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A PARAMETER STUDY OF THE BEHAVIOR OF LINEARLY EXPANDING LASER PLASMAS

By

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A Parameter Study of the Behavior of Linearly Expanding Laser Plasmas

Introduction

In order to perform an experimental investigation of the interaction of laser light with the plasmas characteristic of laser driven implosions, a well behaved plasma target must first be prepared with a pulse of low intensity radiation. The behavior of the plasma may then be observed by irradiation with second, high intensity laser pulse. The plasma phenomena resulting from this interaction of high intensity radiation with the plasma, are generally dependent upon the electron density gradient (scale height), and the electron temperature of the target. A detailed analysis of the plasma phenomena requires that these characteristic parameters be known.

The purpose of this report is to develop an approximate analytic description of the expansion of disk shaped laser targets in terms of these parameters. Estimates are derived from computer simulations and hydrodynamic theory for the prediction of time scales, peak densities, scale heights and temperatures for the laser plasmas. The most useful results are summarized in the conclusion. Furthermore, a list of crude thresholds for the most important known plasma instabilities is compiled which may be utilized in conjunction with the expressions for the characteristic parameters to design preliminary experiments for the investigation of plasma phenomena.

A. Three Dimensional Expansion of a Hot Plasma Sphere

In general, analytic solutions for the expansion of a hot plasma do not exist. However, as a preliminary to the investigations of the expansion of laser targets, the case of an expanding hot, spherical drop will be examined for which exact solutions have been derived.

Using a homogeneous, isothermal, spherical model, Kidder^{1,2,3} reaches the following results from a separation of variables solution of the hydrodynamic equations. In the Lagrangian description the specific energy e is described by:

$$e(t) = \frac{e_0}{1 + \tau^2} \quad \text{A.1}$$

$e_0 = \frac{\text{Energy}}{\text{Mass}}$ is the initial specific energy, $\tau^2 = \frac{c_0^2 t^2}{R_0^2}$, where

$2\pi R_0^2 = \left[\frac{M}{\rho_0} \right]^{2/3}$, M is the mass, and e_0 is the initial energy density; and $c^2 = c_0^2/\gamma$ where c_0 is the initial speed of sound in the plasma.

$$c_0^2 = \gamma(\gamma-1)e_0 \quad \text{A.2}$$

The expression for the material density of the plasma is:

$$\rho(r_0, t) = \frac{\rho_0}{(1+\tau^2)^{3/2}} \exp[-r_0^2/2R^2(t)] \quad \text{A.3}$$

r_0 is the Lagrangian position variable and the plasma is assumed to have expanded far enough such that the radial density profile is gaussian.

For a fully ionized plasma (ideal, Maxwellian gas) the specific energy can be expressed in terms of the plasma temperature θ :

$$e(\text{jergs/gm}) = \frac{3}{2} \cdot 0.0965 \frac{(\bar{Z}+1)}{A} \theta \text{ (keV)} \quad \text{A.4}$$

$\bar{A}$ is the average atomic weight, $\bar{Z}$ is the average atomic number, and the electron and ion temperatures are assumed to be equal $\Theta_e = \Theta_i = \Theta$. Also the electron density n of the plasma can be written as:

$$n = \frac{\rho \bar{Z}}{\bar{A}} N_0, \quad N_0 = 6.023 \times 10^{23} \quad \text{A.5}$$

Thus, in terms of a highly ionized plasma, Kidder's results become:

$$\Theta(t) = \frac{\Theta_0}{1+\tau^2} \quad \text{A.6}$$

$$n(r_0, t) = \frac{n_0}{(1+\tau^2)^{3/2}} \exp \left[-\frac{r_0^2}{2R^2(t)} \right] \quad \text{A.7}$$

These results were confirmed by a computer simulation of an expanding spherical plasma using LASNEX. Comparisons of the computer results with the analytic solutions are illustrated in Figures 1 and 2.

The above expressions may be used, as a zeroth approximation to the expansion of a general plasma, to estimate the values of the initial parameters which correspond to specific expansion time scales. As an example, the mass of a hot spherical plasma will be calculated for 1 nanosecond time scales.

The maximum density ρ_m lies at the center of the expanding sphere at $r_0 = 0$. Therefore:

$$\rho_m = \frac{\rho_0}{(1+\tau^2)^{3/2}} \quad \text{A.8}$$

for $\gamma = 5/3$, $c^2 = (\gamma-1)e_0 = 2/3 e_0$, and

$$\tau^2 = \frac{4\pi}{3} \left[\frac{\rho_0}{M} \right]^{2/3} \frac{E}{M} t^2 \quad \text{A.9}$$

where E is the total initial energy and M is the total mass of the plasma. From Equations A.8 and A.9 the time at which the plasma goes underdense to laser light t_c can be solved for in terms of the initial parameters:

$$t_c^2 = \left[\frac{\rho_o}{\rho_c} \right]^{2/3} - 1 \left[\frac{4\pi}{3} \left[\frac{\rho_o}{M} \right]^{2/3} \frac{E_o}{M} \right]^{-1} \quad A.10$$

ρ_c is the material density for critical electron density. For a Nd laser the critical electron density is 10^{21} electrons/cm³ and $\rho_c = 2.9 \times 10^{-3}$ gm/cm³ for a fully ionized CH₂ plasma.

If the time at which an expanding plasma goes underdense is considered as a characteristic time, then for a fully ionized CH₂ plasma ($\bar{A} = 4.67$, $\bar{Z} = 2.67$), at an initial density of 1 gm/cm³, and initial energy of 2 joules (2×10^{-9} jerks) with a time scale of 1 nanosecond (10^{-1} sh) the total mass must be:

$$M = \left[\frac{4\pi E t^2}{3 \left[\left(\frac{1}{\rho_c} \right)^{2/3} - \left(\frac{1}{\rho_o} \right)^{2/3} \right]} \right]^{3/5} \quad A.11$$

$$\approx 8.9 \times 10^{-8} \text{ gm}$$

B. The Expansion of a Hot Plasma Disk

There are no exact solutions for the expansion of a hot plasma disk. However, approximate solutions for the behavior of an expanding disk can be obtained from the elementary hydrodynamic theories of one dimensional and three dimensional expansion, with the aid of computer simulations.

Consider a hot plasma disk of uniform density and temperature, with its axis along the z direction and its center at $z_0 = 0$ which is suddenly permitted to expand.

