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NONLINEAR BEHAVIOR OF LIGHT-DRIVEN PLASMA INSTABILITIES

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NONLINEAR BEHAVIOR OF LIGHT-DRIVEN PLASMA INSTABILITIES*

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A common feature of light-plasma coupling is the excitation of plasma instabilities by the intense laser light. The instabilities give rise to light absorption, reflection and the production of energetic particles. The nonlinear behavior of many of these instabilities has been examined in computer simulations and theory. We present a nonlinear theory for the turbulent heating due to parametric instabilities near the critical density (where the laser light frequency equals the electron plasma frequency). The linear instability is saturated by mode coupling of unstable electron plasma waves into Landau damped ones, and the electrons are heated by velocity space diffusion in the plasma waves. The theoretical predictions are in good agreement with computer simulation results. We also present computer simulations of laser light back-scattering in the underdense plasma outside the critical density due to the Raman and Brillouin instabilities. The light reflection is shown to decrease rapidly as the ratio of the light pressure to the plasma pressure is made small. Several techniques to suppress these reflective processes are demonstrated. Finally we briefly discuss recent simulations of the $2\omega_{pe}$ decay instability and self-focusing.

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Introduction

The coupling of intense laser light with plasmas is a subject of great interest in laser-fusion studies. The excitation of plasma instabilities is a common feature of the light plasma interaction. The oscillating two stream and ion acoustic decay instabilities occurring near the critical density and the $2\omega_{pe}$ decay instability occurring near one-fourth the critical density primarily enhance the light absorption. The Raman and Brillouin instabilities occurring in the underdense plasma beyond the critical density enhance the light reflection. An additional instability, self-focusing, can lead to a filamenting of the light as it propagates into the plasma. We have studied important features of the nonlinear evolution of all these plasma instabilities. Let us discuss a number of specific examples, which illustrate the instability effects.

Instability-Generated Light Absorption Near the Critical Density

Anomalous absorption plays an important role in the heating of hot, nearly collisionless plasmas by intense laser light. Anomalous (or turbulent) heating is simply the transfer of energy from the laser light into the particles via plasma waves. The richness of this process should not be understated. Energy can flow from the laser light into plasma waves either by means of instabilities or by direct coupling mechanisms. For example, near the critical density instabilities are excited because the light couples electron plasma waves and ion waves⁽¹⁾. Direct driving of plasma waves occurs when a component of the light field acts across a density gradient. This takes place near the critical density when the light is incident at an angle to the density gradient⁽²⁾. Energy flows from the plasma waves into the particles by collisional damping and by Landau damping (and particle trapping).

In order to understand the basic features of anomalous absorption near the critical density, we have studied in some detail a simple but very basic model problem. We consider the heating of a uniform collisionless plasma by an intense spatially homogeneous electric field oscillating near the electron plasma frequency. For laser pellet applications this pump field is due to incident laser light and is in general partly electromagnetic and partly electrostatic. The detailed make-up of this effective field oscillating near ω_{pe} depends on the angle of incidence of the light, the plasma density gradient, and the diameter of the laser light beam. But in any case the fundamental problem is how such an intense,

very long wavelength field (compared to the electron Debye length) heats the collisionless plasma. It is important to understand both the efficiency of the heating and the nature of the turbulently heated velocity distributions. The velocity distributions are particularly important, since very energetic particles degrade pellet performance both due to preheat⁽³⁾ and to decoupling⁽⁴⁾.

The plasma heating by an imposed pump field oscillating near ω_{pe} has been extensively investigated in computer simulations⁽⁵⁾. The basic feature is that the long wavelength pump field (assumed $k_0 = 0$ in the simple model) drives up a spectrum of short wavelength electron plasma waves via the oscillating two stream and electron-ion decay instabilities. These plasma waves then heat the particles by either trapping them or by diffusing them in velocity space.

