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A Model for the Instability Absorption in
a Laser Heated Plasma

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A MODEL FOR THE INSTABILITY ABSORPTION IN A LASER HEATED PLASMA

by

R. C. Mjolsness and H. M. Ruppel

ABSTRACT

A model is presented for the evaluation of the laser flux absorbed by a plasma instability mechanism. The model is suitable for inclusion in hydrodynamic codes. The absorbed flux is proportional to a two-parameter function which is given in phenomenological form. While the theory has been made consistent with existing theoretical information on the behavior of instability absorption, a considerable extrapolation of existing theory is necessary to evaluate the absorbed flux uniquely.

I. INTRODUCTION

There are a number of preliminary experimental reports of absorption in laser heated plasmas apparently in excess of the absorption which can be attributed to inverse bremsstrahlung. Partially in order to explain this phenomenon theorists have considered a number of mechanisms which could contribute additional absorption (usually referred to as anomalous absorption) to the plasma. These include electrostatic instabilities,¹⁻² ion-ion collisions modified by the laser field³ and resonance absorption mechanisms.⁴⁻⁵ In this work the absorption due to the electrostatic instability is evaluated in a simple way to provide a useful estimate of the effects of one of the more significant mechanisms for anomalous absorption.

With simplifying assumptions the absorption is expressed in terms of one two-parameter function, f , presented here both in a table and as a phenomenological expression which should be both suitable and convenient for use in hydrodynamic codes. The calculations¹⁻² which the present work seeks to model contain a saturation of the absorption when a certain temperature is reached. Although growth continues in the so-called saturation region, it does so at a considerably lower

rate. In order to take this into account it is simplest to change discontinuously at the saturation temperature an overall constant which appears in the expression for the absorbed flux. With this two-regime approach to the problem it becomes possible to apply the model in all regions of interest. The expressions given here for the instability absorption are limited primarily by our incomplete knowledge of the instability process except under highly idealized physical conditions. In particular, calculations do not include the effects of density gradients, and nonlinear information is sparse even for homogeneous systems. Moreover, our evaluation of the spatial distribution of the laser electric field will be a non-self-consistent one. On the other hand, the approximation involved in replacing our table of values for the absorbed laser flux by a two parameter phenomenological expression is known to be a minor one.

II. DISCUSSION OF THE MODEL

A basic assumption of the present model is that the density of the plasma target increases linearly as a function of distance into the plasma. Although it is unlikely that this will be realized

in practice, there are reasons why this approximation may work more satisfactorily than appears at first. Consider that the target be broken up into three regions: The subcritical region in which the plasma frequency is less than the laser frequency, the critical region in which $\omega_p = \omega$ and the supercritical region in which the electric field is rapidly attenuated regardless whether the electric energy is primarily reflected or absorbed. The reason that the assumption of a linear density profile is sensible is that in most applications the instability absorption will be confined to the one zone of a hydrodynamic code which contains the critical density, and in any one zone the density profile is necessarily taken to be linear. The theory will thus be strictly applicable only to cases where instability absorption occurs in one zone.

To evaluate the absorption due to the electrostatic instability it is proposed that a term denoted symbolically by "(instability absorption)" be added to the equation for electron energy transfer and similarly be subtracted from the equation for laser energy transport. Then the electron energy equation would have the symbolic form

$$\left(\frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla\right) \left(\frac{3}{2} n k T_e\right) = (\text{e-e heat conduction}) + (\text{bremsstrahlung emission}) + (\text{e-i heat transfer}) + (\text{inverse bremsstrahlung absorption}) + (\text{instability absorption}) \quad (1)$$

This theory is concerned with presenting a simple expression for "(instability absorption)" as an energy flux to be added to the electrons in the critical zone, the last zone to which the laser energy penetrates. This expression is then quite suitable for use in modifying existing hydrodynamic codes.