Initially the thin disk expands principally in the axial direction along the positive and negative z axes⁴. According to the theory of linear (1 dimensional)

expansion; the sides of the disk blow off and an axial rarefaction wave propagates toward the center of the disk from each side with a maximum velocity equal to the initial speed of sound in the plasma c_0^5 .

$$\begin{aligned} c_0^2 &= \gamma(\gamma-1)e_0, \quad \text{for } \gamma = 5/3 \\ &= .16 \frac{(\bar{Z}+1)}{A} \Theta_0(\text{keV}), \end{aligned} \quad \text{B.1}$$

Moreover, for adiabatic expansion, the sides rapidly achieve their asymptotic velocity v :^{6,7}

$$\begin{aligned} v &= \frac{2}{\gamma-1} c_0, \quad \text{for } \gamma = 5/3 \\ &= 3c_0 \end{aligned} \quad \text{B.2}$$

For a CH_2 plasma at an initial temperature of 300 ev:

$$\begin{aligned} c_0 &= 1.94 \times 10^{-1} \text{ cm/sh} \\ v &= 5.82 \times 10^{-1} \text{ cm/sh} \end{aligned}$$

The computer results gave $v = 5.9 \times 10^{-1} \text{ cm/sh}$.

The maximum density lies at center of the expanding disk. A heuristic argument shows that the peak density is inversely proportional to time for linear expansion.

Initially, the disk expands very little in the radial direction; and the radius can, therefore, be considered as essentially constant. Consequently, the surface area of the sides of the disk remains constant and the volume V increases as a function of the thickness and subsequently of time alone.

$$\begin{aligned} V &= 2\pi R_0^2 z \\ &= 2\pi R_0^2 vt \end{aligned} \quad \text{B.3}$$

Thus, for a plasma behaving as a perfect gas $\rho \propto \frac{1}{V} \propto \frac{1}{t}$. Computer simulations reveal that this is in fact the correct time dependence for the peak density during the initial stages of expansion.

The peak density may, therefore, be written in the following form similar to the corresponding expression for three dimensional expansion (A.8).

$$\rho_m(t) \cong \frac{1}{2} \frac{\rho_0}{1+\tau} \quad \text{B.4}$$

$\tau = t/t_0$, where t_0 is an appropriately chosen time characteristic of the particular disk; and the 1 in the denominator is included for normalization at $t = 0$. The expression for the peak density contains the factor of 1/2 as the disk is expanding linearly in two directions at once. A reasonable value for the characteristic time t_0 is the time required for the axial rarefaction wave to propagate to the center of the disk. A comparison of this expression with the results of computer simulations show that this choice of t_0 yields fairly accurate results. (See Fig. 3)

A CH_2 plasma disk with a radius of 10^{-2} cm and a total thickness of 2.84×10^{-4} cm at a temperature of 300 ev and initial density of 1 gm/cm^3 has a sound speed $c_0 = 1.94 \times 10^{-1}$ cm/sh which yields a time:

$$\begin{aligned} t_0 &= \frac{z_0}{c_0} \\ &= \frac{1.42 \times 10^{-4}}{1.94 \times 10^{-1}} = .73 \times 10^{-3} \text{ sh} \end{aligned} \quad \text{B.5}$$

There is also a simple approximation for the time dependence of the peak temperature θ_m for linear expansion. An analysis of computer models leads to an expression of the form:

$$\theta_m(t) \approx \frac{1}{2} \frac{\theta_0}{(1+\tau)^{1/2}} \quad \text{B.6}$$

as previously defined. This approximate relation compares well with the behavior of θ_m in the computer simulations. (See Fig. 4)

Thus far, only the initial stages of the expansion have been considered, for which the radius may be assumed constant. However, the radial surface also expands and a rarefaction wave propagates in the radial direction toward the center of the disk at a maximum instantaneous velocity equal to the sound speed at the wave front. The radial velocities take somewhat longer to reach their asymptotic values (See Figure 5) and consequently their effects upon the initial expansion of the disk may be neglected.

The time required for the radial rarefaction wave to travel to the center of the disk t_r is easily calculated from the instantaneous sound velocity at the wave front which is a function of the temperature of the relatively undisturbed region in front of the wave, the expression for which has been previously derived (B.6):

$$\int_0^{t_r} c(t) dt = R_0$$

where

$$c^2(t) = \frac{c_0^2}{2(1+\tau)^{1/2}} \quad \text{B.7}$$

Thus;

$$\frac{c_0}{\sqrt{2}} \int_0^{t_r} \frac{1}{(1+\tau)^{1/4}} dt = R_0 \quad \text{B.8}$$

If $t_r \gg t_0$ ($R_0 \gg z_0$) then B.8 leads to the following expression for the time t_r at which the radial rarefaction wave reaches the center:

$$t_r = \frac{t_0}{4} \left(\frac{3R_0}{t_0 c_0} \right)^{4/3} = \frac{z_0}{4c_0} \left(\frac{3R_0}{z_0} \right)^{4/3} \quad \text{B.9}$$

For the previously described CH_2 plasma disk $t_r = 2.80 \times 10^{-1} \text{ sh} \gg 9 \times 10^{-4} \text{ sh} = t_0$.

Eventually a time is reached when the radial component of the expansion begins to have a significant effect on the total expansion; and the mode of the expansion crosses over from one dimensional to three dimensional expansion.

The crossover occurs when the rate of hydrodynamic work in the radial direction is equal to rate of hydrodynamic work in the axial direction. It can be shown that the crossover time t_x is closely related to the time at which the radial rarefaction wave reaches the center of the disk.

Basically, the condition for crossover can be written for an isotropic pressure as:

$$A_R u' = A_Z v' \quad \text{B.10}$$

A_R is the radial surface area of a small Lagrangian volume element about the center of the disk; A_Z is the axial surface area which can be considered to be constant; v' is the asymptotic axial expansion velocity which is

proportional to the edge velocity⁸; and u' is the radial expansion velocity which is a function of time (Fig. 5) evaluated at t_x .

At the rarefaction wave interface, matter is rapidly accelerated from extremely low velocities to its maximum expansion velocity. (Fig. 5) Therefore, if a small volume element is chosen about the center which expands with the disk, the radial velocity u' at the edge of the volume element will be very much smaller than the corresponding axial velocity until the radial rarefaction wave reaches the volume element. After the arrival of the rarefaction front at the center, u' accelerates quickly to its asymptotic value which is comparable in magnitude to v' (Fig. 5).

