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Laser-Plasma Interaction in a Theta-Pinch Geometry

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LIST OF SYMBOLS

SYMBOL

$\hat{e}$	direction vector
$\parallel$	parallel to incident wave vector
$\perp$	perpendicular to incident wave vector
$\hat{n}$	scatter direction
n	order of diffraction maximum
N	fringe number
n_{ef}	number of e-foldings of SRS signal from noise
e	electron charge
q_j	species charge
Z	ion charge
j	current
m	mass
ρ	mass density
n	particle density
n_c	cutoff density for given em wave frequency
V	volume
$\underline{r}$	position
$\delta \underline{r}$	displacement

r, R	radius
$r_{\min}$	radius of density minimum
Δr	radial, refractive deflection of incident beam
L	plasma column length
z	axial position
f	focal length
λ	wavelength
λ_D	electron Debye length
r_0	classical electron radius
W	width of grating
d	groove spacing
Δs	fringe spacing
ℓ_{abs}	inverse bremsstrahlung absorption length
$\tilde{\ell}_{\text{abs}}$	dynamic, inverse bremsstrahlung absorption length
ℓ_{int}	SBS interaction length
ℓ_{conv}	convective gain length
ℓ_f	Faraday length for 90° rotation between incident and backscattered polarization planes
θ	angle between incident and backscattered polarization planes
θ	angle between ray and axis
θ_0	incident ray angle
θ_c	critical angle of incident rays for containment
$\Delta\theta$	angular, refractive deflection of incident rays
Ω	solid angle
Θ	angle between scattering source and scatter direction

t	time
τ	confinement time
t_+	duration of positive growth rate (gain)
t_-	duration of negative growth rate (damping)
v	velocity
v_e	electron thermal velocity
v_i	ion thermal velocity
c_s	ion acoustic speed
c	speed of light
v_g	pulse velocity of Sbs disturbance
v_o	jitter velocity
v_{oT}	convective threshold for jitter velocity
ν	frequency (Hz)
ω_o	incident or pump em wave frequency
ω	wave frequency (associated with plasma wave)
ω'	wave frequency (associated with scattered em wave)
ω_{es}	actual plasma wave frequency ($Re\omega$)
$\Delta\omega$	resonant plasma wave frequency from phase matching
ω_{ia}	ion-acoustic wave frequency
ω_{pe}	electron plasma wave frequency
ω_k	Bohm-Gross frequency
ω_{pi}	ion plasma wave frequency
ν_{ei}	electron-ion collision frequency
ω_r	real part of

ω_i	imaginary part of ω
γ	instability growth rate ($\text{Im}\omega$)
γ_{max}	maximum SBS growth rate
γ_0	infinite homogeneous growth rate
γ_a	absolute growth rate
γ_T^{conv}	convective threshold for infinite homogeneous growth rate
γ_T^{abs}	absolute threshold for infinite homogeneous growth rate
γ_+	positive growth rate (gain)
γ_-	negative growth rate (damping)
γ_p	plasma wave damping rate
γ_{em}	em wave damping rate
$\underline{k}_0$	incident wave vector
$\underline{k}$	wave vector (associated with plasma wave)
$\underline{k}'$	wave vector (associated with scattered em wave)
k_{es}	wave number (associated with filamentation)
k_r	real part of k
k_i	imaginary part of k
Δk	phase mismatch
κ'	rate of phase mismatch
κ	inverse bremsstrahlung absorption coefficient
ϕ	es potential function
$\underline{A}$	em potential function
$\underline{E}$	electric field
$\underline{B}$	magnetic field

$\vec{E}_0$	maximum electric field of pump
χ_j	species susceptibility
$D(\omega, k)$	dispersion relation
ϵ	wave energy quanta
$\vec{F}_{NL}$	nonlinear ponderomotive force
$\vec{f}_{NL}$	nonlinear ponderomotive force per unit volume
P	power
P_0	incident pump power
P_{th}	convective threshold for power
$\underline{S}$	Poynting vector
n	quantized wave amplitudes
S	spectral density function
S_{DN}	diagnostic noise level
S_N	initial signal, noise level
S_B	calculated SBS signal
S_I	monochromator instrument function
S_T	calculated, observable SBS signal
S_D	SBS data
S_1	wavelength resolved photovoltaic detector signal
S_2	wavelength integrated photovoltaic detector signal
P	pressure
T	temperature
δ	ratio of specific heats
μ	index of refraction

S	CO_2 laser stress level
β	ratio of local plasma kinetic energy to external magnetic field energy
α	$1/k\lambda_D$
α	alpha particle
$\sigma_T(\theta)$	Thomson differential cross section
Λ	maximum impact parameter in units of impact parameter necessary for 90° deflection in Coulomb scattering

LASER-PLASMA INTERACTION IN A THETA-PINCH GEOMETRY

by

William Thomas Armstrong

ABSTRACT

Prompt stimulated Brillouin scatter (SBS) is studied in an experiment wherein a high power, pulsed CO_2 laser irradiates an independently produced, theta-pinch plasma. SBS does not significantly affect laser heating of the plasma.