In order to evaluate the "(instability absorption)" flux under conditions of physical interest in pellets or foils--which may contain strong density gradients and may have energy removed from the unstable region by electron-electron heat conduction and other processes--it is necessary to make a considerable extrapolation from highly idealized calculations.^{1,2,6} It is assumed here that at each point of the unstable region energy is removed from the laser electric

energy density at the rate at which the most unstable electrostatic mode of a spatially homogeneous system at the local density increases its energy, as long as the electron temperature is small (unsaturated growth). When the electron temperature exceeds a critical value (saturated growth) this rate of decrease is altered by a definite ratio to agree with one particular nonlinear plasma simulation. This is essentially all the nonlinear information on the evolution of the instability which is presently available. The absorbed flux is then calculated by integrating over space this time rate of change of local energy density. However, to affect the integration simply, a highly idealized evaluation of the local laser electric field is employed.

Consider the division of the growth of the electron temperature into two regions: an early time during which the growth is exponential with rate 2γ (unsaturated growth) and a later time which is characterized by a nearly linear rate (saturated growth). The perturbation calculation of Freidberg and Marder predicts a $\gamma(\kappa, n, T)$ where κ is the wave number of the mode of the instability. Here we use n and T without subscripts to refer to the local electron density and temperature, as there should be no danger of confusion. When γ is obtained for given n as a function of T , the most rapidly growing mode can be fit by the function

$$\gamma_{\max}(T) = \frac{\gamma_1}{(1 + b_0(kT)^{1/2})} \quad (2)$$

Actually, the fit will yield 0.53 instead of $1/2$ in the exponent in the denominator, but this difference is not significant for the evaluation of the flux and it will be ignored. This relation can be justified only for small amplitudes of the instability such that perturbation calculations apply. However, an examination of the nonlinear results of Nielson and Freidberg⁶ indicates that in this spatially homogeneous problem the electron temperature can be characterized by the above growth rate even in the saturated regime, if the constant γ_1 is changed at the temperature which marks the transition from one regime to the other. Although it is not clear why this result should be valid, it will be assumed here as a general pro-

property of the nonlinear behavior of the instability. Some assumption concerning the nonlinear behavior must be made; the existing information is simply insufficient to permit a complete treatment of nonlinear effects. The present assumption is at least simple and in agreement with existing information. It has a consequence that the form of Eq. (3) is used throughout with a discontinuity of approximately a factor of thirty in γ_1 . This discontinuity could easily be replaced by a smooth interpolating function if desired, but nothing would be contributed to the physical description of the process by doing this.

Let us return now to a discussion of the plasma density. As mentioned in the beginning of this section, the assumption of a linear density is particularly suitable for the zoned calculations current in hydrodynamic codes. However, it does lead to some error: Because of the rapid change of plasma density near the critical density, the density gradient will often be evaluated incorrectly in actual calculation. Moreover, the unstable region will frequently have sufficient physical extent in the supercritical region that the linear approximation introduces errors.

The local instability absorption rate depends on the value of the local electric field intensity and an approximation is introduced into the calculation of this intensity: The field is assumed to propagate from $z = -\infty$ in response to the real plasma dielectric constant $\epsilon = [1 - (\omega_p^2/\omega^2)]$, of the linear plasma profile only, and to be unaffected by the inverse bremsstrahlung absorption, the instability absorption and the actual shape of the plasma profile. Here ω is the laser frequency in radians and ω_p is the electron plasma frequency. A simple approximation allows this technique to yield essentially the correct electric energy to the boundary of the unstable region, but the instability absorption and the inverse bremsstrahlung absorption in the unstable region remain unaccounted for. The value of the maximum instability growth rate as a function of density is also idealized as a piecewise linear function. These approximations may well introduce an uncertainty of, say, a factor of two into the evaluation of the absorbed flux for a given set of plasma parameters. However, the absorbed flux

depends rather sensitively on parameters, varying by many orders of magnitude in the parameter range of possible interest. The main trends of the absorption are certainly correctly reproduced by this model and such things as the boundary in parameter space defining, let us say, 90% absorption are probably also fairly well determined.

