2(b)). The contour plot (Fig. 2(b)) displays the time evolution of the spectral intensity of Brillouin backscattered radiation normalized to the pump intensity at $\omega = \omega_0$. The shift of red and blue maxima corresponds approximately to the ion acoustic

Figure 2: Numerical simulations for the parameters of Fig. 1(b). (a) normalized Fourier spectrum of density fluctuations at time $t = 53.29$ ps (cf. Fig. 1(b)); (b) contour plot of the Brillouin backscattered spectral intensity normalized to pump intensity.

frequency $2k_0 c_s$. Since the enhanced part of ion wave spectrum is dominated by almost discrete peaks small change in plasma parameters can move k_A^B component closer to the mode propagating along the pump and produce strongly enhanced red shifted spectrum. There is also a continuous part of the density spectra, which will enhance blue and red shifted electromagnetic waves to the similar levels. The presence of blue and red shifted components in the spectra of the scattered light is a characteristic feature, which can help to identify this process experimentally.

SCATTERING INSTABILITIES INVOLVING ION WAVES

In the previous Section we have discussed the stimulated Raman scattering which is a resonant instability involving electron plasma waves. We turn our attention now to parametric processes corresponding to the interaction between electromagnetic and ion waves. They are: stimulated Brillouin scattering (SBS), filamentation (FI) and self-focusing (SFI) instabilities [1, 2]. The ion wave instabilities are of a major concern for the laser fusion schemes, because of their potential capability to reflect a substantial part of the incident laser energy, spoil the compression symmetry, and stimulate other undesirable nonlinear processes [15, 16]. Experiments and theoretical analysis of these instabilities on the short time scale in well controlled plasmas [5] are helping to understand physics of these processes, including their mutual interaction and nonlinear evolution. In this review we will emphasize the common origin of SBS, FI and SFI, related to the ion dynamical response to ponderomotive forces of an electromagnetic and/or a thermal origin. The differences correspond to the nature of ion response: SBS excites resonant ion acoustic waves, while FI and SFI produce non resonant ion density disturbances.

The existing theories of ion wave scattering instabilities fail to explain such commonly

observed phenomena as low level saturation of backscattered SRS reflectivity, pulsations of reflectivity, side scattering, and recent observations of forward SRS [17]. Also the growing experimental evidence indicates [18, 19, 20] that many parametric instabilities are excited in narrow localized regions of the plasma and hence, these instabilities may be the result of filamentation or self focusing instabilities.

The lack of agreement between the experimental and theoretical studies of the ion wave parametric instabilities is partly related to the complicated hydrodynamic behavior of plasmas in a nanosecond time scale, which had never been properly accounted in the theory. The short laser pulse experiments, unaffected by the plasma hydrodynamic evolution, allow for the first time a meaningful comparison between theory and experimental observations. The other origin of the discrepancy are the complicated spatial profiles of the laser produced plasmas, due to inhomogeneities of laser beam intensity over the spot size. The new experimental technique of laser beam smoothing [21] enables us to overcome this difficulty and to perform plasma with well controlled and rather plain profiles. Therefore the combination of the laser beam smoothing with the subpicosecond plasma diagnostic and short laser pulses opens quite new possibilities in the understanding of ion wave parametric instabilities.

In this section we will describe the main characteristics of the ion mode parametric instabilities at the linear stage and possible mechanisms of their nonlinear saturation, together with the results of numerical simulations.

Basic Equations

We consider an homogeneous plasma slab in cartesian geometry with the ion density n_0^i and electron density $n_0^e = Zn_0^i$, where Z is the average ion charge. A beam of s-polarized electromagnetic wave of frequency ω_0 is incident on the plasma slab at $x = 0$ and propagates along x -direction. The time enveloped of the electric field E of the transverse electromagnetic wave polarized along z -axis, and the ion density perturbation $\delta N = n_i/n_0^i - 1$, satisfy the system of equations derived from the standard set of the Maxwell and plasma hydrodynamic equations. Neglecting the dispersion and nonlinear effects on ion acoustic wave propagation, one obtains [22, 23]:

$$\left(2i\frac{\omega_0}{c^2}\frac{\partial}{\partial t} + \nabla^2 + k_0^2 + i\frac{\nu_{ei}\omega_0}{c^2}\frac{n_0^e}{n_c}\right)E = \frac{\omega_0^2}{c^2}\frac{n_0^e}{n_c}E\delta N, \quad (49)$$

$$\left(\frac{\partial^2}{\partial t^2} + 2\hat{\Gamma}_s\frac{\partial}{\partial t} - c_s^2\nabla^2\right)\delta N = \frac{Z}{16\pi m_i n_c}\nabla^2|E|^2, \quad (50)$$

where $k_0 = (\omega_0/c)\sqrt{1 - n_0^e/n_c}$ is the wavenumber of the laser pump, the operator $\hat{\Gamma}_s$ models the damping of ion density fluctuations, $c_s = \sqrt{(ZT_e + 3T_i)/m_i}$ is the ion sound speed, c is the speed of light, $n_c = m_e\omega_0^2/4\pi e^2$ is the critical density, and ν_{ei} is the electron-ion collision frequency describing the absorption of the electromagnetic waves.

The two equations (49), (50) describe the coupling between high frequency electromagnetic waves and low frequency plasma density via the ponderomotive force term on the right hand side of Eq.(50). It is known, that in addition to ponderomotive effects the ion waves can be excited by the inhomogeneous plasma heating. Thermal effects can be important for SRS and filamentation instabilities especially in the regime of parameters where the electron mean free path is comparable with the characteristic wavelength of ion fluctuations. The importance of these processes has been recognized recently as a result of the better understanding of nonlocal thermal transport effects [24],[25]. However in this review we will deal only with the ponderomotive case and leave for the future the discussion of the thermal instabilities, which we are currently studying within the context

of a two dimensional model similar to Eqs. (49), (50)

Linear Theory of Ion-Wave Scattering Instabilities

Equations (49)–(50) can be reduced to Eqs. (15), (17), which describe evolution of the resonant three wave coupling processes like for example SBS. With the straightforward generalization to the two dimensional planar geometry these equations describe evolution of slowly varying amplitudes of the waves participating in backward or side scattered SBS. The linear analysis of these interactions lead to the results summarized in a previous section. The scattering in the forward direction grows slower than the scattering producing electromagnetic radiation propagating in the backward direction. However because the forward SBS involves ion waves of the relatively smaller wave numbers they are weakly damped and the scattering can reach higher level in the asymptotic linear regime characterized by stationary amplification coefficient (27).

Our model (49), (50) describes sound waves in the linear approximation, but accounts for the competition between the backward and side scattered SBS, as well as the filamentation instability. In order to understand the coexistence of these instabilities, one must describe not only the linear growth but also the initial levels of fluctuations. The results of our 2D model show that the thermal levels of sound waves underestimate the real noise amplitudes, in particular for the stationary fluctuations with the transverse wave vectors to the direction of laser pump propagation. These fluctuations, which are increased by the incident laser beam participate in the filamentation instability. They can enhance this relatively slowly growing interaction to the level were it competes with SBS, even on the scale of short pulse laser experiments.