The SBS signal consisted of a single, strong pulse occurring within ± 15 ns of the incident beam's peak power. The signal displayed some intermittency and shot-to-shot power fluctuation. Phase mismatching in the SBS coupling is attributed to a strong axial temperature gradient. A brief period of reduced temperature gradient, and corresponding enhanced SBS coupling, occurs near peak incident power. The constraint of significant SBS gain to this time period is consistent with the single pulse characteristic and power fluctuations of the SBS signal.

The SBS signal was red shifted by $\sim 51 \text{ \AA}^0$ which corresponds to the ion-acoustic frequency. A theoretical fit of the SBS line shape gave a lower bound for the number of e-foldings of the backscattered signal from the noise level, $n_{ef} \geq 14$. SBS power measurements established an upper bound of $n_{ef} \leq 24$. With a gain length of 0.9 cm, these limits on n_{ef} correspond to an interaction length of $13 \text{ cm} \leq \ell_{int} \leq 22 \text{ cm}$. This interaction length is consistent with the high power focus region where the SBS threshold is exceeded. The possible effects of Faraday rotation on SBS are discussed.

Measurements of density profiles and temperature histories permitted examination of laser refraction, local heating and net absorption. Refractive containment of the CO_2 laser beam by an on-axis density minimum was observed at early times during the laser pulse. However, refractive containment was lost at late times due to the diffusive loss of the density minimum. Classical modeling of the expected heating required "bleached" absorption to account for the observed heating. A plasma absorptivity of $\sim 46\%$ was inferred from calorimetry measurements at 250 mtorr fill pressure. These results confirm that classical heating and refraction dominated the laser-plasma interaction.

1. INTRODUCTION

In the study of plasma physics the interaction of waves and plasmas is of fundamental importance. Significant studies of wave interaction with plasmas have been performed with ionospheric, Q-machine, Tokamak and numerous other plasmas. In the last ten years great interest has developed in the use of lasers to prepare thermonuclear plasmas. In such laser-plasma systems the detailed study of the interaction of an electromagnetic (em) wave in an underdense plasma has continued to be of central interest. The coupling of the laser energy to the plasma can be enhanced or reduced depending on the dominance of various wave interactions in the plasma. The work presented in this thesis results from an experimental effort at the Los Alamos Scientific Laboratory to study the laser-plasma interaction as a possible alternative heating method for theta-pinch plasmas.¹

The experiment consisted of a high power CO_2 laser directed axially into the meter-long Scylla I-C theta-pinch.² The facility was located within the CTR Division and saw active use from September 1975 through November 1976. The original areas of study delineated in the experimental proposal were: 1) Propagation, 2) Absorption, 3) Heating and 4) Anomalous Backscatter. The usage of "anomalous scatter" refers to deviation from Thomson (thermal) scattering. The study of anomalous backscatter as a decoupling mechanism in laser heating of plasmas, was designated for this thesis research and forms the central theme of this work. However, during the evolution of the experiment it was found necessary to investigate

various plasma parameters which were of direct interest to the anomalous backscatter study. Through these related examinations the first three topics listed above came into consideration and will also be discussed in this thesis.

The second chapter presents a synopsis of the experiment. A general discussion of the laser-plasma interaction is followed by a description of the experimental arrangement and diagnostic implementation. A brief summary of the experimental results concludes the chapter. The next four chapters address each of the topics mentioned above. The chapters begin with anomalous backscatter and proceed through propagation, absorption and heating. Within each of these chapters a detailed description is given of the diagnostic apparatus, data and analysis pertinent to the respective topic. In the seventh chapter results and conclusions are discussed.

II. EXPERIMENT SYNOPSIS

A. Motivation

The use of high power lasers in heating plasmas to thermonuclear conditions can be presently divided into two general classes.