Figures 1a, 2 and 3 summarize typical results from a computer simulation in the regime in which the energy density in the pump field is much less than the plasma kinetic energy density; i.e. $n^2 \equiv E_0^2/4\pi n_0 e \ll 1$. Figure 1a shows the evolution of the wave energy in electron plasma waves driven unstable by the imposed pump field oscillating at the electron plasma frequency. The unstable plasma waves exponentiate in agreement with linear theory and then saturate at a level a few times the pump field energy. As shown in Figure 2 a broad spectrum of plasma waves has been generated in the nonlinear state. Concomitant with saturation, a turbulent heating of the plasma onsets. This turbulent heating is quite efficient, corresponding to an effective collision frequency (ν^*) of $\sim .3$ the ion plasma frequency. For this collisionless plasma the heating primarily consists of the generation of suprathermal tails on the electron distribution function. A typical

turbulently heated velocity distribution function, $f(v)$, is shown in Figure 3, where we plot $\ln f(v)$ vs velocity. Note the central, relatively unheated Maxwellian distribution plus the long suprathermal tails.

We have developed a nonlinear theory⁽⁶⁾ which allows us to understand the instability saturation and the turbulent heating in some detail. The simulations show that the plasma waves become comparable in amplitude to the pump field driving them. Hence in the theory we allow the unstable plasma waves to themselves act like pump fields and so drive other plasma waves unstable. In this way the initially unstable electron plasma waves (which are only very weakly Landau damped) can transfer their energy into higher wavenumber plasma waves which are Landau damped. These damped waves serve as a sink for the energy flow from the imposed pump field and thereby saturate the instability.

The mode coupling can be described by a set of nonlinear equations which incorporate the interaction between the high frequency electric fields (E_k) and the low frequency ion density fluctuations (n_k).

$$\ddot{E}_k + \omega_{ek}^2 E_k + \nu_{ek} \dot{E}_k = -\omega_{pe}^2 \sum_{k' \neq 0} n_{k'} E_{k-k'}, \quad (1)$$

$$\ddot{n}_k + \omega_{ik}^2 n_k + \nu_{ik} \dot{n}_k = -\frac{k^2}{8\pi n_0 M} \sum_{k'} E_{k'} E_{k-k'}. \quad (2)$$

Here ω_{ek} (ω_{ik}) is the Bohm-Gross (ion acoustic) frequency and the ν 's are the damping rates for these waves (assumed to be Landau damping). v_{0s} is the electron oscillation velocity in the pump field and $\frac{m}{M}$ is the electron-ion mass ratio. These equations are a generalization of Nishikawa's

analysis⁽¹⁾ to the regime in which the fluctuating fields (E_k) are comparable in amplitude to the pump.

The left-hand side of Equations 1 and 2 describe the (damped) oscillation of each fluctuation at its characteristic frequency. The right-hand side shows the nonlinear coupling which has a simple physical origin. The high frequency fields oscillate electrons across the density fluctuations produced by the low frequency ion waves to drive other high frequency fields. In turn, the wave pressure due to the high frequency fields (ponderomotive force, has a low frequency component which drives ion waves.

To complete the theory we need to follow the plasma heating by evolving the velocity distribution of the particles. In the weak pump field regime (no electron trapping), a simple description for the wave-particle interaction can be obtained. A broad spectrum of randomly phased electron plasma waves is driven up in the nonlinear state. These waves impart random kicks to the nearly resonant particles (those with velocities nearly equal to the phase velocity of the wave). Under the action of this stochastic force, the electrons random walk in velocity space. This basic nature of the wave-particle interaction has been confirmed by following the motion of test particles in the computer simulations⁽⁷⁾. As is well known, we can then describe the evolution of the velocity distribution function (f) by the quasi-linear diffusion equation.

$$\frac{\partial f}{\partial t} = \frac{\partial}{\partial v} D \frac{\partial f}{\partial v} \quad , \quad (3)$$

$$D(v) = \frac{\pi e^2}{m} \sum_k E_k^2 \delta(\omega - kv). \quad (4)$$

The diffusion coefficient, D , is the step size of the random walk and is a function of the plasma wave fields.

Equations (1-4) constitute our theoretical description of the turbulent heating. The coupled mode equations determine the wave fields, and the distribution function is given in terms of these fields by the diffusion equation. These coupled nonlinear equations are straight-forward to solve using a computer. For example, we solve equations 1 and 2 by considering a large number of Fourier modes and then computing the evolution of the many coupled ordinary differential equations. Typically 50 modes are used, with spot checks using more modes (a finer subdivision of k space). Equation 3 is solved by a standard finite differencing of the diffusion equation.