Very crudely, at the time t_r that the rarefaction wave reaches the small volume element at the center, the surface areas in the radial and axial directions will be of the same order of magnitude. t_r is roughly a small multiple of the time required for a sound wave to propagate radially to the center at the initial sound speed. The axial velocity is also a small multiple of the sound speed. Therefore, at t_r the disk will have expanded axially a distance roughly of the same magnitude as the radius ($z = 3c_0 t \sim 3R_0$). Assuming that the expansion of the small Lagrangian volume element scales as the expansion of the disk itself, the radial and axial surface areas (A_R and A_Z) at this time will be comparable in magnitude as well as radial and axial velocities (u' and v'). Thus, the rates of hydrodynamic work will also be of comparable magnitude.

This crude estimate of the crossover time is confirmed by computer simulation. A comparison of the results for three separate computer runs are tabulated below:

Run Name	$t_x(\text{sh})$ computer	$t_x(\text{sh})$ theory
DISK 08	$\sim 2.5 \times 10^{-1}$	2.8×10^{-1}
DISK 09	$\sim 6.5 \times 10^{-1}$	6.9×10^{-1}
DISK 10	$\sim 4.48 \times 10^{-1}$	4.5×10^{-1}

At the crossover time the expressions for the peak density and temperature change abruptly. Kidder's calculations for spherical expansion should approximate these relations for the three dimensional expansion of the disk. However, the spherical results do not wholly apply to the case of an expanding disk as the expansion is cylindrical rather than spherical. Heuristically, the relation $\rho_m \propto \frac{1}{t}$ for linear expansion should change to $\rho_m \propto \left(\frac{1}{t}\right)^3$ for three dimensional expansion. This result agrees with the equations for spherical expansion; but, by similar reasoning $\theta_m \propto \frac{1}{t^{1/2}}$ changes to $\theta_m \propto \frac{1}{t^{3/2}}$ for the three dimensional plasma expansion. The computer results verify that for cylindrical expansion

$$\rho_m \cong \frac{\rho_x}{(t/t_x)^3} = \frac{1}{2} \frac{\rho_o}{(1+t_x/t_o)} \frac{1}{(t/t_x)^3} \quad \text{B.11}$$

$$\theta_m \cong \frac{\theta_x}{(t/t_x)^{3/2}} = \frac{1}{2} \frac{\theta_o}{(1+t_x/t_o)^{1/2}} \frac{1}{(t/t_x)^{3/2}} \quad \text{B.12}$$

ρ_x and θ_x are the densities and temperatures at crossover. These relations are plotted with the computer results in Figures 3 and 4.

C. Low Intensity Laser Irradiation of a Plasma Disk

The behavior of the characteristic parameters of an expanding plasma shall now be examined for the more realistic case of the expansion of a plasma disk heated with a low intensity $\sim 10^{14}$ w/cm² pulse of laser light. Furthermore, approximate relations for describing the behavior of the laser plasma parameters of principal interest--the peak density, the scale height, and the maximum temperature, shall be derived.

A great deal of research has been done on the comparatively simple, one dimensional problem of the laser heating and subsequent expansion of a spherical target^{9,10,11}. However, little work has been done on the two-dimensional problem of the investigation of the behavior of the heating and expansion of a plasma disk¹².

Computer simulations of the laser heating and expansion of their CH₂ disks were performed in order to determine approximate relations for the behavior of the three characteristic parameters of the expanding plasma. The initial dimensions and parameters of the successful computer "experiments" are tabulated below:

Run Name	TARGET			LASER			
	Radius R ₀ (cm)	Thickness Z ₀ (cm)	Mass M (gm)	Input Energy (joules)	Energy Absorbed (joules)	Intensity Profile (σ)	Pulse Length tp (sh)
LASD09A	.005	5.68x10 ⁻⁴	4.45x10 ⁻⁸	.828	.828	Gaussian (36.5μ)	1x10 ⁻²
LASD10A	.01	1.42x10 ⁻⁴	4.45x10 ⁻⁸	.828	.828	Gaussian (75μ)	2x10 ⁻²
LASD12A	.01	2.84x10 ⁻⁴	8.9x10 ⁻⁸	.828	.828	Gaussian (75μ)	1x10 ⁻²
LASD20A	.01	5.68x10 ⁻⁴	17.8x10 ⁻⁸	4.8	3.45	Gaussian (75μ)	1x10 ⁻²
LASD22A	.01	1.42x10 ⁻⁴	4.45x10 ⁻⁸	4.8	3.37	Gaussian (75μ)	1x10 ⁻²
LASD24A	.01	1.42x10 ⁻⁴	4.45x10 ⁻⁸	1.2	.988	Constant Linear	1x10 ⁻²
LASD01Q	.01	1.42x10 ⁻⁴	4.45x10 ⁻⁸	1.2	.88	Gaussian (75μ)	1x10 ⁻²

The laser intensity profiles for all of the runs were approximately gaussian except for LASD24A which had a constant linear profile. Moreover, the last four runs utilized the 10 laser RAYTRACE option in LASNEX 13 which gave a somewhat smoother intensity distribution than the preceeding runs.

Very little is known about the actual process by which a material is converted into an ionized plasma by a beam of high intensity radiation¹¹. However, it has been assumed that the process of heating, vaporization and ionization of a laser target is completed in a fraction of the laser pulse length and can therefore be neglected. All of the computer runs listed above performed partial ionization physics with an initial material density of 1 gm/cm^3 corresponding to an electron density of $5.3 \times 10^{22} \text{ electrons/cm}^3$ at an initial temperature of 10 ev.

A real plasma absorbs light through inverse bremsstrahlung and by the stimulation of various absorptive instabilities in the vicinity of the critical electron density n_c . All radiation which reaches the critical density, that is not absorbed, is reflected back through the plasma. For Nd laser light $n_c = 10^{21} \text{ electrons/cm}^3$.

In the non-RAYTRACE computer simulations all of the laser energy incident upon the plasma is absorbed, either by inverse bremsstrahlung or by anomalous absorption. Furthermore, a highly efficient absorption mechanism is assumed which absorbs all light energy which reaches the critical density. For RAYTRACE lasers the laser light is refracted and reflected by the plasma as well as absorbed. $\sim 90\%$ of the laser energy

that is not refracted and reaches the region of critical density is absorbed; and $\sim 10\%$ is reflected. Due to the steep initial density gradients, only a small fraction $\sim .5\%$ of the total laser energy is absorbed by inverse bremsstrahlung.