1. Laser Imploded Pellets

The first is laser imploded pellets (LIP) or inertial confinement; wherein, the laser energy is deposited in the outer layers of a target microsphere with the resultant ablation imploding the inner layers to sufficiently high densities and temperatures for substantial thermonuclear burn of the D-T fuel. Termination of the burn occurs when internal pressures generated by the extreme temperatures overcome the inertial forces of the fuel with resultant rapid expansion and cooling of the fuel. This disassembly time is just the transit time for a rarefaction wave to propagate across the pellet radius. With this inertial confinement limiting the burn time, a breakeven condition for net energy production is $\rho R > 0.2 \text{ g/cm}^2$.³ In addition to this necessary compression criterion, there is the requirement common to all CTR schemes of reaching ignition temperature. The ignition temperature is typically defined as that temperature necessary for thermonuclear α particle kinetic energy to exceed radiation energy losses. For D-T fuel this temperature is between 3 and 7 keV depending on the density.⁴

Various pellet designs with differing pusher-tamper-fuel construction can optimize the coupling of the available laser energy into the necessary compression and heating, thus reducing constraints

on laser energy and pulse shape.⁵ Beyond these target optimizations, there is still the necessity of preventing extraneous decoupling mechanisms. The laser energy must pass through an underdense plasma layer to the absorption region (before or near the critical density). The absorbed energy must then be conducted to the ablation region. Any superthermal electrons produced in the absorption process could penetrate the pellet core, preheating the target, and preventing the required compression.⁶ Strong magnetic fields could be produced in the laser-plasma interaction reducing significantly the conduction of energy to the ablation front.³ Anomalous backscatter in the underdense region could prevent the laser beam from ever propagating to the absorption region. It is this last concern that has motivated theoretical and experimental investigations into backscatter coupling.⁷⁻¹⁰

2. Laser Heated Solenoids

The second class of laser-plasma devices for CTR is the laser heated solenoid (LHS) which employs magnetic confinement of the thermonuclear plasma. Originally proposed by Dawson, et.al., the LHS concept incorporates a solenoidal plasma which is axially heated by a long pulse, long wavelength laser.¹¹ Relying on classical inverse bremsstrahlung as the heating mechanism, the CO₂ laser is well matched to a plasma with density $\sim 10^{18} \text{ cm}^{-3}$ and length of several hundred meters. The LHS, as with other magnetically confined thermonuclear plasmas, is governed by the "Lawson Criterion" for breakeven energy production, $n\tau > 10^{14} \text{ s/cm}^3$.¹² Assuming 33% efficiency in thermonuclear energy recovery, the $n\tau$ lower limit is

set by requiring the recovered energy to be greater than the input energy. This criterion presumes operation at ignition temperature with complete escape of the reaction products. The burn time, dictated by the unstable characteristics of the plasma confinement, depends on the particular magnetically confined system under consideration.

The competitive advantages of the LHS system rest on its similarity to other linear reactor geometries: neutral plasma stability and high power density for efficient use of input energy. Furthermore, axial heating avoids the first wall problems of implosion-heating and allows high field operation through the use of a DC superconductor outside the reactor blanket. In addition to recognized problems of linear systems (axial confinement, magnetic field diffusion, heat conduction, etc.¹³) the LHS must also address the possible difficulties of laser heating: refractive containment, bleaching and anomalous backscatter. Again, anomalous backscatter from an underdense plasma is of concern as a possible decoupling mechanism.

The present experiment investigates laser interaction with a theta-pinch plasma. The geometry and parameters of this experiment are quite similar to LHS studies. Thus, the results of this investigation are particularly relevant to LHS development. Though the plasma conditions are quite different from LIP plasmas, it was thought that a study of the fundamental aspects of anomalous scattering would also be of value to LIP studies. While this work was in progress, LIP research indicated that the dominant mechanisms

apparently occur near the critical density point: extreme density steepening due to the ponderomotive force and resonant absorption.³ The large density gradients presumably prevent significant anomalous scatter and concern about this mechanism diminished. However, recent work with solid planar targets indicates significant anomalous scatter may occur with shaped laser pulses.¹⁴ Shaped laser pulses are anticipated in future LIP experiments. Hence, detailed understanding of anomalous scatter may still be of practical significance to LIP research.

B. Laser-Plasma Interaction

Theoretical work on the laser-plasma interaction has been varied and extensive.¹⁵⁻²⁰ To examine the different coupling mechanisms in a coherent manner, it is first useful to review the response of a single particle to a high frequency electromagnetic field.