The expression for the filamentation growth rate follows from Eqs. (49), (50) if we assume that the scattered waves propagating obliquely to the pump have no frequency shift ($k_{\perp} \ll k_0$):

$$E(\mathbf{r}, t) = [E_0 + E_- \exp(-ik_{\perp} \mathbf{r}_{\perp}) + E_+ \exp(ik_{\perp} \mathbf{r}_{\perp})] \exp(ik_0 x).$$

The filamentation instability, similarly to SBS, exhibits complicated spatial-temporal behavior before the convective saturation. In the transient stage it is difficult to separate the filamentation from the forward SBS, because their characteristic growth rates are comparable with the Brillouin frequency shift and therefore, both of them lead to the non stationary disturbances of the pump beam. However, the stationary amplification coefficient for FI differs considerably from that of SBS (27). According to Ref. [15, 27] in the case of the thermal filamentation and classical heat transport ($k_{\perp} \lambda_e < 1$, where λ_e is the electron mean free path) the amplification takes place for $k_{\perp} < K_{cr}^{(th)}$ where $K_{cr}^{(th)} = \sqrt{k_0/\lambda_e} \sqrt{n_0^e |E_0|^2 / 4\pi n_e^2 T_e}$. In that case the coefficient of the exponential amplification in the plasma slab of the length L has the following form :

$$G_{th} \approx (K_{cr}^{(th)})^2 L / k_0. \quad (51)$$

The ponderomotive FI takes place for the small-scale inhomogeneities ($k_{\perp} \ll \lambda_e$) when $k_{\perp} \approx K_{cr}^{(p)} \approx k_0 \sqrt{n_0^e |E_0|^2 / 4\pi n_e^2 T_e}$. The corresponding amplification coefficient is $G_p \approx (K_{cr}^{(p)})^2 L / k_0$.

All these expression for the filamentation instability are relevant if the laser beam width a is much larger than the characteristic filamentation scale, $k_{\perp} a \gg 1$, and FI results in splitting of the whole beam into several nearly independent filaments. If the beam size is smaller than the filamentation scale than the whole beam self-focusing (SFI) can take place. The condition of the thermal SFI onset is $G_{th} > 1$ for $a \gg \lambda_e$. The ponderomotive SFI is in effect for narrower beams $a \ll \lambda_e$ and starts to work if $K_{cr}^{(p)} L > k_0 a$. In the

particular case of the Gaussian laser beam in a homogeneous plasma the interaction length L in the last expression equals to the beam Rayleigh length $k_0 a^2$ where a is the minimum radius of the beam. In that case the condition for the onset of the self focusing can be written as $\alpha_{sf} > 1$ where

$$\alpha_{sf} = (k_0 a)^2 (n_e |E_0|^2 / 8\pi n_{ec}^2 T_e) \quad (52)$$

is the self focusing parameter [27]

Two-dimensional Simulations of Ion-Wave Parametric Instabilities

In most theoretical studies the different ion wave scattering instabilities have been investigated separately, while in experiments, they can occur simultaneously. Therefore, in this section the competition between SRS in different directions, filamentation, and self focusing instabilities will be discussed for the case of a ponderomotive coupling. At first we will describe forward and backward SRS competition and next, we will discuss effects of the laser beam self-focusing on SRS.

The system of equations (49),(50) has been solved in two dimensional cartesian geometry with non periodic outgoing boundary conditions. The incident Gaussian-like laser beam enters at the boundary $x = 0$ and propagates along the x -axis. The length of interaction region L is assumed to be less or comparable to the Rayleigh length of the incident laser beam $k_0 a^2$. This choice corresponds approximately to the hot spots in the laser beams in an underdense plasma corona of ICF targets.

For the adequate comparison with experiments the results of simulations have been represented in terms of the emission field outside the interaction volume. According to the standard Holmholtz-Kirchhoff representation of the radiated field we have calculated the electric field $E(\rho, t)$ in a far field (Fraunhofer) zone at a distance $\rho \gg k_0 a^2$ from the interaction region. The emitted power in a unit angle has been represented as $Q(\theta, t) = c\rho |E(\rho, t)|^2 / 8\pi$, where θ is the emission angle with respect to the propagation axis of the incident beam. Along with the angular distribution Q , the integral quantity of the scattered light – the reflectivity in backward semicircle – have also been calculated:

$$R_b(t) = P_0^{-1} \int_{\pi/2}^{3\pi/2} d\theta Q(\theta, t), \quad (53)$$

where P_0 is the incident beam power. R_b is a convenient characteristics of SRS, describing its most detrimental feature for ICF applications, and allowing for interesting comparisons with one dimensional calculations.

Forward-backward SRS interaction.

The numerical solutions [26] to Eqs. (49),(50) have shown strong dependence of the angular distribution of SRS emission on the ion acoustic mode damping and the geometry of the interaction region. The characteristic temporal dependence of the backscattered reflectivity R_b is shown in Fig.3(a) for the case of a relatively strong damped ion acoustic waves $\Gamma \equiv \Gamma_s(2k_0)/2k_0 c_s = 0.06$, and in Fig. 3 (b) for the case of weak damping $\Gamma = 0.03$. As we will see later, the angular distribution of the scattered radiation can be very complicated and only in the case of a narrow pump beam the main signal is emitted in the backward direction. The dotted curves in Figs. 3 represent results obtained from the one dimensional (1D) version of the code. For a comparison, the dashed curves displayed in Fig 3 have been obtained from 1D runs exploiting driven Korteweg-de Vries (KdV) equation [28] for the ion density perturbations. The one-dimensional results correspond closely to the case of narrow pump ($a < L$), whereas the wide beam runs ($a > L$) exhibit much

Figure 3: Temporal dependence of the reflectivity in rear semicircle $R_b(k_0 c_s t)$ for the case of (a) strong ($\Gamma = 0.06$) and (b) weak ($\Gamma = 0.03$) ion acoustic wave damping. The other parameters of the simulations are the following: $|E_0|^2/8\pi n_{ec} T_e = 0.045$, $n_i/n_{ic} = 0.2$, interaction length $L = 35.5 \lambda_0$; $a = 10 \lambda_0$ for the curves 1, and $a = 50 \lambda_0$ for the curves 2. The dotted lines correspond to 1D results, and dashed curves has been obtained from 1D KdV theory for the same parameters. The insert in panel (a) displays long-term behavior for curve 1 in linear scale.

Figure 4: Angular distribution of the electromagnetic wave power emitted from the plasma slab in the far field zone $\hat{Q} = Q(\theta, t)/(n_{ie} m_i c_s^2 c/k_0)$ for the parameters of curve 1 in the Fig. 3 (a) at time $k_0 c_s t = 15$ (a), 45 (b), 165 (c), and 285 (d)

higher reflectivity. It is interesting to note that 2D simulations display more complicated long time behavior: cf. insert panel in Fig. 3 (a), where a backscattered signal starts to grow up again after a long period of the apparent saturation.