As an approximate model, consider a CH_2 plasma disk at an initial temperature of 10 eV with a uniform electron density of 5.3×10^{22} electrons/cm³ (1 gm/cm^3 material density), which is exposed to a 100 picosecond Nd laser pulse incident upon one side of the disk along the negative z axis. For LASD01Q the laser pulse has a trapezoidal shape, rising to a maximum power of 1.6×10^{-7} jerks/sh.

This corresponds to a total energy of 1.2×10^{-9} jerks (1.2 joules). For a gaussian intensity profile with $\sigma = 75 \mu$ only $\sim 85\%$ of this energy enters the plasma of which .88 joules is absorbed.

Initially, most of the plasma is overdense and absorption occurs only near the surface of the exposed side. The thin layer of plasma along the irradiated surface is strongly heated and it ablates off sending a shock wave into the plasma disk. Behind the shock front the plasma is rapidly heated; the plasma expands quickly reaching an asymptotic expansion velocity; and the plasma density drops permitting the deeper penetration of laser light.

For a linear incident intensity profile (LASD24A), the irradiated side of the disk is uniformly heated, except near the edges; and the expansion of the disk approximates a linear expansion in a single direction. At first the thin plasma disk expands almost entirely in the direction of the incident beam. A laser induced shock wave propagates through the plasma in the negative z direction. At the shock front interface compression and heating occur. The compression at the shock front increases the density to several times the initial density depending upon the intensity of the shock. The peak density, thus, moves with the interface through the plasma during the pulse. Figure 6 illustrates the propagation of the shock wave. Behind the shock front the plasma is rapidly expanding; and the peak density comes to rest near the far side. After the end of the laser pulse the maximum density drops rapidly. For a linear intensity profile the density falls as the inverse of time-- $\rho \propto \frac{1}{t^{1.2}} \sim \frac{1}{t}$.

However, for a more realistic gaussian intensity distribution, the plasma is not heated uniformly. The gaussian profile heats the center of the exposed side considerably more than the edges. This

results in a greater expansion velocity and a stronger, faster shock wave at the center. These effects create an indentation in the radial density profiles. The processes of shock wave propagation and plasma expansion are similar to those for the linear intensity profile; but the relations describing the behavior of the plasma are no longer approximated by the expressions for a linear expansion. Computer simulations show that, although ρ is initially proportional to $\frac{1}{t}$ for the linear profile, $\rho \propto \frac{1}{t^2}$ for a gaussian intensity distribution. (See Fig. 7).

The region of peak density in the expanding disk lies roughly in the shape of a torus (see Fig. 8). The volume of a torus is given by:

$$V = 2\pi^2 b r^2$$

C.1

b is essentially constant for times less than the time required for the radial rarefaction wave to reach the center t_r . Furthermore, the radial and axial expansion velocities near the edge of the disk are comparable in magnitude.

Thus, $r \propto t$ and $\rho \propto \frac{1}{V} \propto \frac{1}{t^2}$ as indicated by computer results. For the models considered, the initial compression at the shock front interface increases the maximum density by a factor of 2 or 3; but by the end of the pulse, the peak density is approximately its initial value. Consequently,

estimates for the peak material and electron densities may be written:

$$\rho_m \approx \frac{\rho_o}{(t/t_p)^2} \quad C.2$$

$$n_m \approx \frac{n_o}{(t/t_p)^2} \quad C.3$$

ρ_o and n_o are the initial material and electron densities and t_p is the pulse length. These estimates are compared with the computer results in Figure 7.

As the shock front propagates into the disk, the electrons are heated to a peak temperature. As a consequence of the long electron-ion thermalization times $\sim t_p^{15}$, the electron temperature is initially somewhat greater than the ion temperature. After the laser is turned off the electrons and the ions cool rapidly due to the expansion of the plasma. For a gaussian intensity profile the computer results show that $\Theta_e \sim \frac{1}{t^{.8}} \sim \frac{1}{t}$. (See Fig. 9).

The maximum electron temperature, which occurs during the laser pulse, appears to be mostly dependent upon the incident laser intensity and the magnitude of the resulting shock. As a rough estimate, the maximum temperature at the end of the pulse will be comparable to the maximum average temperature in the plasma. Thus, the peak electron temperature may be estimated as:

$$\Theta_e = \frac{\Theta_o}{t/t_p} \quad C.4$$

Θ_0 is the maximum average temperature derived from the total energy absorbed by the plasma and its total mass (Eq. A.4).

Furthermore, computer simulations indicate that the maximum expansion velocity is a function of the peak temperature. By the time the ablating side reaches its asymptotic velocity, laser heating will only have occurred in a thin layer of the disk which is raised to a temperature near Θ_m . The ablation velocity is determined by the high temperature of this linearly expanding layer. Thus the maximum expansion velocity for a laser heated disk may be approximated by:

$$v_m = 3c_m \quad \text{C.5}$$

$$c_m^2 = .16 \frac{\bar{Z}+1}{A} \Theta_m \text{ (kev)} \quad \text{C.6}$$

Run Name	$\Theta_{\max}$ (kev)	Computer v_m (cm/sh)	Theory v_m (cm/sh)
LASD09A	.8	9.8×10^{-1}	9.5×10^{-1}
LASD10A	.58	8.7×10^{-1}	8.1×10^{-1}
LASD12A	.45	8.0×10^{-1}	7.14×10^{-1}
LASD20A	2.1	1.7	1.54
LASD22A	~ 2.1	2.0	1.54
LASD01Q	1.9	1.3	1.01

In general, it is apparent that the approximations for the behavior of the expanding plasma are more accurate for the lower intensities and weaker shocks.

As the plasma disk expands the axial density gradient quickly assumes the shape of an exponentially decreasing function of z ; (See Fig. 10). The ablating side expands at its asymptotic velocity; in the positive z direction and the region of peak density recedes in the negative direction due to the momentum imparted by the laser pulse.