1. Linear or Classical Effects

The laser-plasma interaction may be divided into classical (linear) effects and higher order, anomalous (nonlinear) effects. Following Schmidt's use of the Oscillation-Center Approximation,²¹ one begins with the equation of motion for an electron:

$$m_e \dot{\underline{v}} = -e[\underline{E}(\underline{r}, t) + (\underline{v}/c) \times \underline{B}(\underline{r}, t)]. \quad (\text{II-1})$$

For harmonic time variation, the $\underline{E}$ field is

$$\underline{E}(\underline{r}, t) = \underline{E}(\underline{r}) \cos \omega_0 t \quad (\text{II-2})$$

and the magnetic field is given by Maxwell's equation as

$$\underline{B}(\underline{r}, t) = -(c/\omega_0) \nabla \times \underline{E}(\underline{r}) \sin \omega_0 t. \quad (\text{II-3})$$

If we at first neglect field inhomogeneities and the magnetic field term, then (II-1) reduces to

$$\dot{\underline{y}}(0) = -(e/m) \underline{E}_0 \cos \omega_0 t. \quad (\text{II-4})$$

This has the straightforward solution of

$$\underline{y}(0) = -(e/m\omega_0) \underline{E}_0 \sin \omega_0 t \quad (\text{II-5})$$

which represents the 0th order response of a plasma to a laser; ie, the electrons are harmonically jiggled with a maximum velocity or "jitter" velocity given by

$$v_0 = eE_0/m\omega_0. \quad (\text{II-6})$$

The AC resistance of a plasma is due to binary collisions between electrons and ions driven in the above manner. This AC resistance represents the classical absorption process for a laser in a plasma and is typically referred to as inverse bremsstrahlung absorption (IBA). The work of Dawson and others^{22,23} has led to the high frequency, AC power absorption coefficient:

$$\kappa = \frac{16\pi Z^2 n_e n_i e^6 \ln \Lambda_0}{3c \nu_0^2 (2\pi m_e T)^{3/2} (1 - \nu_{pe}^2 / \nu_0^2)^{1/2}} \text{ cm}^{-1}, \quad (\text{II-7})$$

or an absorption length for $1/e$ attenuation of

$$l_{\text{abs}} = 1/\alpha \approx \frac{1.03 \times 10^{35} T^{3/2}}{n^2 \ln \lambda_0} \text{ cm} \quad (\text{II-8a})$$

$$\approx 2 \times 10^{34} T^{3/2} / n^2 \text{ cm}, \quad (\text{II-8b})$$

where $\omega_0 = 1.78 \times 10^{14}$ rad/sec for the CO_2 laser, ω_{pe} is the plasma frequency, $\omega_0 \gg \omega_{pe}$ and $\ln \lambda_0 \approx 5$ ($\lambda_0 = v_e / \omega_0 p_{\min}$, where v_e is the electron thermal velocity and $p_{\min}$ is the minimum impact parameter for electron-ion collisions).

In analogy to classical absorption, classical scattering is just the immediate re-emission of radiation from particles accelerated by an incident em field. This is generally referred to as Thomson scattering (or Compton scattering at higher frequencies). The instantaneous power radiated in the unit solid angle centered on $\hat{n}$ is given by Larmor's formula²⁴:

$$\frac{dP}{d\Omega} = \frac{e^2}{4\pi c} |\hat{n} \times (\hat{n} \times \dot{\underline{y}}/c)|^2 \quad (\text{II-9a})$$

$$= \frac{e^2}{4\pi c^3} |\dot{\underline{y}}|^2 \sin^2 \theta \quad (\text{II-9b})$$

where θ is the angle between $\dot{\underline{y}}$ and $\hat{n}$. Substitution of (II-5) into (II-9b) and averaging over time gives

$$\begin{aligned} \frac{dP}{d\Omega} &= |\underline{S}_0| r_0^2 \sin^2 \theta \\ &= |\underline{S}_0| \sigma_T(\theta) \end{aligned} \quad (\text{II-10})$$

where r_0 is the classical electron radius, $|\underline{S}_0|$ is the magnitude of the incident wave's Poynting vector

$$|\underline{S}_0| = \frac{c}{8\pi} |\underline{E}_0|^2$$

and $\sigma_T(\theta)$ is the Thomson differential scattering cross section.