Additional information about physical processes leading to the time evolution of the reflectivity (Fig. 3) are shown in Figs. 4 and 5. They show the angular distribution of the scattered electromagnetic radiation in the far field (Fraunhofer) zone. The transmitted light is emitted in a narrow cone $\Delta\theta \approx \lambda_0/a$ which is less than 10° for the case of narrow pump (curves 1 in Fig. 3), and about 1° for wider pump (curves 2 in Fig. 3). The rest of the electromagnetic radiation in the far field zone is produced due to scattering by the ion density perturbations. The initial level of scattering by the random density fluctuations is about $Q_0 \sim 10^{-5} I_0$ and does not depend on the angle θ . For the case of relatively long interaction regions ($a < L$) (cf. Fig. 4 (a), (b) and 5 (a)), the backward SBS grows first and is emitted in a relatively large cone angle $\delta\theta_b \approx \tan^{-1}(a/L)$. Note that the small-scale ($\sim 1^\circ$) modulation pattern in far field plots is the consequence of finite grid size used in the calculations. The large scale ($\sim 10^\circ$) modulations are the result of scattered light diffraction due to density channel formed by the incident pump beam. A very important parameter, responsible for the evolution of scattering, is the stationary amplification coefficient G_b for backward SBS (eq.(27)). If the ion acoustic damping

Figure 5: Angular distribution of electromagnetic wave power emitted from the plasma slab in the far field zone $Q = Q(\theta, t)/(n_{ie} m_i c_s^2 c/k_0)$ for the parameters of curve 1 in Fig. 3 (b) at time $k_0 c_s t = 30$ (a) and 180 (b)

is small enough and $G_b \geq 10$, SBS exhibits simple temporal behavior (cf. Fig. 5 (b)) backscattered light depletes the pump beam and saturates at a relatively high level. In that case, two dimensional effects are not very important and the total backward reflectivity is closer to one dimensional calculations as can be seen from Fig. 3 (b). The difference between 2D results (curve 1), or dotted curve corresponding to linear 1D simulations, and dashed curve in this figure may be attributed to the ion hydrodynamic nonlinearities. The latter inhibit the growth of the ion waves and are taken into account in the 1D KdV theory.

After the backward SBS saturation, side scattering grows to the levels about two orders of magnitude higher than initial ones, and eventually saturates. The cause of this saturation is demonstrated in Fig. 6 which shows the electric field and density distributions inside the plasma. In Fig. 6 (a), one can see that regions of back scattering and side scattering are separated, and due to the pump depletion, side scattering cannot grow to the level that the linear theory predicts.

If the damping of ion acoustic mode is strong enough and backscattering saturates at a low level without noticeable pump depletion, then the scattering in other directions have a chance to grow. This is illustrated in Fig. 4 (c) where the forward scattering in angles $\theta \sim 40^\circ - 60^\circ$ reaches even larger levels than backscattering. The characteristic angles of forward emission are defined by the Fourier spectrum of density disturbances produced by the pump beam during the process of the density channel formation.

The growth and saturation of forward SBS has important effect on backscattering. One can see in Fig. 4 (d) that additional back reflection shown in insert panel of Fig. 3 (a) has qualitatively different angular shape: its angular divergence is approximately the same as the pump beam's one. The near field distribution can help to understand the cause of this collimated backscatter. The forward and backward scattering processes are spatially separated because of convective nature of these instabilities and due to the propagation of scattered waves in opposite directions. The density disturbances related to backscattering are localized near the entrance boundary (see Fig. 6 (b)). The density fluctuations produced by forward SBS reach their maximum near the forward boundary and can act as an additional seed for backscattering. Due to diffraction of the scattered waves on these density inhomogeneities, the preference in amplification among all possible

Figure 6: Spatial distribution of electric field amplitude $|E|/\sqrt{4\pi n_e m_e c_s^2}$ (a) and the ion density fluctuations δN (b) in the plasma slab for the parameters of curve 1 in Fig. 3 (b) at time $k_0 c_s t = 180$.

Figure 7: Temporal dependence of the reflectivity in the rear hemisphere $R_b(k_0 c_s t)$ for the different ion acoustic wave damping: $\Gamma = 0.03$ (curve 1), 0.1 (2), 0.2 (3), and 0.3 (4). The other parameters of the simulations are the following: $|E_0|^2 / 8\pi n_{ec} T_e = 0.045$, $n_i / n_{ic} = 0.5$, $k_0 a = 20$; interaction length $L = 55.5 \lambda_0$ for curve 1, and $L = 45.5 \lambda_0$ for curves 2, 3, 4.

structures has the backward propagating wave, whose shape is approximately the same as the pump beam shape. This wave has smaller losses and therefore can gain more energy. These arguments, which are similar to that used in the light phase conjugation theory [29, 30], elucidate the reasons for the collimated backscattering generation at the late stage of SBS temporal evolution. In other words, the forward SBS produces the conditions for the effective electromagnetic mode competition, which are absent in an initially homogeneous plasma. Therefore, these simulations demonstrate the possibility and provide the necessary conditions of the phase conjugation effect in laser-plasma interaction.

SBS from intense localized beam structures.

In this section we will discuss simulations when the ponderomotive self-focusing instability is allowed in the interaction region. Figure 7 demonstrates the temporal dependence of the backward reflectivity (53) for several runs with different ion acoustic damping. The self-focusing parameter (52) $\alpha_{sf} = 9$ is the same for all curves. The damping parameter Γ controls backward SBS amplification coefficient and doesn't affect self-focusing. In the case of weak damping (cf. curves 1 and 2 in Fig. 7), backscattering apparently dominates and leads to a high level reflectivity, which is determined by the pump depletion. Self-focusing plays the secondary function in this run. This is why curves 1 and 2 in Fig. 7 look similar to the curves displayed in Fig. 3. Because of the fast growing backscattering, the intense electric fields are localized near the entrance line and the conventional self-focusing has not enough length for its excitation. However, the far field distributions exhibit few angular sidebands in forward and backward directions which can be attributed

Figure 8: Temporal dependence of electromagnetic emission from the plasma slab $\hat{Q} = Q(\theta, t)/(n_{ie}m_i c^2 c/k_0)$ in straight forward ($\theta = 0$), (a) and backward ($\theta = \pi$), (b) directions. The parameters of this run correspond to curve 4 in Fig. 7.

to the filamentation instability of two counter propagating light beams – incident and backscattered [31]. This instability, in contrary to the conventional filamentation, is an absolute one and can operate even for the small beam overlapping region if the beam's intensities are comparable. The time history of the far field distribution shows that this absolute filamentation instability arises after the saturation of backward SBS and develops in a longer time scale.