Finally it is required that the plasma density change little over one vacuum wavelength of the laser radiation; otherwise the theory is undeveloped.

III. EVALUATION OF THE ABSORBED FLUX

Consider a slab of plasma with zero density at $x = -L$ and critical density at $z = 0$. The density, of course, is assumed to vary linearly in the z -direction. A steady laser beam with vacuum wave length λ is directed at the slab from $z = -\infty$. Using the fit to $\gamma_{\max}$ discussed in section II we can model the instability absorption with the phenomenological expression

$$\frac{d}{dt} \left(\frac{3}{2} nkT \right) = \frac{2\gamma_0 \left(\frac{3}{2} nkT \right) g(z)}{(1 + b_0 (kT)^{3/2} |E(z)|^{-1})} \quad (3)$$

In this equation k is the Boltzmann constant, and the product $\gamma_0 g(z)$ represents the γ_1 of Eq. (2); the variation of the growth rate with density is given by $g(z)$ while the normalization of the growth rate is given by the present γ_0 . Definitions of these two quantities are given below. On the left hand side of Eq. (3) n and T should be interpreted as the zone averaged values for use in hydrodynamic codes. On the right hand side T should similarly be taken to be the zone averaged quantity, but because of the spatial locality of the instability process n is probably best replaced by n_c , the critical density at which $\omega = \omega_p$. Strictly, this density should be included in the spatial integration for the absorbed laser flux, but this minor correction we do not make. The parameter b_0 is defined by

$$b_0^2 = \frac{\pi}{2} \frac{me^3}{0.0974e^{-2}\lambda^2 I} \quad (4)$$

where λ is the vacuum wavelength of the incident laser flux and

$$I = (\ell_N/L)^{1/2} I_N \quad (4a)$$

where I_N is the flux incident on the unstable zone and $z = -\ell_N (\ell_N < L)$ gives the boundary of the unstable zone. The quantities I_N and ℓ_N will be available from other portions of the hydrodynamic calculation and I_N will differ from the incident flux I_0 principally due to inverse bremsstrahlung absorption in the underdense region of the plasma.

The function $g(z)$ is obtained phenomenologically in the following way. The piecewise linear function γ , which approximates the curves of Freidberg and Marder for γ_1 has the form

$$\begin{aligned} \gamma_1 &= 0, \quad n \leq n_0 \\ &= \gamma_0 \left(\frac{n - n_0}{n_c - n_0} \right), \quad n_0 \leq n \leq n_c \\ &= \gamma_0, \quad n \geq n_c \end{aligned} \quad (5)$$

where n is the electron density, n_c is the critical density and n_0 is the density at the onset of the unstable region. Note that n_c is slightly mass dependent; but this dependence, which is approximately the tenth power of the ion mass, will be ignored. Now since we are assuming a linear density profile such that

$$n = n_c (1 + z/L),$$

we obtain

$$\gamma = \gamma_0 \left[1 + \frac{z}{L} \frac{1}{1 - n_0/n_c} \right], \quad n \leq n_c \quad (6)$$

Hence we can write $g(z)$ in Eq. (3) as

$$\begin{aligned} g(z) &= 0, \quad z < -L(1 - n_0/n_c) \\ &= 1 + \frac{z}{L} (1 - n_0/n_c)^{-1}, \quad -L(1 - n_0/n_c) \leq z \leq 0 \\ &= 1, \quad z > 0 \end{aligned} \quad (7)$$

For an ion of mass equal to the mass of the hydrogen ion the beginning of the unstable regime ($n = n_0$) occurs at $n_0 = 0.8464 n_c$. This is the value of n_0 which will be used for the numerical results presented here. As mentioned above, the value of n_0 approaches n_c as the ion mass increases; however this effect will be ignored.