In a plasma, local scattering occurs in a similar fashion, but net scattering only occurs in the presence of density fluctuations.²⁵ A volume of plasma radiating with a certain phase will always have a separate, complementary volume of plasma radiating out of phase with it. With a uniform density distribution, these two volumes would have scattered field component intensities that exactly cancel. With density fluctuations, the resulting time averaged scattered power in the frequency range $\omega \rightarrow \omega + d\omega$ in the solid angle $d\Omega$ centered on $\hat{n}$ is

$$P(\hat{n}, \omega) = \frac{|\underline{S}_0|}{2\pi} \sigma_T(\theta) n V S(\underline{k}', \omega') d\Omega d\omega \quad (\text{II-11})$$

where $\omega' = \omega_0 - \omega$, $\underline{k}' = \underline{k}_0 - \underline{k}$ and $S(\underline{k}', \omega')$ is the spectral density function giving the ensemble averaged density fluctuations:²⁶

$$S(\underline{k}', \omega') = \lim_{t \rightarrow \infty} \frac{\lim_{V \rightarrow \infty}}{tV} \left\langle \frac{|n_e(\underline{k}', \omega')|^2}{n_e} \right\rangle. \quad (\text{II-12})$$

It should be noted that most Thomson scattering arrangements give very low scattered intensity levels into typical detection geometries (of the order 10^{-12} x incident). Enhanced classical scattering can occur at frequencies corresponding to normal modes of the plasma. The anomalous scatter levels observed in this experiment typically were 6 to 7 orders of magnitude greater than the Thomson scatter levels. Thomson scattering will be assumed as the noise source for stimulated Brillouin backscatter in Ch. 3.

The classical effect of refractive bending of a laser beam in a plasma density gradient can also be seen as a 0th order result of (II-4).²⁷ The induced current in a plasma is

$$\mathbf{j} = eZn_i\mathbf{v}_i - en_e\mathbf{v}_e$$

so that (II-5) gives

$$\begin{aligned}\mathbf{j} &= \frac{e^2}{\omega_o} \left(\frac{n_i Z}{m_i} + \frac{n_e}{m_e} \right) \mathbf{E}_o \sin \omega_o t \\ &= \frac{n_e e^2}{m_e \omega_o} \mathbf{E}_o \sin \omega_o t.\end{aligned}\tag{II-13}$$

By inserting (II-13) and the expression for $\mathbf{E}$ into the field wave equation (from Maxwell's equations):

$$\nabla^2 \mathbf{E}(t) - \frac{4\pi}{c^2} \mathbf{j} - \frac{1}{c^2} \ddot{\mathbf{E}}(t) = 0,\tag{II-14}$$

and Fourier analyzing the resulting wave equation, one obtains the normal plasma dispersion relation for transverse waves:

$$\omega^2 = \omega_{pe}^2 + k^2 c^2. \quad (\text{II-15})$$

From this expression one can subsequently obtain the index of refraction for a plasma:

$$\begin{aligned} \mu^2 &= \frac{k^2 c^2}{\omega^2} = \frac{\omega^2 - \omega_{pe}^2}{\omega^2} \\ &= 1 - \omega_{pe}^2 / \omega^2. \end{aligned} \quad (\text{II-16})$$

Since $\omega_{pe}^2 \propto n_e$, macroscopic density fluctuations may then lead to ray bending.

2. Nonlinear or Anomalous Effects

The preceding discussion has shown the origin of IBA heating, Thomson scattering and refractive scattering to be classical or "0th" order laser-plasma interactions. To find the higher order interactions we return to (II-1) and solve for $\underline{y}$, allowing for the wave magnetic field and slight nonuniformity in the electric field. The inclusion of these next order terms is particularly important at high field strengths.

Substituting the Taylor expansion of $\underline{E}(\underline{r})$:

$$\underline{E}(\underline{r}) = \underline{E}(\underline{r}_0) + (\delta \underline{r} \cdot \underline{\nabla}) \underline{E}(\underline{r}_0) + \dots$$

into (II-1) and keeping 1st order terms, we may write the next higher order or nonlinear force equation as

$$m\dot{\underline{v}}^{(1)} = -e[(\delta\underline{r} \cdot \nabla)\underline{E}(\underline{r}_0) + (\underline{v}^{(0)}/c) \times \underline{B}(\underline{r}_0)], \quad (\text{II-17})$$

where $\underline{v} = \underline{v}^{(0)} + \underline{v}^{(1)}$ and

$$\delta\underline{r} = \int \underline{v}^{(0)} dt = (e/m\omega_0^2) \underline{E}_0 \cos \omega_0 t. \quad (\text{II-18})$$

Using the expressions for $\delta\underline{r}$, $\underline{v}^{(0)}$ and $\underline{E}$ in (II-17) and averaging over a cycle, one finds the secular, nonlinear force acting on the electron:²⁸

$$\begin{aligned} \underline{F}_{NL} &= m \langle \dot{\underline{v}}^{(1)} \rangle \\ &= - \frac{e^2}{2m\omega_0^2} [(\underline{E}_0 \cdot \nabla)\underline{E} + \underline{E}_0 \times (\nabla \times \underline{E})]. \end{aligned} \quad (\text{II-19})$$