Runs with higher ion acoustic damping $\Gamma = 0.3$, curve 4 in Fig. 7, show less reflectivity and more complicated temporal behavior. The reflectivity pulsation's are characterized by the much longer period than the backward SBS saturation time and they illustrate the competition between SBS and self-focusing. This competition is even more visible if the temporal history of the narrow cone emission has been traced, as it is shown in Fig 8. The brightness of light emission in straightforward ($\theta = 0$) and backward ($\theta = \pi$) directions oscillate with the period, which is approximately equal to the time the ion acoustic wave needs to pass the incident beam cross-section. This is very clear indication of the interaction between backward SBS and self-focusing instabilities. We should note here that the competition between self-focusing and backward SBS was reported recently in Ref. [32] in the paraxial optics approximation. In their example the backward SBS restricts self-focusing instability and results in the periodic flashes of backscattering. In our simulations we see a different scenario: SBS has no serious effect on self-focusing saturation, but mainly follows the electromagnetic field intensity variations inside the plasma. This situation is illustrated in Fig. 9, which corresponds to the large value of self-focusing parameter $\alpha_{sf} = 16$. At time $k_0 c_s t = 30$, due to self-focusing high intensity electric field has been generated and boosts backscattering reflectivity more than two orders of magnitude with respect to convective saturated level. Later, the focus position moves forward along x -direction and a new focus starts to appear. This dynamic behavior of self-focusing in-

Figure 9: Spatial distribution of the electric field amplitude $|E|/\sqrt{4\pi n_{ic} m_i c_s^2}$ (a) and the ion density fluctuations δN (b) in the plasma slab at time $k_0 c_s t = 30$ for the following parameters: $|E_0|^2/8\pi n_{ec} T_e = 0.08$, $n_i/n_{ic} = 0.5$, $k_0 a = 20$; interaction length $L = 45.5 \lambda_0$.

stability produces strong oscillations in backward SBS reflectivity as we already discussed earlier.

Comments on 2D studies on ion wave parametric instabilities

Two dimensional simulations of laser interaction with ion plasma modes demonstrate a variety of regimes depending on the ion acoustic wave damping, geometry of the interaction region and the laser beam intensity. There are two main dimensionless parameters which control the interaction processes. The first one is the backward SBS gain coefficient G_b and the second one is the self-focusing parameter α_{sf} . In the case of small $\alpha_{sf} \leq 1$, the competition between backward, side, and forward SBS takes place. The backscattering unavoidably dominates for the gain $G_b \geq 10 - 12$. For lower gains, forward SBS grows to a higher level due to the saturation of the backscattered SBS. The characteristic angle of forward scattering is directly related to the smallest spatial scale in the incident beam cross-section. In the long term evolution, signatures of phase conjugation has been observed in the backscattered light signal.

In the case of large values of the parameter $\alpha_{sf} \geq 10$, the main issue is the competition between self-focusing and backward SBS. If the gain coefficient is sufficiently large $G_b \geq 15$, SBS definitely wins, producing strong pump depletion and arresting self-focusing. For smaller values of G_b , backscattering cannot prevent self-focusing, and the latter has a dramatic effect on SBS temporal evolution. This regime is characterized by strong temporal oscillations of the backscattered and transmitted light and production of high amplitude electric field spikes inside the interaction region. The field maxima have a short life-time, they dissolve, and grow again in a different location. This behavior has a lot in common with the "flicker" mode of self-focusing reported in Ref. [33] with the difference that in our case it happens even for a smooth Gaussian incident beam and backscattering plays an important role in the evolution of the instability. The transmitted light in that case is splitted into few beamlets with time-varying intensities and propagation directions.

This study has been restricted here to the case of a smooth incident beam and a homogeneous plasma slab. This has been made with the idea to elucidate the characteristic regimes of the interaction, and provide a simple analytical description of any of the processes. The more realistic models of strong amplitude and/or phase modulated incident electric field which are relevant to the experiments with RPP, ISI and SSD laser beams [21] are now under investigation. But as a first order approximation, the present paper may be regarded as the study of temporal behavior of a one speckle of a multimode laser beam in a plasma. From that point of view, the control parameters G_b and α_{sf} can be used for the description of multimode beams also, if the corresponding parameters will be evaluated to any speckle or "hot spot" of the laser beam.

SCATTERING INSTABILITIES IN SHORT TIME SCALES: EXPERIMENTAL EVIDENCE

In this section we will review results from experiments performed with a laser pulse of 600 ps duration, and in which the temporal resolution of the diagnostics was sufficient to make evident many features occurring in time scales of the order of ten picoseconds.

Experimental Configuration

Preformed plasmas experiments were developed at Ecole Polytechnique to investigate laser-plasma interactions in long and quasi-homogeneous plasmas [34]. These experiments were performed using two beams of the LULI Nd:glass laser facility [35]. The approach

consist of using one laser beam to preform a plasma with known characteristics, and a second laser beam to pump the instabilities. This arrangement permits independent control of plasma conditions and laser parameters. The plasma was formed using one beam at either $1.053\mu\text{m}$ or $0.527\mu\text{m}$, and focused with a combination of spherical and cylindrical lenses. The focal spot was rectangular with approximate dimensions of $1.3\text{mm} \times 150\mu\text{m}$. The laser pulse was Gaussian in time with a full width at half maximum of 600 ps, with no temporal modulation. The incident intensities on target was $1 - 2 \times 10^{13}\text{W}/\text{cm}^2$. The targets were CH foils with thickness between 200nm and 800nm . To avoid edge effects, the targets were approximately $700\mu\text{m}$ wide, smaller than the length of the focal spot. The initial foil thickness and time delay between the two beams determined the plasma parameters at the time of the interaction. Plasmas with $n_e/n_c < 0.2$ and $T_e < 0.5\text{keV}$ were obtained 600 ps after the peak of the pulse [36]. The second beam (either at the fundamental frequency or converted to second harmonic) was focused with an $f/2.7$ lens along the long axis of the plasma. The focal spot diameter was $70\mu\text{m}$ for 80% of the laser energy, with a maximum averaged intensities of $2 \times 10^{15}\text{W}/\text{cm}^2$. Refraction effects from the interaction beam were negligible. The length of the interaction region ($\sim 300\mu\text{m}$) was set by the focal depth of the focusing lens and not by the plasma length.

A diagnostic for the plasma electron density is provided by the SRS light itself. From the linear dispersion relations and the wave-matching conditions for the Raman process, there is a unique correspondence between the plasma density (n_e/n_c) where the Raman light originates, and the scattered wavelength (λ_s). For a low plasma density this is given by

$$\frac{n_e}{n_c} = (1 - \lambda_o/\lambda_s)^2 [1 - 5.9 \times 10^{-3} T_e f(x)] \quad (54)$$

$$f(x) = \frac{(1+x)^2}{(1-x)^{-2}} \left(\frac{1}{x}\right) (3x - x^2 - 1) \quad (55)$$

where $x = \lambda_o/\lambda_s$ and T_e is the electron temperature in keV.

Stimulated Raman Scattering

Stimulated Raman Scattering (SRS) was studied both in backward and in forward directions. The backward Raman emission was collected by the main focusing lens, while forward Raman was collected by a second $f/2.7$ lens in the forward direction. Typical time-resolved spectra for both forward and backward emission, recorded from the same shot, are shown in Fig. 10. The temporal resolution is determined by the combination of the resolution of the streak camera and the temporal dispersion introduced by the diffraction grating spectrometer.