An important parameter of expanding plasmas is the scale height L , which is a measure of the slope of the density gradient. It is defined for exponentially decreasing density gradients as:

$$-\frac{1}{L} = \frac{1}{n} \frac{\partial n}{\partial z} \quad \text{C.7}$$

Lelevier¹⁶ shows that, for an isothermal, one-dimensional expansion, the density is a decreasing exponential of z .

$$n(z) = n_0 e^{-\frac{z+c_0 t}{c_0 t}} \quad \text{C.8}$$

c_0 is the maximum average speed of sound in the plasma;

$$c_0^2 = .16 \frac{\bar{Z}+1}{A} \frac{E_0}{M} \quad \text{C.9}$$

where E_0 is the total energy absorbed. Lelevier also shows that this result approximates the axial density profile for an adiabatic expansion for which $\gamma \sim 1 \sim 5/3$. Consequently, the previous expression may be used to estimate the scale height for the expanding laser heated plasma. From C.7 and C.8:

$$\frac{1}{n(z)} \frac{\partial n(z)}{\partial z} = - \frac{1}{c_0 t} \quad \text{C.10}$$

This result for the slope of the density gradient is independent of the density at which it is evaluated. Therefore, the scale height for a linearly expanding disk is:

$$L = c_0 t \quad \text{C.11}$$

A comparison of the values for the scale height predicted by this result with those evaluated from the density gradients in the computer simulations, show that this is a fair estimate for the scale height (See Fig. 11). As before, this approximation is most accurate for low incident intensities and weak shocks.

D. High Intensity Irradiation of a Prepared Plasma Target

Computer simulations were performed of the high intensity irradiation of plasma targets, such as those prepared in Part C, in order to determine the behavior of the characteristic parameters that may be expected for real experiments.

The plasmas prepared in the previous section go underdense in approximately 1 nanosecond; and they cross over from linear to three-dimensional expansion somewhat later $t_x \gtrsim 3 \times 10^{-1}$ sh. For nanosecond time scales the scale heights are long $\gtrsim 50 \lambda_0$ (the wavelength for Nd light is $\lambda_0 = 1.06\mu$); and the electron temperatures are low ≤ 50 eV. Computer simulations were executed in which these linearly expanding plasma targets were irradiated with a high intensity, 100 picosecond, trapezoid laser pulse of ~ 42 joules directed along the negative z axis. The prepared targets were pulsed at times corresponding to maximum electron densities of $2n_c$, n_c , and $1/4 n_c$ with both the parallel and focussed ray options of RAYTRACE. For the target LASD01Q, the incident times were 5.5×10^{-2} sh, 1.0×10 sh and 2×10^{-1} sh, respectively. The intensity profiles for the high intensity laser beams were roughly gaussian with a σ of 75μ . The peak intensities ranged up to $\sim 2 \times 10^{16} \text{ W/cm}^2$, with a maximum average intensity of $\sim 4 \times 10^{15} \text{ W/cm}^2$. The successful runs are tabulated below:

Run Name	Target Name	n_m	Incident time t_i (sh)	Initial (μ) Scale height	Peak $\frac{jerk}{sh}$ Power	Beam	Intensity Profile
LASD02S	LASD01Q	$2n_c$	5.5×10^{-2}	96	6.6×10^{-6}	Parallel ray	Gaussian $\sigma=75\mu$
LASD03R	"	$2n_c$	5.5×10^{-2}	96	"	Focussed ray	"
LASD04R	"	$2n_c$	1.0×10^{-1}	167	"	Parallel ray	"
LASD08S	"	$\frac{1}{4}n_c$	2.0×10^{-1}	298	"	Parallel ray	"

The laser energy which enters the plasma is either reflected, refracted, or absorbed. As before laser light is absorbed by the plasma through inverse bremsstrahlung, anomalous absorption in the vicinity of the critical density; and by an artificial absorption mechanism, DUMPALL, which absorbs 90% of all light which reaches the critical density and reflects 10%. In the computer simulations, between 23% and 56% of the incident light energy is absorbed by the plasma targets and the rest escapes. Moreover, on examination of the computer calculations of the spatial distribution of the escaping light energy shows that most of the light is refracted and only 10 - 20% is reflected. (See Fig. 12).

By the times that the plasma target is exposed to the main pulse, the initial disk has expanded a considerable distance in the positive z direction $\sim 10^{-1}$ cm. The density is a smoothly decreasing exponential function of z with a long scale height; and the electron temperature is fairly uniform in front of the peak densities. (See Fig. 10 and Fig. 13).

Heat conduction is very efficient in the expanded plasma due to the long mean free paths of the thermal electrons. Consequently, the high intensity laser pulse is absorbed fairly uniformly throughout the plasma. Initially, most of the laser energy is absorbed by inverse bremsstrahlung until the electron temperature is raised to a sufficient level $\theta_e \gtrsim 200$ ev for $L \sim 100\mu$. Then, for regions of the plasma at or near critical density, anomalous absorption and DUMPALL dominate; and strong absorption occurs in the vicinity of the critical density. Although the heat is quickly dispersed throughout the plasma, some compression occurs in this region resulting in greater electron densities, and higher electron temperatures. (See Fig.'s 14 and 15).

In all of the high intensity computer runs there was a fairly significant increase in the peak electron density due to compression. In fact, for the $1/4 n_c$ run, LASD08S, the maximum electron density rises above critical

The compression at peak density also tends to steepen the density gradient and shorten the scale height. However, because of the smooth absorption of the laser energy and the short duration of the pulse, the scale heights at or below n_c do not change appreciably during the pulse except at the very center of the disk where the effect is most pronounced. (See Fig. 14). Consequently, the scale height of the plasma target at the time of incidence t_i of the main pulse may be used as a close approximation to the actual scale height during the pulse for most of the target. Therefore, from equation C.11

$$L = c_0 t_i$$

c_0 is the maximum average sound speed of the prepared target calculated from C.9. The computer results confirm that the scale height does not deviate from this estimate by more than $\sim 10\%$.

Due to uniform heating, the plasma becomes nearly isothermal except for the region of high absorption near n_c where the peak temperature rises to a factor of 2 or 3 greater than the base electron temperature. (See Fig. 15). An estimate for a characteristic plasma temperature may be derived from the total energy absorbed and the total mass of the plasma. The maximum average temperature Θ_A for the plasma heated by the main pulse is:

$$\Theta_A = 6.91 \frac{E_0}{M} \text{ (keV)} \quad \text{D.2}$$

Here, E_0 is the total light energy absorbed during the high intensity pulse. If no energy were lost through expansion and if the electron and ion temperatures were equal this result would be exact. However, for real laser plasmas energy is lost; and, because of the long electron-ion thermalization times, the electrons are considerably hotter than the ions except in the region near peak density where the thermalization times are shorter. A comparison of Figures 15 and 16 shows the differences in temperature distribution. In any case Θ_A is smaller than the peak temperature and roughly comparable to the instantaneous, base electron temperature in the plasma. The temperatures for the computer simulations are compared below:

Run Name	E_0 (joules)	Mass M (gm)	Peak (keV) Temperature	Average base (keV) Temperature	Θ_A (keV)
LASD02S	17.4	4.45×10^{-8}	5.2	2	2.70
LASD03R	9.5	4.45×10^{-8}	2.2	1.1	1.48
LASD04R	26.6	4.45×10^{-8}	9.1	3.0	4.13
LASD08S	23.5	4.45×10^{-8}	4.5	2.78	3.95

Thus, the scale height and electron temperature of the laser plasma can be fairly accurately specified. These two parameters may be used to designate a laser plasma "phase space". The plasma targets that are heated by a high intensity laser pulse are characterized by the regions of the phase space that they occupy. Figure 18 shows the regions of laser plasma phase space occupied by the computer simulations. This graph also illustrates the accuracy of the approximations for the scale heights and the electron temperatures.