The first term is due to finite excursions of the oscillating electron into regions of different field intensity and drives the electric parametric instabilities. Some of these instabilities represent anomalous absorption mechanisms since the decay products are electrostatic waves which are contained in the plasma. The second term gives rise to a drift of the electrons in the direction of the wave. This is due to the magnetic field distortion of the electron orbit.²⁹ It is this term that drives the parametric backscatter instabilities (stimulated Brillouin scatter and

stimulated Raman scatter) which represent the anomalous scattering processes.²⁸

Equation (II-19) may be reduced with a vector identity to give the nonlinear force per unit volume:

$$\begin{aligned} \underline{f}_{NL} = n_e \underline{F}_{NL} &= - \frac{1}{\omega_0^2} \frac{4\pi n_e e^2}{m} \nabla \cdot \left(\frac{\underline{E}^2(\underline{r})}{16\pi} \right) \\ &= - \frac{\omega_p^2}{\omega_0^2} \frac{\nabla \cdot \langle \underline{E}^2 \rangle}{8\pi}, \end{aligned} \quad (\text{II-20})$$

which can be recognized as the ponderomotive force.³⁰ Though the electrons exhibit greater response to the high frequency driving field (due to their lower mass), the resulting, secular ponderomotive force is transmitted to the ions through self-consistent fields.

As described above, this nonlinear force drives many of the parametric or anomalous coupling mechanisms. A more detailed description of this coupling process for stimulated Brillouin scatter will be presented later. In the next section, reference to stimulated scatter growth rates and thresholds will be made. To obtain these expressions one must use a complete set of fluid⁸ or kinetic⁷ equations which contain the ponderomotive force. It is sufficient now to understand the physical origin of the scatter mechanisms and to recognize their relative coupling strength (of the order v_0/c , see (II-17)). With the increased intensity available in high power pulsed lasers, v_0/c may rise to a level where stimulated

backscatter and other nonlinear effects begin to dominate the interaction.

3. Interaction Summary

To summarize the various laser-plasma interactions, a list of coupling mechanisms is now made. Classical and anomalous mechanisms will be described under the general headings of absorption and scattering. The pertinent rates for classical effects and the thresholds and growth rates for anomalous effects will be given. Where appropriate, quantitative estimates of these parameters will be made using the tabulated experimental parameters from the experimental section of this chapter (p. 28, fill pressure = 250 mtorr).

Classical Absorption

Inverse Bremsstrahlung Absorption (IBA)-

This is the high frequency, AC power absorption mechanism in an underdense plasma due to binary collisions between accelerated electrons and ions. For the CO₂ laser frequency the absorption length is

$$\begin{aligned} \ell_{\text{abs}} &\approx 2 \times 10^{34} T_e^{3/2} / n_e^2 \\ &\approx 50 \text{ cm.} \end{aligned}$$

Fig. 1 gives plotted values of ℓ_{abs} (after Dawson, et.al.³¹). This is the dominant absorption process in the present work.

Figure 1. l_{abs} for IBA as a function of n_e and T_e
 (after Dawson, et.al.;³¹ eq. II-7 is used).

Inverse Cyclotron Absorption (ICA)-

This process requires proper phasing between the incident wave and thermal cyclotron motion and results in a net increase of transverse kinetic energy through increased cyclotron motion.³² The absorption is strongest for propagation nearly perpendicular to $\hat{B}$.³¹ Propagation parallel to $\hat{B}$ and low temperature, in the present experiment, make this process negligible.

Resonant Absorption-

This is the limiting case of collisional absorption for $\omega_0 \rightarrow \omega_{pe}$. It requires a density gradient and the wave polarized with a finite component of $\underline{E}$ along the gradient for significant absorption.³³⁻³⁶ It is considered the primary absorption mechanism in LIP systems³ but is considered unimportant in the present underdense system.

Anomalous Absorption

Oscillating Two Stream Instability^{37,38} (OTS)-

Following Chen's discussion,²⁸ this is the decay of an incident em wave into stationary ion density striations transverse to the incident direction of propagation and a standing electron plasma wave. The ponderomotive force due to the combined electrostatic (es) and electromagnetic fields drive the stationary density fluctuations. The proper phase relationship for this to occur requires $\omega_0 \leq \omega_k$ where $\omega_k^2 = \omega_{pe}^2 + 3k^2 v_e^2$. The time averaged velocity distribution of the electrons oscillating in these density fluctuations is doubly

peaked; hence the name, oscillating two stream. For the present work $\omega_0 \gg \omega_k$ and this process is neglected.