The results of this theory are in good agreement with the simulation results. The predictions are shown in Figures 1b and 3, for the same parameters as those used in the typical simulation we have shown. The plasma wave energy (Figure 1b) first exponentiates in agreement with linear theory and then saturates at a level of approximately twice the pump field energy, in agreement with the simulation. As shown in Figure 3, the turbulently heated distribution function computed from the diffusion equation also agrees well with the results of the simulation. These comparisons confirm that our theoretical description captures the dominant physical effects: mode coupling of the unstable plasma waves into Landau damped waves plus diffusion of the electrons in velocity space.

We have discussed the model problem in some detail in order to elucidate the basic physics of anomalous absorption. Important features of the turbulent heating, such as suprathermal tail formation, can then be understood in their simplest context. The production of non-Maxwellian

velocity distributions has been directly measured in recent laboratory experiments on plasma heating by microwaves⁽⁸⁾. In addition, very energetic particles have been observed in laser plasma experiments by monitoring the hard X-rays which are produced⁽⁹⁾. Our simulations and theory show that non-Maxwellian distributions are to be expected, since they are a characteristic feature of turbulent heating of collisionless plasmas by high frequency fields.

Of course, it is clear that for pellet applications our model is quite ideal and should only be expected to give a coarse upper limit for the energetic particle production. In simulations and theory we have found many effects which reduce the number of energetic particles, such as the limited region over which the turbulence extends in a plasma with a density gradient⁽¹⁰⁾ and the presence of short-wavelength ion waves⁽¹¹⁾. The improved models predict fewer energetic particles, but the velocity distributions are still non-Maxwellian. It should be noted that an appreciable level of electron-ion collisions also leads to fewer energetic particles. There is a regime above instability threshold in which the unstable plasma waves are stabilized by scattering into collisionally damped waves⁽¹²⁾--a complementary limit to the collisionless case we here considered. Particle simulations, using a collision operator, are now underway to determine how far above instability threshold this "collisional" regime extends⁽¹³⁾. Detailed quantitative estimates of the turbulently heated distribution functions await a better understanding of other aspects of the light-plasma interaction, such as self-focusing and light scattering. In the interim, it is clearly prudent to be cautious of pellet designs which are extremely sensitive to energetic particles.

Instability-Generated Laser Light Reflection Outside the Critical Density

It is well-known that instabilities can also lead to a reflection of laser light. This instability-generated reflection occurs in the underdense plasma outside the critical density and is due to the decay of the laser light into a scattered light wave plus either an electron plasma wave⁽¹⁴⁾ (the Raman instability) or an ion wave⁽¹⁵⁾ (the Brillouin instability). We have studied these instabilities in computer simulations in order to assess their importance and to find techniques to suppress their effects. Computer simulations were first applied to the reflective instabilities by Forslund, Kindel and Lindman⁽¹⁶⁾.

Since the linear theory of these instabilities has been treated in many papers, we will directly discuss our simulation results⁽¹⁷⁾. In general one should allow for the possibility of both back-scatter and side-scatter⁽¹⁸⁾ of the laser light. To begin with the most tractable case for particle simulations, let us confine ourselves to examining back-scatter. The techniques we present for reducing the instability effects will also apply to side-scatter. We then use a $1\frac{1}{2}$ dimensional electromagnetic particle simulation code. This code works directly with the electric and magnetic fields and is fully relativistic. The plasma density profile rises linearly from zero to .7 of the critical density (10^{21} cm^{-3} for a Nd laser) in a distance of 10 free space laser wavelengths. This model is similar to that used in valuable studies recently reported by Forslund, et al.⁽¹⁶⁾ with the exception that this code is relativistic. The critical density is not included in this simple model in order to isolate the effects of the reflective instabilities from the usual para-

metric ones occurring near the critical density. Many free-space wavelengths of vacuum are included on either side of the plasma to allow for plasma blow-off. The laser light is introduced from the left, and electromagnetic waves are allowed to propagate freely into space from either simulation boundary. The simulation particles are charge-weighted to provide an optimum representation of the low and high density regions of the plasma.