They both show a narrow spectral emission corresponding to a narrow range of plasma densities present in the interaction region. The Raman spectra corresponds to an electron densities $\sim 0.05n_c$ at the beginning of the pulse, and shifts later in time towards lower densities, down to $0.03n_c$ before the end of the pulse. At these low electron densities, Landau damping of the electron plasma waves quenches the SRS instability. The lowest electron density at which backward Raman light is observed is referred as the Landau cut-off and, because of its dependence on electron temperature, it is often used as a temperature diagnostic. For the experiments described here, the cut-off corresponds to an electron temperature of 0.4 keV. The mean electron density deduced from backward and forward Raman spectra were in agreement for different time delay between the beams and for different target thickness.

Both backward and forward Raman spectra exhibit temporal structures which occur in burst lasting $< 20\text{ps}$. Similar temporal structures are observed in the SBS light, which will be discussed in the next section. Although the streak cameras did not have an absolute

Figure 10. Time-resolved spectra of stimulated Raman scattering in backward and forward directions from a preformed homogeneous plasma. The interaction beam is at 526nm and the intensity is 10^{13} W/cm^2 .

time mark as a reference, there was enough evidence from the intensity modulations to indicate correlation between the instabilities. If a correlation exists, it is through coupling of ion sound waves from SRS and electron plasma waves from SRS. Such a coupling is supported by results of experiments described in the following sections. These observations also suggest that forward Raman is linked to backward Raman as it has been observed in other experiments [37], [38] and in numerical simulations [39].

Upshifted and downshifted Raman scattering. Time-resolved spectra of the up- and down-shifted SRS light collected in the backward direction are shown in Fig. 11. The two spectra show a narrow range of emission corresponding to a narrow range of plasma densities in the interaction region. The up and down components are well correlated in time and symmetrically shifted from the fundamental frequency at all times. Using the dispersion relations for the different waves, and the phase-matching conditions (Eqs (1) and (2)), the same range of plasma densities were obtained from the upshifted and down-shifted spectra. They show a similar temporal evolution, determined by the hydrodynamic expansion of the plasma. One interesting point is that both emissions show short bursts in time which are well correlated in time.

The observation of up- and down-scattered SRS emission in the backward direction can only be explained in terms of mode-coupling between the SRS-generated Langmuir wave and an ion wave [40]. Because upscattering from either backward or forward SRS will be emitted in the forward direction, one needs another mechanism to explain the observed emission. It is not possible for light to be reflected back, since these experiments did not have a critical density that could reflect either the pump or any component of SRS light. The upscattered light, with frequency $\omega_{up} = \omega_0 + \omega_L$, can only be explained by Thomson scattering of the pump EM wave (ω_0, k_0), from counter-propagating Langmuir waves with

Figure 11 Time-resolved spectra of the upshifted and downshifted Raman components from a preformed plasma. Corresponding electron densities are plotted on the right-hand scale.

almost the same frequency and opposite wave vector ($k_{up} = k_{SRS} - 2k_0$) than the SRS driven Langmuir wave. Several mechanisms have been proposed for the origin of a counter propagating electron plasma wave. The correlation between the upshifted SRS light and the SBS backscattered observed in two different plasma configurations (preformed plasma and thin target foils) supports the conjecture of mode coupling through the SRS-driven ion sound wave. This coupling could be responsible for the temporal structures that have been observed in both up- and downshifted time-resolved spectra.

Stimulated Brillouin Scattering

Backscattered Stimulated Brillouin Scattering emission was studied using a similar optical arrangement as described in the previous section. The spectra were recorded using a combination of high resolution spectrometer and streak camera, with a wavelength resolution of 0.01 nm and a temporal resolution of 30ps. The plasma was imaged onto the slit of the spectrometer, thus allowing some spatial resolution. This permitted to collect Brillouin light coming only from a small region in the center of the plasma. This is different from other experiments where light coming from a larger volume of plasma is collected on the spectrometer, leading to temporal smoothing of the spectra due to spatial integration. Since the plasma was under-critical at the time of arrival of the interaction beam, the backreflected light was due to SBS, with no contribution from specular reflection.

A time-resolved spectrum is shown in Fig.12. The spectrum is very narrow at all times and red shifted from the laser wavelength by a constant value of around 0.7 nm. These spectral features confirm that SBS is emitted from an homogeneous plasma and with very little flow velocity along the laser axis. The spectra is completely different from those observed in thin exploded foil experiments, where red shifts increasing as a function of

Figure 12 Time-resolved spectrum of stimulated Brillouin backscattering from a preformed homogeneous plasma. The interaction beam is at 526nm and the intensity is 10^{15} W/cm^2

time are attributed to Doppler effect from the recoiling foil. The combination of a long interaction region and low flow velocity along the laser axis explains the high level of SRS reflectivity observed in this experiment, as compared to exploding foil targets, which have an intrinsic larger velocity gradients and lower reflectivities. Under similar irradiation conditions, time-integrated reflectivity's were of the order of 10^{-2} in preformed plasmas, and around 10^{-4} in thin exploded foils.

A feature we want to underline in this review is the fast evolution of observed reflectivity. SRS emission appears as a collection of very short bursts, shorter than 30ps, indicating the rapid growth, saturation and quenching of the instability. This temporal modulation do not show any periodicity and vary randomly from shot to shot. Instantaneous reflectivity's are then much higher than time-integrated reflectivity's, demonstrating the difficulty in comparing experimental data and theoretical modeling. These modulations are not coming from the incident laser pulse, and are evidence of non linear effects on the ion waves excited by the instability. One possible hypothesis for these structures is the coupling of ion sound waves due to Brillouin and electron plasma waves due to Raman [41]. Stimulated Raman emission recorded during this same experiment showed similar temporal behavior.

Filamentation

In laser produced plasmas, filamentation is a difficult instability to diagnose. All experimental evidence of filamentation have been inferred from features observed indirectly. In the experiment described here, evidence of filamentation was found from signatures observed in the transmitted laser light, in Raman emission, and in second harmonic emission. The observed temporal structures indicate the presence of density channels in the

Figure 13. Space and time resolved pictures of the transmitted light at the fundamental laser frequency: (a) without plasma, (b) with the preformed plasma. The laser intensity is $1.5 \times 10^{15} \text{ W/cm}^2$ for the two pictures.

plasma that can be interpreted as filaments.

Transmitted light. Two direct consequences of filamentation are intensity modulation of the laser beam as it propagates through the plasma and electron density depression produced by the expulsion of the electrons from the high intensity regions. The first feature provided a diagnostic for filamentation, by time resolving the transmitted laser light. The transmitted light was collected by an $f/2.8$ lens and imaged through a narrow band filter onto the slit of a streak camera. These one dimension images of a cross section of the plasma yielded time resolved measurements of the spatial profile of the light as it traverses the preformed plasma. The spatial resolution was $3\mu\text{m}$.

In the presence of plasma, strong modifications of the spatial intensity distribution of the beam were observed, compared to the initial intensity pattern [12]. The transmitted light evolve during the laser pulse, changing from temporal modulations to spatial modulations. Fig 13 shows time resolved pictures of the transmitted light with the slit of the streak camera placed across an image of the focal plane. Fig 13(a) was recorded without plasma, to characterize the incident laser beam itself. The laser irradiance was $1.5 \times 10^{15} \text{ W/cm}^2$ for both pictures. The observed irregular spatial non uniformity's are due to the far field pattern of the laser beam, and there was no temporal structure expected nor observed from the laser pulse itself. Fig 13(b) was recorded with the preformed plasma, showing a different behavior in time and space. Two different regimes can be identified and will be described in the following paragraph.