Parametric Decay Instability^{39,40} (PDI)-

For $\omega_0 \gtrsim \omega_k$, displacement of the electron perturbation considered in the OTS case, is out of phase with the driving field and serves to damp the ion density fluctuations. However, if the density perturbations constitute an ion-acoustic wave, then in the moving frame of the wave the effective pump frequency, ω'_0 , is again less than ω_k . Thus, the OTS mechanism is possible and the waves grow. This represents a decay into an electron plasma wave and an ion-acoustic wave. Again, this only occurs well above the present density regime.

Dual plasma Wave^{41,42} ($2\omega_p$)-

In this mechanism the incident wave decays into two electron plasma waves with the required phase matching condition of $\omega_0 \leq 2\omega_{pe}$ or $n_e \geq n_c/4$, where n_c is the critical density. This again is outside the present parameter regime.

Classical Scattering

Thomson Scattering-

This is the immediate re-emission of radiation from accelerated particles in an em field. The scattered power is (see (II-11) and (II-12)):

$$P(\hat{n}, \omega) = \frac{|S_0|}{2\pi} \sigma_T(\theta) n V S(\underline{k}', \omega') d\Omega d\omega.$$

Thomson scattering allows for resonant scattering off normal modes of the plasma; ie, electron plasma waves and ion-acoustic waves. In addition, coherent scattering could also occur off drift waves, Bernstein waves and other electrostatic perturbations. In the present work, only Thomson scattering is considered and only as the thermal noise source for stimulated Brillouin scatter.

Anomalous Scattering

Stimulated Brillouin Scatter (SBS)-

This is a parametric decay with one product being an electromagnetic wave and the other being an ion-acoustic wave. The ponderomotive force resulting from the beating of the incident and the scattered wave drives the electrostatic perturbation, which provides the scattering source for the scattered em wave. The infinite homogeneous growth rate, γ_0 , the convective threshold, P_{th} , and the length for convective gain, ℓ_{conv} , have been given by several authors: 7, 8, 16

$$\gamma_0 = \frac{v_0}{c} \left(\frac{\omega_0}{2\omega_{ia}} \right)^{1/2} \omega_{pi}$$

$$\approx 10^{10} \text{ sec}^{-1}$$

$$P_{th} > 2.7 \times 10^{-5} \frac{\gamma_p \gamma_{em}}{\omega_{pi}^2} \frac{\omega_{ia}}{\omega_0} c$$

$$\approx 5 \times 10^8 \text{ w/cm}^2$$

$$l_{\text{conv}} = \frac{\gamma_{\text{pc}}}{2\gamma_0^2} \approx 1 \text{ cm.}$$

The threshold and gain length were significantly exceeded by the CO_2 laser power capability of $\sim 3.5 \times 10^{10} \text{ w/cm}^2$ and focus region of $\sim 50 \text{ cm}$. An extensive discussion of this, the most likely parametric decay, constitutes the main effort of this thesis.

Related ion mode scattering-

Two limiting cases of SBS have been given additional names. "Resistive quasi-ion mode scattering" (or "stimulated Compton scattering", "induced Compton scattering", "nonlinear Landau growth") refers to the non-collective limit of $T_e \lesssim T_i$ (ion-acoustic mode heavily damped) or $k\lambda_D \gtrsim 1$. "Reactive quasi-ion mode scattering" (or "strong coupling limit") refers to the case where the pump wave is so strong as to drive the ion waves off resonance ($\omega_{\text{es}} > kc_s$).⁴²

It might also be mentioned that the cutoff for SBS occurs at $\gtrsim \omega_{\text{pe}}$ or $\lesssim n_c$ at which point the scattered wave would see a critical density and the PDI mechanism would be expected to occur. However, both these mechanisms may be dominated at the critical density by the ponderomotive force from the reflection of the incident wave (causing extreme density gradients)⁴⁴ and by resonant absorption.