It is instructive to separate and study each instability individually. Since the Raman instability depends only on the electrons, it can be isolated by using immobile ions. Let us briefly describe the basic features of these simulations. In agreement with linear theory, the Raman instability maximizes where the plasma density $n = \frac{1}{4} n_{cr}$ (the critical density), where it drives an electron plasma wave with a wavelength nearly equal to the laser wavelength plus a long wavelength backscattered transverse wave. For intense light the instability is simply saturated by electron trapping, whereas for less intense light other nonlinearities play a role (such as the resonant decay of plasma waves when ion motion is included). Two features of the nonlinear evolution are of particular interest: the light reflection and the generation of very energetic electrons.

As noted in Ref. (16) the net back-scattering due to the Raman instability is not severe. Figure 4a shows the percentage of light energy reflected as a function of η^2 ($\eta^2 = E_0^2 / 4\pi n_{cr} \theta_e$, where E_0 is the laser electric field, n_{cr} is the critical density, and θ_e is the electron temperature). For this data the laser light intensity was fixed, and the

electron temperature was varied. The observed net reflection is at most 9% (for $\eta^2 \gg 1$) and decreases rapidly as the electron temperature is increased. For high electron temperatures ($T_e \gtrsim 5$ keV), the driven plasma waves are appreciably damped and so are not driven to large amplitudes. For low electron temperatures, strong electron trapping occurs, heating the electrons until the driven plasma waves again become heavily damped. In these latter cases, light energy is strongly reflected only for a short time while the plasma is heating.

An important feature of the Raman instability is the energetic electrons which are produced. Since the phase velocity of the plasma wave driven near $1/4 n_{cr}$ is $\sim 1/2 c$, electrons acquire velocities nearly equal to c when they are trapped by the wave. Hence we use a fully relativistic code. Since these electrons can preheat the core of a pellet⁽³⁾, it is important to explore means by which they can be suppressed. Finite amplitude, short wavelength ion fluctuations efficiently couple high phase velocity plasma oscillations into slow, damped ones which produce less energetic electrons⁽¹⁹⁾. Such ion waves can be produced in the underdense plasma surrounding a pellet due to, for example, counter-streaming between different ion species in the blow-off or steepening of the ion waves produced by other instabilities. In order to demonstrate their beneficial effect, we have simply introduced a short wavelength ion wave as an initial condition into our simulations of the Raman instability. Figure 4b shows the percentage of the kinetic energy in electrons with velocities greater than $.5 c$ after the instability saturation as a function of ion wave amplitude. In this example the laser light intensity is 3×10^{16} Watts/cm². There is substantial reduction for ion wave amplitudes $\delta n/n \approx 20\%$. For

less intense light, smaller amplitude ion waves are equally effective.

One final point concerning the Raman instability should be noted. As Rosenbluth points out⁽²⁰⁾, the instability threshold depends on density gradient, because the resonance conditions for the laser decay can only be satisfied locally in the plasma. This is particularly important for back-scatter. For a 10μ density gradient and a Nd laser, we find an abrupt shut-off of the Raman instability at an intensity of $\sim 10^{16}$ Watts/cm². This is in reasonable agreement with the theoretical estimates^(16,20) for back-scatter.

This suppression of the Raman back-scatter by density gradient allows us to isolate and study the Brillouin instability. We now use an electron-ion mass ratio of 300. This mass ratio well separates the time scales for ion and electron motion and also allows us to run many ion periods in the simulations. We have varied the mass ratio to spot-check our results.

Again let us briefly describe the basic features of these simulations. Ion waves, with half the wavelength of the laser light, exponentiate as expected from linear theory. Correlated with the buildup of these ion waves, light reflection occurs. We measure this reflection by monitoring the Poynting vector on the left side of the system. This reflection exhibits relaxation oscillations which persist a number of cycles, as is typical in coupled mode interactions. We average the reflected power over the oscillations, since an average over these time scales is inherently carried out in experimental measurements.

The time-averaged reflection due to the Brillouin instability versus n^2 is shown in Figure 5a. For $n^2 \gg 1$ the reflection is quite large ($\sim 50\%$

even for our 10μ plasma). However, it rapidly decreases as n^2 becomes small. This is particularly important since laser fusion schemes operate in the $n^2 \ll 1$ regime. The data was obtained by varying the laser light intensity and/or the thermal energy of the plasma, indicating that n^2 is a reasonable parameter. This decrease of the reflection with n^2 is consistent, for example, with results found by Forslund, et al.⁽¹⁶⁾ that less reflection occurs when the electrons become heated due to either the Raman or other instabilities.