During the first half part of the pulse (200-300ps), one can observe short bursts of emission which last less than 20 ps. These bursts are randomly distributed along the slit of the streak camera. This can be interpreted as produced by refraction of the laser beam outside the acceptance angle of the collecting lens, in a repetitive manner. The

Figure 14: Backward time resolved Raman spectra from the preformed plasma showing Raman light emitted from the background density and Raman growing in very narrow density regions. The slope of these Raman bursts is smaller at the end of the pulse than at the beginning of the pulse. The laser wavelength is 527nm and the nominal laser intensity is $1.5 \times 10^{15} \text{ W/cm}^2$.

sudden refraction of the beam could be explained by a fast change of the plasma density, as expected from the nonlinear evolution of filaments inside the plasma.

The latter part of the pulse shows spatial structures which are relatively stable in space and time, indicating that the laser beam was divided into small beamlets, possible due to small plasma density channels. The size of these channels are typically of the order of 5 to 10 μm , corresponding to regions of lower plasma density, as one would encounter in filamentation. Some of these features have been observed in numerical simulations of the nonlinear evolution of filamentation, although the two regimes observed early and late in the laser pulse are not clearly understood.

Raman spectra. An example of time resolved Raman spectra is shown in Fig.14. The background Raman emission displays the characteristics which have been discussed above. Superimposed to this emission one can observe structures with temporal evolution which interpretation requires a different mechanism. They exhibit a temporal slope which is different from the mean slope of the spectrum. They indicate that Raman light is emitted from very narrow density channels with an electron density that is decreasing in time. This can be interpreted as if the Raman instability is growing in filaments induced by hot spots due to the laser intensity distribution. The Raman emission follows the electron density evolution inside the filament, as indicated by the shift towards shorter wavelength in time. The slope of each of these narrow structures is larger than the mean slope of the spectrum, showing a faster decrease of the density in these regions. This temporal evolution is different at the beginning and at the end of the laser pulse. The slopes indicate faster density evolution in the beginning of the laser pulse, consistently with the rapidly varying modulations observed in the transmitted images. Different theories [43] predict modifications of Raman growth rates in filaments. These theories have been used to explain Raman emission with laser intensities below threshold, as well as to explain structures observed in the Raman spectra. In this experiment the laser intensity is above the Raman threshold, so there is no need to invoke filamentation to explain the Raman-background emission. However, the very narrow structures observed in the spectra are

Figure 15: Time-resolved photograph of the transmitted light at second harmonic of the fundamental laser frequency

consistent with the density drop expected from filamentation. The final electron density is around 90% of the initial electron density, which corresponds to a density perturbation produced by filamentation of $dn/n \sim 10\%$, in agreement with the theoretical size of the density fluctuation given by [44]

$$dn/n = 1 - \exp[-0.375(v_o/v_{th})^2] \quad (56)$$

where v_o is the electron quiver velocity and v_{th} is the electron thermal velocity. The time needed to modify the electron density varies from a few picoseconds at the beginning of the pulse to ~ 100 ps at the end of the pulse.

Second harmonic emission. Second harmonic emission was collected in the forward direction, using a $1.053\mu m$ interaction laser beam. A two dimensional cross section of the interaction region was imaged with a combination of two lenses into film with a total magnification of 35. A one dimensional strip across the interaction region was imaged into the entrance of a streak camera, thus providing a time resolved record. A time-resolved photograph of the transmitted light at the second harmonic is shown in Fig.15.

The width of the streak camera slit corresponded to about $150\mu m$ in the plasma. The emission is observed to have strong temporal and spatial modulation. The observed short temporal bursts have a duration of the order of 20 ps, with a spacing between bursts of a few tens of picoseconds. The spatial structures have a dimension between 6 and $10\mu m$. The laser energy converted into second harmonic emission, and collected into the f/2.5 lens, was $3 - 6 \times 10^{-6}$.

Generation of second harmonic emission in the forward direction in underdense plasma requires the presence of a radial modification of the plasma density in order to introduce a strong electron density gradient perpendicular to the beam direction. The filamentation of the laser beam propagating in the underdense plasma produces transverse density gradients, which in the presence of the strong laser electric field, give rise to localized electron currents at $2\omega_o$. The structures observed in time-integrated images show emissions which are well defined in space and which can correspond to filaments.

The dimensions of the observed structures have transverse wavenumbers which maximize the linear spatial growth rate of filamentation [46]. For the experimental conditions, $\lambda_0 = 1.053 \mu\text{m}$, $I = 3 \times 10^{15} \text{W/cm}^2$, $T_e = 0.5 \text{keV}$ and $n_e/n_c = 0.3$, the wavelengths that maximize the spatial growth rate for both ponderomotive and thermal filamentation are $\lambda_{c,p} = 2.5 \mu\text{m}$ and $\lambda_{c,T} = 2.3 \mu\text{m}$ respectively.

The temporal modulations of the transmitted light at the second harmonic are consistent with the formation of filaments, as it has been observed in numerical simulations. In the supersonic regime, the maximum growth rate γ_{max} of the filamentation instability is given by $\gamma_{max} = \sqrt{2}\omega_{pe}v_{os}/c$. For the experimental conditions described earlier, the formation time of the filament, $\tau = 1/\gamma_{max}$, is about 2ps. These temporal structures are consistent with theoretical models for the nonlinear evolution of filaments inside the plasma.

STIMULATED BRILLOUIN SCATTERING WITH 10 PS LASER PULSES

In this section we will describe a study of SBS using a laser pump with a duration of 8-10 ps. This was the first experimental attempt to characterize SBS with a short pulse laser, in order to have a small number of e-folding of growth. The short duration of the interaction also permitted to neglect effects due to hydrodynamics of the plasma as well as filamentation of the laser beam, thus permitting a more meaningful comparison with analytical and numerical interpretations. In quite different conditions, the SRS instability has been studied with a ps laser pulse, both experimentally [49], and theoretically [50].

The experiment was conducted in a preformed plasma, using a long (3 ns) laser pulse to ablate a thin foil target. The interaction pulse (8-10 ps) arrived after a 10 ns time delay to eliminate any coupling with SRS ion waves produced by the first laser pulse. A preformed plasma has the advantage of providing a longer density scale length than what would be obtained with the short pulse alone, as well as providing plasma conditions which are independent of the short pulse intensity. The time integrated reflectivity of the SBS scattered light was recorded with an S-1 photomultiplier, while a detector array provided a time integrated spectra. The photomultiplier have sufficient sensitivity to measure low levels of SBS, as well as the EM noise level of the plasma, with a temporal resolution sufficient to separate the SBS emission of the short pulse from that of the long pulse.