In the exponentially growing region γ_0 can be scaled with mass using the phenomenological approximation to the Freidberg and Marder¹ theory

$$\gamma_0 = 0.035 (M(\text{hydrogen})/M(\text{ion}))^{1/3} \omega, \quad (8)$$

where ω is the laser frequency. In the saturated

regime no comparable result exists.

It was noted in section II that it is necessary to change the value of γ_0 when the temperature reaches saturation. Again this is a phenomenological result in that such a change continues to make Eq. (2) valid but with altered γ_0 . The saturation temperature can be obtained from the curves of Freidberg and Nielson⁶ as

$$d_{\text{kev}} = 1.1 \times 10^{-17} I \lambda_\mu^2 \quad (9)$$

Recall that I is related by Eq. (4a) to the laser flux (in joules/sec/cm²) incident on the unstable zone and that λ_μ is the vacuum wave length in microns. In order to know more precisely what to do in the saturated regime, it would be necessary to extend the work of Freidberg and Nielson to include the effects of inhomogeneities, collision processes, density variations and ions of various masses. In lieu of such information, what can be done at present is to examine the ratio of γ_0 's in the saturated and unsaturated regions for one particular density and mass ratio and use this ratio to change γ_0 in all cases. From the curves of Freidberg and Nielson one obtains the ratio

$$\gamma_0(\text{saturated})/\gamma_0(\text{unsaturated}) = 1/30 \quad (10)$$

The physical electric field in the unstable region is the real part of the expression $E = E(z)e^{-i\omega t}$ and $E(z)$ is calculated in the present approximation from the expression

$$\left[\frac{d^2}{dz^2} + k^2 \epsilon \right] E(z) = 0, \quad (11)$$

or

$$\left(\frac{d^2}{dz^2} - k^2 \frac{z}{L} \right) E(z) = 0 \quad (11a)$$

The field is, of course, subject to the boundary conditions of being an exponentially damped disturbance for $z > 0$ and of being the sum of two waves --incoming and outgoing--of equal amplitude at the zone boundary, $z = -\ell_N$ of the unstable zone. In addition, $E(z)$ can be and is taken to be a real function, since these boundary conditions define a standing wave. Inspection of Jahnke, Ende and Lösch, p. 156 discloses that $E(z)$ must have the form

$$E(z) = C \sqrt{z} Z_{1/3} \left(\frac{2}{3} i \beta \left(\frac{z}{L} \right)^{3/2} \right), \quad (12)$$

where $\beta = kL$ and $Z_{1/3}(w)$ is a linear combi-

nation of Bessel function of argument w and order $(1/3)$. The requirement of exponential dumping for $z > 0$ forces $Z_{1/3}$ to be a Hankel function of the first kind, i. e., $Z_{1/3}(w) = H_{1/3}^{(1)}(w)$. The application of the boundary condition at $z = -\ell_N$ to determine C is more awkward to do exactly, although it certainly presents no difficulty in principle. Here we will carry out this evaluation asymptotically, obtaining results which are accurate when $w_N \equiv (2/3)\beta(\ell_N/L)^{3/2} \gg 1$ and not too bad when $w_N \simeq 1$. Then the electric field which has incoming and outgoing waves of unit amplitude at $z = -\ell_N$ turns out to be

$$E_N(z) = e^{i(\frac{2\pi}{3})(\frac{\ell_N}{L})^{1/4}[\frac{\pi}{3}\beta(\frac{z}{L})]^{1/2}} H_{1/3}^{(1)}(\frac{2}{3}\beta(\frac{z}{L})^{3/2}). \quad (13)$$

With this electric field the I of Eq. (4) should be taken to be I_N , the flux incident on the boundary of the unstable zone. However, this choice is a poor one, since the integral defining the absorbed flux then depends on the zoning of the problem and must be redone for each zoning. Thus we have preferred to work with

$$E(z) = e^{i(\frac{2\pi}{3})(\frac{\ell_N}{L})^{1/4}[\frac{\pi}{3}\beta(\frac{z}{L})]^{1/2}} H_{1/3}^{(1)}(\frac{2}{3}\beta(\frac{z}{L})^{3/2}) \quad (14)$$

and to take I of Eq. (4) to be $(\ell_N/L)^{1/2} I_N$. We have tested the approximation involved in Eqs. (13) and (14) and find it adequate whenever $L > \lambda$ and $\ell_N > 0.1L$.