A simple physical argument for the observations can be given. The energy density of ion waves is $(\frac{\delta n}{n_0})^2 n_0 k T_e$, where $\delta n/n_0$ is the root-mean-square ion wave amplitude. The electron temperature enters since the electrons provide the pressure in a ion wave. Given intensity laser light will then produce smaller amplitude ion waves when the electron thermal energy is larger. Hence the reflection decreases, since it depends on the level to which the ion waves are driven.

In the small n^2 regime we have doubled the size of the plasma (to 20μ), and the reflection approximately doubles. When we increase the electron-ion temperature ratio to 10, the reflection increases several-fold. This again is as expected physically. Fixed intensity laser light produces larger ion waves if their damping is less. In this regard, it should be noted that the electron and ion temperatures remain closer together ($T_e/T_i \sim 3-5$) when one uses a short wavelength laser (heating at higher plasma density) than when one uses a long-wavelength laser such as CO_2 .

A nonlinear theory is necessary to quantitatively predict the reflection. An analytic solution for the back-scatter of laser light in a spatially homogeneous medium has been carried out, assuming the only nonlinearity is pump depletion⁽²¹⁾ and that the instability is in its convective regime. This theory illustrates some characteristic features of the spatial back-scatter in the limit in which the phonon is heavily damped. For the present applications, however, one must allow a density gradient, plasma expansion, and nonlinear damping.

The instability generated reflection can be substantially reduced by using laser light with a finite bandwidth. Figure 5b shows the reflected power versus the light bandwidth for a sample case in which the net laser light intensity is 10^{16} Watts/cm² (Nd) and the electron temperature is 1 keV. The light bandwidth is simply modeled by distributing the laser energy in many individual fields, each with a slightly different frequency and random phases. Similar results have been obtained by using laser light with random phase modulation. This decrease in the reflection agrees with simple estimates. Basically only that fraction of the power within a given frequency width (proportional to the linear growth rate of the instability) acts coherently. This is also found in linear instability calculations for finite bandwidth laser light⁽²²⁾. For weaker intensities even smaller bandwidths are effective at suppressing the reflection. In another interesting limit we find that the reflection is also reduced when one uses two intense monochromatic laser fields with different frequencies, due to destructive interference between the unstable plasma waves.

The simulations of back-scatter are by no means a definitive treatment of light scattering in plasmas, but they do indicate several important points. First, nonlinear effects such as particle trapping and plasma wave decay can play a central role in determining how large the reflection becomes. It is important to understand how to model these nonlinearities in terms of, for example, effective damping rates which are a function of wave amplitude. Secondly, it is important to examine the reflection in the limit in which the energy density of the laser light is much less than the plasma energy density ($n^2 \ll 1$). The net reflection is a strong function of this characteristic parameter, as we have shown for the Brillouin instability. Finally, we have demonstrated several physical effects (such as laser light bandwidth) which reduce deleterious effects of the reflective instabilities.

Other Instability Investigations: the $2 \omega_{pe}$ Decay and Self-Focusing

Several other instabilities are also under study. The first of these is the $2 \omega_{pe}$ decay instability⁽²³⁾ which occurs near one-fourth the critical density. This instability is due to the decay of the light into two electron plasma waves and represents an additional absorption mechanism. The instability is inherently a two-dimensional effect, since transverse currents cannot couple to longitudinal ones in only one dimension.

Let us briefly summarize our studies of this instability carried out by Langdon, Lasinski and Krue⁽²⁴⁾. First we find that the previous linear

theory has to be modified to include the $\mathbf{v} \times \mathbf{B}$ force on the electrons. This extension is necessary to describe those unstable plasma waves which have wave numbers comparable to the wave number of the laser light. Simulations using a 2-D electromagnetic, relativistic particle code show good agreement with the extended linear theory.