The short pulse was generated from a 100 ps mode-locked pulse, and then chirped and dispersed in a 2 km long optical fiber. At the output of the optical fiber the pulse had a spectral width of approximately 2.5 nm and a duration of 300 ps. This was then amplified in a YAG regenerative amplifier to a pulse energy of a few mJ. Due to the narrow gain bandwidth of YAG (0.2-0.3 nm), only the central part of the pulse was amplified, giving an output pulse with a duration between 8 and 10 ps, and a near transform limited spectral width. This novel technique permitted to obtain short pulses (8-10 ps) at TW power levels, without having to use pulse compression schemes with diffraction gratings, as is normally used in standard chirped pulse amplification (CPA) systems. The pulse was then amplified using a combination of rod and disc glass amplifiers with a maximum aperture of 10 cm. The maximum energy of the short pulse was 10 J. The pulse duration was measured using a single-pulse autocorrelator. A separate autocorrelator was used to monitor the level of prepulse, which was less than 10^{-4} and sufficient for this experiment. The long (3ns) pulse was generated in a separate single longitudinal-mode oscillator, and amplified in a separate amplifier chain to an energy of approximately 10 J. Both pulses were brought together with opposite polarization so that they could be combined using a polarizer, and then focused on target with a single f/10 lens.

Plasma parameters desirable for the experiment were obtained by irradiating a thin plastic 1.5 mm thick foil, with the 3 ns laser pulse at an intensity of 10^{12}W/cm^2 , 10 ns

Figure 16: Stimulated Brillouin scattering reflectivity of a 10 ps laser pulse from a preformed plasma, as a function of the laser intensity

before the arrival of the short pulse. This arrangement produce a plasma with a long electron density scale-length, maximum electron density of approximately $10\%n_{cr}$, and high electron temperature. The 10 ns delay provided enough separation between the two pulses to avoid any seeding of SRS in the short pulse due to SRS produced by the long pulse. The plasma temperature at the time of interaction was determined by the intensity of the short pulse. The increase in plasma temperature was calculated assuming inverse bremsstrahlung absorption, and compared to fluid code simulations. For a short pulse intensity of approximately 10^{15} W/cm^2 , the temperature at the time of the interaction was approximately 1 keV.

The reflectivity of the 10 ps pulse is shown in Fig.16 as a function of the laser intensity. While at low laser intensities the reflectivity has a well defined non zero slope, at high intensities there is clear evidence of saturation. The reflectivity measurements are time integrated, and one expect that the instantaneous reflectivity will be higher at the end of the short pulse. The temporal resolution was sufficient to separate the plasma emission at the time of interaction from any emission or reflection occurring when the plasma forming laser pulse was present. Figure 17 shows the temporally integrated spectrum of the backscattered light. The SRS contribution from the long pulse was less than 10% of the component from the short pulse. For low laser intensities ($I < 10^{14} \text{ W/cm}^2$) the backscattered light consists of two well defined components: one centered close to the incident laser wavelength, and a second component shifted by approximately 0.4 nm and with a spectral width of approximately 1 nm. The spectral width of the unshifted component is approximately 0.2 nm, which corresponds to the instrumental resolution. For higher laser intensities, the red shifted component has a similar shift, but is wider and becomes the predominant part of the spectrum. At intensities of about 10^{15} W/cm^2 a second component is observed which is superimposed on the main component and has a wider spectral width. The observed spectral widths, close to the value of the ion acoustic frequency, indicate that for all laser intensities the interaction is weakly coupled. In a

Figure 17: Time integrated SRS spectra, for three laser intensities; a) $I = 1 \times 10^{13} \text{ W/cm}^2$; b) $I = 1 \times 10^{15} \text{ W/cm}^2$. The vertical axis are in arbitrary units. The high frequency modulation is an artifact of the detector.

strongly coupled regimen, the shifts of the spectrum are a function of the laser intensity.

The dependence of the reflectivity on the laser intensity did not agree with results from 2D modeling [6] [5]. At low laser intensity, the model predicted reflectivity's two order of magnitude higher than observed experimentally. To explain this discrepancy the authors invoked the coupling between the Raman and Brillouin instabilities, a process that was accounted for in the theory.

One may conjecture that although SRS develops from a lower level of noise, it has a significantly larger growth rate; the plasma waves can then couple to the ion acoustic waves, give rise to the parametric decay instability which leads to cascading and/or collapse. These nonlinear processes may in turn generate ion acoustic wave turbulence with ion acoustic waves characterized by a broad spectrum. All these processes can take place on a time scale of a few picoseconds, which is shorter than the time necessary for the SRS gain factor to be larger than unity. Thus the coupling of plasma waves to ion acoustic waves could represent an efficient mechanism for giving rise to an anomalous level of ion acoustic wave noise from which SRS can then develop. It would consequently be of interest for future short pulse experiments to investigate the SRS/SBS correlation. In particular, future experiments performed in plasmas with a lower density in order to suppress SRS would make possible to examine whether the SRS reflectivity becomes the Thomson scattering reflectivity TS, as theoretically predicted. The interaction between these two instabilities have been observed in a number of experiments under different plasma and laser conditions [20].

The utilization of short laser pulses with ps duration offers the possibility of studying the growth and evolution of parametric instabilities more detail than what has been done until now. Ideally one would like to vary the pulse duration so as to control the number of e-folding in the growth. Varying the pulse duration and intensity allows the separate control study of issues related to the growth and noise levels and from nonlinear issues leading to saturation. Using pulses with duration much shorter than 1 ps (e.g. 100 fs) would allow the study of the growth of SRS as well, an instability with faster growth than SBS. Shorter pulses will also permit the study of SRS at much higher laser intensities and still avoid saturation. Further improvements in coming experiments could include the

use of beam smoothing techniques, such as random phase plates, to modify the intensity distribution in the interaction volume, and to use 2ω and 3ω light as a pump, to improve the diagnostic capabilities.

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References

- [1] W. L. Kruer, *The Physics of Laser Plasma Interactions*, Addison-Wesley Publishing Company, Inc., 1988.
- [2] H. A. Baldis, E. M. Campbell, and W. L. Kruer, *Laser Plasma Interactions, Handbook of Plasma Physics*, Vol 3, edited by A. Rubenchick and S. Witkowski (North-Holland, Amsterdam, 1992).
- [3] D. W. Forslund, J. M. Kindel, and E. L. Lindman, *Phys. Fluids* **18**, 1002 (1975).
- [4] A. G. Sitenko, "Fluctuations and nonlinear wave interactions in plasma", (Pergamon, New York, 1972), Chap. 4, p. 126.
- [5] H. A. Baldis, D. M. Villeneuve, B. LaFontaine, G. D. Enright, C. Labaune, S. Baton, Ph. Mounaix, D. Pesme, M. Casanova and W. Rozmus, "Stimulated Brillouin scattering in picosecond time scales : experiments and modeling", submitted to *Phys. Fluids B*.
- [6] Ph. Mounaix, D. Pesme, M. Casanova and W. Rozmus, "Space and time behavior of parametric instabilities for a finite pump wave duration in a bounded plasma", submitted to *Phys. Fluids B*.
- [7] E. A. Williams, R. R. McGowan, ICF Quarterly Report **1**, 147 (1991), Lawrence Livermore National Laboratory, Livermore CA, UCLR-LR-105821-91-4.
- [8] D. F. DuBois (Private communication).
- [9] H. A. Rose, D. F. DuBois, B. Bezzerides, *Phys. Rev. Lett.* **58**, 2547 (1987).
- [10] W. Rozmus, R. P. Sharma, J. C. Samson, W. Tighe, *Phys. Fluids* **30**, 2181 (1987).
- [11] G. Bonnaud, D. Pesme, R. Pellat, *Phys. Fluids B* **2**, 1618 (1990).
- [12] T. Kolber, W. Rozmus, V. T. Tikhonchuk, *Phys. Fluids B* **5**, 138 (1993).
- [13] B. Bezzerides, D. F. DuBois, H. A. Rose, "Saturation of stimulated Raman scattering by the excitation of strong Langmuir turbulence", submitted to *Phys. Rev. Lett.*
- [14] T. Kolber, W. Rozmus, V. Tikhonchuk, Ph. Mounaix, and D. Pesme, "Laser light scattering from fluctuations enhanced by stimulated Raman scattering", *Phys. Rev. Lett.* (in press, 1993).