For most experimental purposes the region of phase space occupied by a particular plasma target determines the plasma instabilities which may be expected to occur as the target is heated by a high intensity laser pulse. The connection between the laser plasma phase space and the thresholds for the occurrence of specific plasma effects shall be discussed in the following section.

E. Plasma Phenomena Resulting From the Interaction With High Intensity Laser Light

There are many distinct plasma phenomena that may be expected to occur when a prepared plasma target is irradiated with a high intensity pulse of laser light. The principal known plasma instabilities that are stimulated by the interaction of high intensity radiation with a plasma are the Raman and Brillouin instabilities, the $2 \omega_{pe}$ instability, the decay instabilities, the oscillating two stream instability, self-focusing and filamentation.

The Raman effect is characterized by the parametric decay of the incident electromagnetic wave into an electron plasma wave and another light wave. Raman scattering occurs only for electron densities at or below $1/4 n_c$. Brillouin scattering is an analogous effect in which the light wave decays into an ion acoustic wave and another light wave. The Brillouin effect occurs for electron densities up to critical. For both the Raman and Brillouin effect, the light is scattered both backward and to the side. The $2 \omega_{pe}$ effect occurs only near $1/4 n_c$ and is a result of the decay of the electromagnetic wave into two electron plasma waves.

The parametric decay instability is characterized by the decay of the electromagnetic wave into both an ion acoustic wave and an electron plasma wave. The oscillating two stream instability is similar to the decay instability except the ion wave has zero frequency. Both the decay and oscillating two stream instabilities occur only in the vicinity of the

critical density. Finally, self-focusing and filamentation of the laser light results from the intensity dependence of the plasma index of refraction.

Several general conditions must first be fulfilled before these instabilities will be observed. First, the absorption of laser light by inverse bremsstrahlung must be negligible so that sufficient laser intensities reach the regions in which instabilities are stimulated.

Taking into account the inverse bremsstrahlung absorption and the swelling factor ¹⁷ near n_c , the intensity may be written as:

$$I = \begin{cases} I_0 3.6 (k_0 L)^{1/3} \exp \left(-\frac{16}{15} \frac{(k_0 L) 5 \times 10^{-4}}{\theta^{3/2} (\text{KeV})} \right), & n \sim n_c \\ I_0, & n < n_c \end{cases} \quad \text{E.1}$$

I_0 is the incident intensity, L is the scale height near the critical density measured in microns (μ), k_0 is the wave number of the incident light, and θ is the electron temperature in KeV. In all further calculations the intensity I is always the value given by E.1.

For neodymium laser heating collisional damping has a negligible effect on the high frequency (Raman and $2 \omega_{pe}$) effects for:²⁵

$$\nu_{ei} \ll .83 \times 10^6 \sqrt{I} \quad \text{E.2}$$

ν_{ei} is the electron-ion collision frequency. E.2 reduces to the requirement that for laser intensities $\sim 10^{15} \text{ W/cm}^2$ and electron densities up to $10^{21} \text{ electrons/cm}^3$:

$$\theta_e \gg 1 \text{ ev} \quad \text{E.3}$$

Landau damping also has an insignificant effect for:²⁵

$$\frac{\lambda_D}{\lambda_0} \ll \frac{1}{4} \quad \text{E.4}$$

$\lambda_D = 2.34 \times 10^4 \sqrt{\frac{\theta}{n}}$ (Kev) is the Debye length. For plasma temperatures of the order of a few keV this condition is met for

$$n \gg 10^{16} \text{ electrons/cm}^3 \quad \text{E.5}$$

For the computer simulations only Landau damping need be considered, but only in the low density regions of the plasma near the edges.

Much work has been done on the calculation of approximate intensity thresholds for the occurrence of the individual instabilities. These threshold conditions are generally dependent upon the scale heights and electron temperatures of the plasma. Thresholds for these plasma phenomena have been evaluated for Nd laser plasmas and are tabulated in Table 1. Some remarks on these results follow.

The intensity thresholds for the Raman effects are functions of the scale height alone¹⁸. The Brillouin thresholds, however, are in general dependent upon the scale height of the velocity gradient L_u as well as the density scale height L ¹⁹. The velocity scale height is defined as the distance through the plasma for which the velocity changes by the sound speed c_0 . For a linearly expanding plasma with an approximately linear velocity gradient²⁰, the scale height L_u is $\sim c_0 t \sim L$.

The decay instability is dependent upon the ratios of the ion and electron temperatures²¹. For long thermalization times, the computer

simulations indicate that $\theta_e > \theta_i$. However, if the Brillouin effects occur, the ion temperature may quickly rise to the level of the electrons²². In any case, the thresholds are evaluated at their maximum values, i.e., for $\theta_e \sim \theta_i$.

Self-focusing has approximate power and length thresholds rather than an intensity threshold²³. In order for self-focusing to occur, the instantaneous power of the incident beam must be sufficient to overcome the effects of diffraction and the plasma must be thick enough for the beam to focus a significant amount. Both the power and length conditions were evaluated for an electron density of $\sim 1/2 n_c$. Calculations for the computer simulations yield a self-focusing distance of z_{sf} of $\sim 80 \mu$ and a critical power P_c of $\sim 10^7$ watts which are easily exceeded. Although the thresholds for self-focusing are still in the stage of development, it is evident that this effect must be taken into account. If the power greatly exceeds the self-focusing threshold, the laser beam may be expected to break up into filaments which self-focus individually.