It is clear that with the use of Eq. (14) the integral defining the absorbed flux can be expressed as a function of the two parameters β and $b \equiv b_0(kT)^{1/2}$. These two parameters depend on the properties of the plasma in the unstable zone and on the properties of the laser flux incident on this zone.

We define an instability absorbed flux, F , to be added to the electron internal energy and subtracted from the laser flux by the expression

$$F \equiv (\text{instability absorbed flux}) \\ = f(b, \beta) [2\gamma_0 L (\frac{3}{2} n_c kT)] \quad (15)$$

where

$$f(b, \beta) = f_-(b, \beta) + f_+(b, \beta) \quad (16)$$

The functions f_+ and f_- are obtained by integrating the time rate of change of the energy, Eq. (3), along the z direction into the slab of

plasma. This then yields the energy flowing into the slab, per unit area, per unit time, i. e. a flux. The subscripts "minus" and "plus" refer to the underdense and overdense regions, respectively. That is:

$$f_-(b, \beta) = \frac{1}{L} \int_{-L/\alpha}^0 \frac{dz [1 + \alpha z/L]}{(1 + b_0(kT)^{1/2} |E(z)|^{-1})} \quad (17)$$

(where, for convenience, we have defined $\alpha \equiv [1 - (n_0/n_c)]^{-1}$) and

$$f_+(b, \beta) = \frac{1}{L} \int_0^\infty \frac{dz}{(1 + b_0(kT)^{1/2} |E(z)|^{-1})} \quad (18)$$

The flux, F , may be evaluated from these equations whenever the calculation yields $F \leq I$. However, since this theory does not include a self-consistent treatment of the electric field, it may result that $F > I$. When this occurs it is necessary to take $F = I$.

IV. SUMMARY OF RESULTS

In order to facilitate inclusion of the instability absorbed flux into hydrodynamic codes a phenomenological fit was made to f as a function of the quantities b and β . The parameter β has been defined previously as

$$\beta = 2\pi L/\lambda \quad (19)$$

and b as

$$b = b_0(kT)^{1/2}$$

The integrals which yield f_+ and f_- were first obtained numerically for 220 pairs of values of b and β spanning the parameter range of primary physical interest:

$$10^{-3} \leq b \leq 10^4$$

and

$$2\pi \leq \beta \leq 2000\pi \quad (20)$$

Then a study was made of the asymptotic behavior of the integrals for large b and large β . Using this analysis as a guide to start off the search a gradient scheme was employed to minimize the sum of the squares of the deviations of the phenomenological fit from the numerical result. The resulting expressions for f_+ and f_- are

$$f_+ = b^{-1} (2\pi/3\beta)^{1/2} [1 - e^{-0.673\beta}] \\ \times [1 - \{0.492 \ln(\beta)/(1 + 0.3\beta)\}]$$

where $\delta = b (2\beta/3)^{-1/6}$, and (21)

$$f_- = 0.248 \frac{\{1 + 0.536[1 - e^{-b}][1 - (2\pi/\beta)^{1/2}]\}}{\{[1 + 0.536(1 - e^{-b})][3.25 + b]\}}.$$

As before $f(b, \beta) = f_+ + f_-$, and F is obtained from $f(b, \beta)$ with the use of Eq. (15). To summarize the prescription for obtaining the absorbed flux due to the electrostatic instability: One picks values of b and β appropriate to the system; then F is calculated from the phenomenological forms, Eq. (21), for f_+ and f_- and from Eqs. (15) and (16), using Eq. (8) for γ_0 at low temperatures. Finally when the temperature reaches saturation, as defined in Eq. (9), γ_0 is lowered by the factor of Eq. (10), and this new instability absorption is used as long as the temperature is above saturation temperature.