- [15] W. L. Kruer, Comments Plasma Phys. Controll. Fusion 9, 63 (1985).
- [16] V. T. Tikhonchuk, Sov. Phys. Uspekhi 34, 903 (1991).
- [17] S. H. Batha, K. S. Bradley, H. A. Baldis, R. P. Drake, K. Estabrook, L. W. Johnston, and D. Montgomery, Phys. Rev. Lett. 70, 802 (1993).
- [18] C. L. Shepard, J. A. Taryn, G. E. Busch, R. R. Johnson, R. J. Schroeder, and R. L. Berger, Phys. Fluids 29, 583 (1986).
- [19] T. Afshar-rad, S. E. Coe, O. Willi, and M. Desselberger, Phys. Fluids B 4, 1301 (1992).
- [20] H. A. Baldis, D. M. Villeneuve, C. Labaune, D. Pesme, W. Rozmus, W. L. Kruer, and P. E. Young, Phys. Fluids B 3, 2341 (1991).
- [21] Y. Kato, K. Mima, N. Miyanaga, S. Aringa, Y. Kitagawa, M. Nakatsuka, and C. Yamanaka, Phys. Rev. Lett. 53, 1057 (1984); R. H. Lehmberg and S. P. Obenshain, Opt. Commun. 46, 27 (1983), X. Deng, X. Liang, Z. Chen, W. Yu, and R. Ma, Appl. Optics 25, 377 (1986).
- [22] P. Frycz, W. Rozmus, J. Samson, V. T. Tikhonchuk, and R. Rankin, Comments Plasma Phys. Controll. Fusion 15, 19 (1992).
- [23] A. A. Andreev, N. E. Andreev, A. N. Sutyagin, and V. T. Tikhonchuk, Sov. J. Plasma Phys. 15, 546 (1989).
- [24] R.W. Short and E.M. Eppelrain, Pys. Rev. Lett. 68, 3307 (1992)
- [25] H.A. Rose and D.F. Dubois, Phys. Fluids B 4 1394 (1992)
- [26] M. R. Amin, C. E. Capjack, P. Frycz, W. Rozmus, and V. T. Tikhonchuk, "Two dimensional studies of stimulated Brillouin scattering, filamentation and self-focussing instabilities of laser light in plasmas", Phys. Fluids , submitted for publication (1993).
- [27] A. J. Schmitt, Phys. Fluids 31, 3079 (1988).
- [28] W. Rozmus, M. Casanova, D. Pesme, A. Heron, and J.-C. Adam, Phys. Fluids B 4, 576 (1992).
- [29] R. H. Lehmberg, Phys. Rev. Lett. 41, 863 (1978).
- [30] B. Ya. Zel'dovich, N. F. Pilipetsky, and V. V. Shkunov, *Principles of Phase Conjugation*, Springer Verlag, Berlin (1985) [Springer Series in Optical Sciences, Vol 42].
- [31] C. J. McKinstrie and M. V. Goldman, J. Opt. Soc. Am. 9, 1778 (1992).
- [32] N. E. Andreev and L. M. Gorbunov, JETP Lett. 56, 144 (1992)
- [33] R. D. Jones, W. C. Mead, S. W. Coggeshall, C. H. Aldrich, J. L. Norton, G. D. Pollak, and J. M. Wallace, Phys. Fluids 31, 1249 (1988).
- [34] F. Amiranoff, C. Labaune, G. Matthieussent, B. Mabilie, P. Mora, D. Pesme, M. Haguenaue, LULI Annual report, 50 (1986)
- [35] D. Bruneau, A.M. Tournade and E. Fabre, Appl. Opt. 24, 3740 (1985).

- [36] C. Labaune, F. Amiranoff, G. Matthieussent, D. Pesme, C. Rousseaux and H.A. Baldis Rev. Phys. Appl. 23, 1505 (1988).
- [37] D.M. Villeneuve, and H.A. Baldis, Phys. Fluids 31, 1790 (1988)
- [38] C. Rousseaux, F. Amiranoff, C. Labaune, and G. Matthieussent, Phys. Fluids (to be publish, 1993)
- [39] G. Bonnaud, D. Pesme, and R. Pellat, Phys. Fluids B 2, 1618 (1990); G. Bonnaud, D. Pesme, and R. Pellat, ibid 4, 423 (1992).
- [40] C. Labaune, H.A. Baldis, S.D. Baton, D. Pesme, T. Jalinaud, Phys. Rev. Lett. 69, 285 (1992)
- [41] J.A. Heikkinen and S.A. Karttunen, Phys. Fluids 29, 1291 (1986)
- [42] C. Labaune, S.D. Baton, T. Jalinaud, H.A. Baldis, D. Pesme, Phys. Fluids. B4, 2224 (1992)
- [43] C.S. Liu and V.K. Tripathi, Phys. Fluids 29, 4188 (1986). R.W. Short, W. Seka and R. Bahr, Phys. Fluids 30, 3245 (1987). H.C. Barr, T.J. Boyd and G.A. Coutts, Phys. Rev. Lett. 56, 2256 (1986)
- [44] P.K. Kaw, G. Schmidt, and T. Wilcox, Phys. Fluids 16, 1522 (1973).
- [45] E.M. Epperlein, Phys.Rev. Lett. 65,2145 (1990).
- [46] C.S. Liu, M.N. Rosenbluth, and R.B. White, Phys. Fluids 19, 1211 (1974).
- [47] C.J. Walsh, D.M. Villeneuve, and H.A. Baldis, Phys. Rev. Lett. 53, 1445 (1984).
- [48] H.A. Baldis, R.P. Drake, W.L. Kruer, K.G. Estabrook, E.A. Williams, and T.W. Johnston, Phys. Rev. Lett. 62, 2829 (1989).
- [49] C.B. Darrow, C. Cloverdale, M.D. Perry, W.B. Mori, C. Clayton, K. Marsh, and C. Joshi, Phys. Rev. Lett. 69, 442 (1992).
- [50] S.C. Wilks, W.L. Kruer, A.B. Langdon, P. Amendt, D.C. Eder, and C.J. Keane, Proc. SPIE "Short-Pulse High-Intensity Lasers and Applications" ,1413, 131 (1991).

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