It must be remembered that all of these thresholds are order-of-magnitude estimates²⁵. In fact, the convention for the relation between the intensity I and the effective electric field of the electro-magnetic wave varies by a factor of 4 from one derivation to another depending on whether the peak or RMS value is used. The results which are graphed and tabulated reflect this additional uncertainty which is somewhat representative of the overall accuracy of the thresholds. Furthermore, the derivation of these thresholds neglect the possibility that other effects may be occurring simultaneously. Competition between two or more instabilities

may significantly affect the threshold values. For example, resonant absorption²⁶, another effect which has no intensity threshold, is strongly competitive with the decay instabilities in the vicinity of critical density. Moreover, if Raman or Brillouin effects are occurring, they can modify the electron and ion temperatures and the scale height and therefore affect thresholds for any other instabilities occurring simultaneously. Finally, other instabilities may exist, which have not yet been discovered, that might also have to be considered.

These approximate thresholds do, however, give some idea of where one should begin looking for particular phenomena. In Figure 18 the intensity thresholds are plotted in laser plasma phase space along with the ranges of the characteristic parameters of the computer simulations. A comparison of the regions occupied by the simulations with the regions designated by the iso-intensity curves, gives an idea of the instabilities that may be expected to occur for similar experiments.

Conclusion:

The conditions for the occurrence of plasma instabilities which result from the interaction of high intensity laser light are dependent upon certain characteristic parameters of the plasma target. Approximate expressions have been derived for a linearly expanding plasma which is produced by the low intensity irradiation of a small, thin disk shaped target. The most useful relations are summarized below:

The approximate size of the plasma target corresponding to specific line scales may be determined by:

$$M \approx \frac{4\pi E_o t^2}{3 \left[\left(\frac{1}{p_c} \right)^{2/3} - \left(\frac{1}{p_o} \right)^{2/3} \right]} \quad \text{A.11}$$

E_o is the total energy absorbed, M is the mass, and p_c and p_o are the critical and initial material densities.

The crossover time t_x from linear to three dimensional expansion may be estimated by:

$$t_x \approx \frac{z_o}{4c_o} \left(\frac{3R_o}{z_o} \right)^{4/3} \quad \text{B.9}$$

z_o is the initial thickness of the disk, R_o is the initial radius, and:

$$c_o^2 (\text{cm/sh}) = .16 \left(\frac{\bar{Z} + 1}{\bar{A}} \right) \theta_o (\text{keV}) \quad \text{B.1}$$

$$\theta_o = 6.9 \frac{\bar{A}}{\bar{Z} + 1} \frac{E_o}{M} \left(\frac{\text{jerks}}{\text{gm}} \right) \quad \text{A.4}$$

The maximum electron densities and temperatures for a linearly expanding disk, that has been heated with a Gaussian-like intensity distribution, are approximated by:

$$n \approx \frac{n_o}{\left(t/t_p \right)^2} \quad \text{C.3}$$

$$\theta_e \approx \frac{\theta_o}{t/t_p} \quad \text{C.4}$$

t_p is the low intensity pulse length and the initial electron density is:

$$n_o = 6.023 \times 10^{23} \frac{\bar{Z}}{\bar{A}} \rho_o \quad \text{A.5}$$

The scale height for the linearly expanding disk is approximated by:

$$L = c_o t \quad \text{C.10}$$

Moreover, the computer simulations show that when these prepared targets are exposed to a high intensity laser pulse, the scale height may be estimated by the scale height of the target at the line of incidence, and the electron temperature level is approximated by:

$$\theta_A = 6.9 \frac{\bar{A}}{\bar{Z} + 1} \frac{E_o}{M} \quad \text{D.2}$$

Here, E_o is the total energy absorbed during the main pulse.

The characteristic parameters of the laser plasma determine a phase space in which the intensity thresholds for the plasma instabilities may be plotted. A comparison of the regions of phase space occupied by particular laser plasma experiments with these plots, gives a rough idea of the plasma effects which might be expected to occur.

Finally, the phase space plots of the thresholds for the instabilities may be used, along with the relations for the behavior of the parameters of a linearly expanding plasma to design experiments of the laser heating of expanding disks for the preliminary investigation of plasma phenomena.

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CAPTIONS FOR FIGURES

- Fig. 1 Central density of spherically expanding plasma as a function of time.
- Fig. 2 Mean temperature of spherically expanding plasma as a function of time.
- Fig. 3 Peak density of thin disk illuminated from one side. Solid line is LASNEX calculation; dashed, heuristic model.
- Fig. 4 Mean electron temperature of thin disk illuminated from one side. Solid line is LASNEX calculation; dashed, heuristic model.
- Fig. 5 Shows the radial velocity (u) and the axial velocity (v) for a small volume element at the center of the disk as functions of time.
- Fig. 6 The density profiles at various times during the pulse in the vicinity of the peak density. The solid line is the initial density profile before irradiation.
- Fig. 7 The results of computer simulations compared with the time evolution of the peak density predicted by Equation C.2.
- Fig. 8 A cross section of the expanded disk cut along the axis. The original disk is illustrated for scale. The toroidal shape of the region of peak density is readily apparent.
- Fig. 9 The computer results for the time evolution of the electron temperature compared with the predictions of Equation C.4.
- Fig. 10 Electron density as a function of Z for $R = 7$ and $26.7 \mu\text{m}$.
- Fig. 11 Predicted density scale heights as a function of time.

- Fig. 12 A typical distribution for the laser light energy reflected and refracted by a prepared plasma. The laser beam is incident along the axis of the expanded disk and θ is measured with respect to this axis.
- Fig. 13 The uniformity of the electron temperature in the plasma at the time of the incidence of the high intensity pulses for various radii (7×10^{-4} cm, 2.67×10^{-3} cm).
- Fig. 14 Axial density profiles for various radii after the incidence of the high intensity pulse. Note the deep indentation at the center of the disk ($R \sim 7 \times 10^{-4}$ cm) and steepening of the density gradient in the vicinity of the peak densities which is a result of compression.
- Fig. 15 A comparison of this graph with Fig. 13 illustrates the characteristic behavior of the plasma when heated by a high intensity laser pulse.
- Fig. 16 A comparison of Fig. 15 with Fig. 14 shows the difference in the electron and ion temperatures.
- Fig. 17 The electron temperature range and rough scale height are plotted at several different times after the incidence of the pulse for different computer simulations. 1 corresponds to the temperatures and scale heights at incidence. Note that the scale height does not change much from the initial value (1). Also plotted on this graph are the predicted values of θ_A for the computer simulations. These are represented by the half filled circles.