Nonlinearly the instability is observed to saturate by electron trapping in the decay unstable electron plasma waves^(24,25). Figure 6 shows a typical momentum space plot from a simulation of intense laser light propagating through a spatially homogeneous plasma. The light frequency is $2.04 \omega_{pe}$. Note that very high energy particles have been produced by the electron trapping ($\gamma v/c \sim 1.3$, where $\gamma = 1/\sqrt{1 - v^2/c^2}$). The electrons in the forward direction have been further accelerated by the electron plasma wave due to the Raman instability, which occurs on the same time scale as the $2 \omega_{pe}$ decay. On the longer time scale associated with ion motion, we observe both the Brillouin instability and the parametric instabilities due to decay of the electron plasma waves which have been generated. These simulations are, of course, ideal (rather intense light in a spatially homogeneous plasma), but they are a first step to understanding the complex interaction between the many instabilities.

Lastly the self-focusing⁽²⁶⁾ of laser light has also been examined in computer simulations. Light concentrates in regions of low plasma density, thereby expelling more plasma. The result is the formation of a narrow filament (or filaments) of intense laser light. Recent simulations by Langdon and Lasinski⁽²⁷⁾ with both a 2-D and a 2½-D electromagnetic code have shown self-focusing of laser light into a narrow channel several

wavelengths wide. They find that there is then a strong heating of electrons near the sides of the channel, provided the light polarization is such that the electric field of the light is directed across the density gradient and the plasma density outside the channel is above one-fourth the critical density. This strong heating is attributed to the instability of electron plasma waves which occurs when electrons are oscillated across a density gradient at nearly twice the local electron plasma frequency. (This instability is similar to the $2\omega_{pe}$ decay, but now the sharp density gradient at the sides of the channel is the source of the plasma currents which drive the waves.) These simulations, which isolate self-focusing, show that the gradients in plasma density which are generated can introduce new mechanisms for plasma heating. We are now studying the competition between plasma heating, self-focusing and light scattering (which in general occurs on a faster time scale).

Conclusion

We see that light-driven plasma instabilities produce a rich variety of nonlinear effects: absorption, reflection, energetic particles, and filamentation. Our studies of the nonlinear evolution of the many instabilities help us both to identify their basic effects and to estimate their importance for laser fusion applications. The quantitative understanding of light-plasma interactions is at an early stage of development, particularly as related to the complex interplay between the many different instabilities. An intensive, world-wide study of these instabilities promises to steadily advance our knowledge of them, with significant

benefits both to laser fusion applications and to other areas of plasma physics. It is especially encouraging that laser plasma experiments are becoming increasingly sophisticated and detailed. A close interaction between such experiments and calculations will greatly enhance our ability to sort out and describe the most important interactions.

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FIGURE CAPTIONS

Figure 1. Electric field energy as a function of time a) from a computer simulation and b) from the coupled mode equations. $\eta^2 = .13$, $\omega_0 = \omega_{pe}$, $m/M = .01$, and $\theta_E/\theta_i = 3.0$.

Figure 2. Saturated turbulent field spectrum from the computer simulation at time $\omega_{pet} = 703$.

Figure 3. Electron velocity distribution from the computer simulation at time $\omega_{pet} = 326$. The temperature at this time is $2.5 T_{e0}$. The dashed line is the distribution calculated from the diffusion equation at the corresponding time.

Figure 4. Simulation results for the Raman instability with immobile ions. a) The reflected light energy (normalized to the incident energy) versus η^2 . b) The percent of energy in electrons with velocities greater than .5C versus ion wave amplitude. The laser light intensity is $3 \times 10^{16} \text{ w/cm}^2$ (Nd), and the imposed ion wave has a wave number $k\lambda_D = .2$, where λ_D is the Debye length at $1/4$ the critical density.

Figure 5. Simulation results for the Brillouin instability. The electron-ion mass ratio is 300, and the temperature ratio is 5. a) The reflected light energy versus η^2 . b) The reflected light energy versus the laser bandwidth. ω_{pi} is the ion plasma frequency at the critical density.

Figure 6. A typical momentum space plot ($\frac{y}{c}$ vs $\frac{y}{c}$) from a two-dimensional simulation of the $2 \omega_{pe}$ decay instability.

FIGURE 1

FIGURE 3

FIGURE 4

FIGURE 5

FIGURE 6