An attempt was made to evaluate the integrals analytically by using the asymptotic form of the Hankel functions. For certain regions of b and β this produced reasonable results for f_- . However, f_+ , which often is the major contributor to f , is less well evaluated; in fact even f_- is poorly represented by this WKB result for large b and β .

Table I is included to document the reliability of the phenomenological fit and to indicate the size of f for the range of b and β . When the gradient search was made 220 points were used rather than only the representative 28 points displayed here. It should be emphasized that the table contains f and not the flux itself which in the text is represented by F .

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TABLE I

VALUES OF THE ABSORPTION FUNCTIONS f_- , f_+ AND $f = f_- + f_+$ AT REPRESENTATIVE POINTS IN THE b , β PARAMETER SPACE. THE LAST COLUMN IS THE PHENOMENOLOGICAL FIT TO f GIVEN IN EQ. (21).

b	$\beta(2\pi)$	Exact f			Fit To f
		f_+	f_-	f	f
.001	1	.17732E+01	.76760E-01	.14500E+01	.14577E+01
.001	10	.10682E+00	.76760E-01	.38359E+00	.38638E+00
.001	100	.68404E-01	.76718E-01	.14512E+00	.14578E+00
.001	1000	.15237E-01	.76712E-01	.91949E-01	.91835E-01
.010	1	.10441E+01	.76400E-01	.11205E+01	.11066E+01
.010	10	.23744E+00	.76511E-01	.31395E+00	.31076E+00
.010	100	.53753E-01	.76178E-01	.12993E+00	.12935E+00
.010	1000	.12140E-01	.76093E-01	.88233E-01	.88131E-01
.100	1	.66188E+00	.72983E-01	.73486E+00	.73257E+00
.100	10	.15722E+00	.74191E-01	.23141E+00	.22937E+00
.100	100	.36921E-01	.72326E-01	.10925E+00	.11030E+00
.100	1000	.86060E-02	.72040E-01	.80646E-01	.82424E-01
1.000	1	.25716E+00	.50454E-01	.30761E+00	.30369E+00
1.000	10	.68588E-01	.57749E-01	.12634E+00	.12362E+00
1.000	100	.17810E-01	.53827E-01	.71637E-01	.75086E-01
1.000	1000	.45276E-02	.52924E-01	.57452E-01	.62496E-01
10.000	1	.43064E-01	.12373E-01	.55437E-01	.52286E-01
10.000	10	.13104E-01	.18379E-01	.31483E-01	.28789E-01
10.000	100	.39242E-12	.17497E-01	.21421E-01	.21738E-01
10.000	1000	.11578E-02	.17139E-01	.18297E-01	.19613E-01
100.000	1	.46688E-02	.14483E-02	.61171E-02	.68336E-02
100.000	10	.14699E-02	.23622E-02	.38321E-02	.37534E-02
100.000	100	.46178E-03	.23226E-02	.27844E-02	.29100E-02
100.000	1000	.14465E-03	.22976E-02	.24422E-02	.25239E-02
1000.000	1	.47092E-03	.14735E-03	.61827E-03	.73029E-03
1000.000	10	.14881E-03	.24319E-03	.39200E-03	.39900E-03
1000.000	100	.46953E-04	.24037E-03	.28732E-03	.29479E-03
1000.000	1000	.14859E-04	.23829E-03	.25315E-03	.26207E-03
10000.000	1	.47134E-04	.14760E-04	.61894E-04	.73767E-04
10000.000	10	.14905E-04	.24391E-04	.39296E-04	.40265E-04
10000.000	100	.47041E-05	.24121E-04	.28825E-04	.29679E-04
10000.000	1000	.14900E-05	.23919E-04	.25409E-04	.26335E-04