Fig. 18 This diagram summarizes the plasma instability thresholds as functions of laser intensity, electron temperature, and scale height. (Note: the self-focusing and filamentation thresholds are omitted.) These plots may be easily reproduced from the equations in Table 1. Also, displayed on this graph, for reference, are the regions of phase space occupied by the computer simulations from Fig. 16.

TABLE 1*

<u>INSTABILITY</u>	<u>ELECTRON DENSITY</u>	<u>GENERAL FORM OF THRESHOLD</u>	<u>THRESHOLD FOR Nd</u>
Raman Sidescatter ¹⁸	$n = \frac{1}{4} n_c$	$(k_o L) \left(\frac{v_{os}}{c} \right)^{3/2} \gtrsim 1$	$\left[I \left(\frac{w}{cm^2} \right) \right]^{3/4} L(\mu) > 6 \times 10^{12}$
Raman Backscatter ¹⁸	$n = \frac{1}{4} n_c$	$(k_o L) \left(\frac{v_{os}}{c} \right)^{3/2} \gtrsim 1$	$\left[I \left(\frac{w}{cm^2} \right) \right]^{3/4} L(\mu) > 6 \times 10^{12}$
Raman Backscatter ¹⁸	$n < \frac{1}{4} n_c$	$(k_o L) \left(\frac{v_{os}}{c} \right)^2 \gtrsim 1$	$I \left(\frac{w}{cm^2} \right) L(\mu) > 2 \times 10^{17}$
Brillouin Backscatter ^{18, 28}	$n \lesssim n_c$	$\frac{2I}{1 c \theta} k_o L \gtrsim 5$	$I \left(\frac{w}{cm^2} \right) \frac{L(\mu)}{\theta(keV)} > 2 \times 10^{15}$
Brillouin Sidescatter ¹⁸	$n \leq n_c$	$\left(\frac{v_{os}}{v_{te}} \right)^2 \left(\frac{L}{L_u} \right) (k_o L_u)^{4/3} \left(1 + \frac{L_u}{L} \right)^{2/3} \gtrsim 1$	$\left[I \left(\frac{w}{cm^2} \right) / \theta(keV) \right] [L(\mu)]^{4/3} > 10^{14}$
$2 \omega_{pe}$ ¹⁸	$n = \frac{1}{4} n_c$	$k_o L \left(\frac{v_{os}}{v_{te}} \right)^2 \gtrsim 1$	$I \left(\frac{w}{cm^2} \right) \frac{L(\mu)}{\theta(keV)} > 4 \times 10^{14}$

*See Appendix for definitions of variables and constants.

<u>INSTABILITY</u>	<u>ELECTRON DENSITY</u>	<u>GENERAL FORM OF THRESHOLD</u>	<u>THRESHOLD FOR Nd</u>
Oscillating Two Stream ²¹	$n \sim n_c$	$\frac{2I}{n c \theta} > 2 \left[1 + \frac{\theta_i}{\theta_e} \right] \left[\frac{4\lambda_D}{L} + \frac{2 v_e}{\omega_c} \right]$	$I \left(\frac{w}{cm^2} \right) > 10^{16} \left[\frac{3 \times 10^{-2} \theta_e^{3/2}}{L(\mu)} + \frac{10^{-3}}{\theta_{(keV)}^{1/2}} \right]$
Decay ²¹	$n \sim n_c$	$\frac{2I}{n c \theta} > 1.6 A_Y \left(\frac{4\lambda_D}{L} \right) \left(\frac{v_i}{\omega_i} \right)^{1/2} + 3.2 \left(\frac{v_i}{\omega_i} \right) \left(\frac{v_e}{\omega_c} \right)$	$I \left(\frac{w}{cm^2} \right) > 10^{16} \left[\frac{.2 \theta_{keV}^{3/2}}{L(\mu)} + \frac{4 \times 10^{-4}}{\theta_{keV}^{1/2}} \right]$
Self Focusing ²³			
Power:	$n \lesssim n_c$	$P \gtrsim P_c \sim \frac{\lambda_0^2 2 n_c c \theta}{8\pi \left(\frac{n}{n_c} \right) \left(1 - \frac{n}{n_c} \right)}$	$P \gtrsim P_c \text{ (watts)} \sim 7.76 \times 10^6 \theta(keV)$
Distance:	$n \lesssim n_c$	$z > z_{sf} \sim \left(\frac{2\pi R^2}{\lambda_0} \right) \left[\frac{P}{P_c} - 1 \right]^{-1/2}$	$z > \left(\frac{2\pi R^2}{\lambda_0} \right) \left[\frac{P \text{ (watts)}}{7.76 \times 10^6 \theta(keV)} - 1 \right]^{-1/2}$

INSTABILITY

ELECTRON DENSITY

GENERAL FORM OF THRESHOLD

THRESHOLD FOR Nd

Filamentation

$$n \lesssim n_c$$

$$\left(\frac{v_{os}}{v_{th}} \right)^2 > \left(\frac{n_c}{n} \right) \left(\frac{k_{\perp}}{k_0} \right)^2$$

$$\left(\frac{v_{os}}{v_{th}} \right)^2 > 10^{-9} \frac{n_c}{n} k_{\perp}^2$$

APPENDIX

$$A \sim 5$$

$$\gamma = \left(1 + \frac{3 \theta_i}{\theta_e}\right)$$

$$\left(\frac{v_{os}}{c}\right)^2 = \frac{I}{1.3 \times 10^{18}} (w/cm^2)$$

$$\left(\frac{2 v_{os}}{v_{te}}\right)^2 = \frac{I}{10^{16} \theta(keV)} (w/cm^2)$$

$$\left(\frac{v_{os}}{v_{te}}\right)^2 = \eta^2 \sim \frac{2I}{n c \theta}$$

ν_e = electron collision frequency

$$\frac{\nu_i}{\omega_i} = \left(\frac{\pi}{2}\right)^{1/2} \left\{ \left(\frac{\theta_e}{\theta_i}\right)^{3/2} \exp \left[-\frac{1}{2} \left(\frac{\theta_e}{\theta_i} + 3\right) \right] + \left(\frac{M_e}{M_i}\right)^{1/2} \right\}$$

R. - radius of the incident beam.

FIGURE 1

FIGURE 2

FIGURE 3

FIGURE 4

FIGURE 5

FIGURE 6

FIGURE 7

FIGURE 8

FIGURE 9

FIGURE 10

FIGURE 11

FIGURE 12

FIGURE 13

FIGURE 14

FIGURE 15

FIGURE 16

FIGURE 17

FIGURE 18

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