Stimulated Raman Scatter (SRS)-

This parametric decay has the same driving mechanism as SBS but with the ion-acoustic wave replaced by an electron plasma wave. The analogous growth rate, threshold and gain length are:^{7,8,16}

$$\gamma_0 = \frac{v_0}{c} \left(\frac{\omega_0 \omega_{pe}}{2} \right)^{1/2} \approx 6 \times 10^{10} \text{ sec}^{-1}$$

$$P_{th} > 2.7 \times 10^{-5} c^2 \frac{\gamma_p \gamma_{em}}{\omega_{pe} \omega_0} \approx 1 \times 10^8 \text{ w/cm}^2$$

$$l_{conv} = \frac{\gamma_p c}{2\gamma_0^2} \approx 0.2 \text{ cm.}$$

SRS appears as likely to occur as SBS. The phase mismatching due to axial density gradients is estimated in section III.A.7. It is found that density gradients are likely to prohibit significant SRS gain. SRS did not appear in low sensitivity scans, integrated wavelength monitor traces or in calorimeter measurements. SRS is therefore considered an unimportant mechanism in the present work.

Related electron mode scattering-

In analogy to SBS, the non-collective and over-driven limits of SRS are resistive quasi-electron mode scattering and reactive quasi-electron mode scattering, respectively. The cutoff for SRS occurs at $\geq 2\omega_{pe}$ or $\leq n_c/4$ at which point the scattered wave would see a critical density and the $2\omega_p$ instability would be expected to occur.⁴⁵

Modulational Instability (filamentation, self-focusing)-

Refraction of the incident wave by macroscopic density striations can channel the em wave into the less dense regions. The resulting ponderomotive force, due to the modulated field intensity, increases these density perturbations. The ponderomotive force must overcome a small electric field due to a finite Debye length which sets the instability threshold:⁷

$$\frac{v_o^2}{c^2} > 4(1+T_i/T_e)k_{es}^2\lambda_D^2, \text{ or}$$

$$P_{th} > 1 \times 10^{-4} k_{es}^2 \lambda_D^2 c^2 \approx 4 \times 10^8 \text{ w/cm}^2,$$

where $k_{es} = r^{-1}$ and r is the focused beam radius. The growth rate is

$$\gamma \approx \frac{1}{2} \frac{v_o}{c} \omega_{pi} \approx 2.7 \times 10^8 \text{ sec}^{-1}.$$

This is a likely process for the present experiment. Evidence for filamentation and channeling (in the programmed, on-axis density minimum) is discussed in Ch. 5. (It is also possible to view filamentation as a four-wave process, in a similar manner to the OTS instability.)

In fig. 2 the various laser-plasma interactions are outlined according to frequency regimes. In Table I, a list of thresholds and growth rates for the parametric mechanisms are given (after Liu and Kaw⁴⁶). The most likely occurring anomalous instabilities are SBS,

Figure 2. Comparison of coupling mechanisms as a function of frequency regimes. The experimental regime is far to the right ($\omega_o \gg \omega_{pe}$).

TABLE I
THRESHOLDS AND GROWTH RATES OF PARAMETRIC INSTABILITIES

<u>Instability</u>	<u>$P_T \left(P = \frac{m^2 \omega_o^2 c}{4 \pi e^2} v_o^2 \right)$</u>	<u>γ_o</u>
SBS	$\frac{v_o^2}{c^2} > \frac{\gamma_p \gamma_{em}}{\omega_{pi}^2} \frac{\omega_{ia}}{\omega_o}$	$\frac{v_o}{c} \left(\frac{\omega_o}{\omega_{ia}} \right)^{1/2} \omega_{pi}$
SRS	$\frac{v_o^2}{c^2} > \frac{\gamma_p \gamma_{em}}{\omega_o \omega_{pe}}$	$\frac{v_o}{c} (\omega_o \omega_{pe})^{1/2}$
PDI	$\frac{v_o^2}{v_e^2} > \frac{4 \gamma_{pi} \gamma_{pe}}{k c_s \omega_{pe}}$	$\frac{1}{2} \frac{v_o}{v_e} (k v_e)^{1/2} \omega_{pi}^{1/2}$
OTS	$\frac{v_o^2}{v_e^2} > 2 \left(1 + \frac{T_e}{T_i} \right) \frac{\gamma_p}{\omega_p}$	$\frac{k v_o}{\left(1 + \frac{T_e}{T_i} \right)^{1/2}}$
$2\omega_p$	$\frac{v_o}{c} > \frac{2 \gamma_p}{\omega_p}$	$\frac{v_o}{2c} \omega_p$

SRS and filamentation. Indirect evidence of filamentation was found and is discussed in a later chapter. It was beyond the capability of this experiment to make a detailed search for SRS. However, coarse searches for SRS indicated it to be an unimportant loss process. Furthermore, it is considered less likely than SBS as a scattering process for reasons mentioned above. SBS appears as the strongest scattering process and the most dangerous anomalous loss mechanism. Therefore, the central theme of this thesis is the observation and study of SBS.

C. Experiment Description

1. Machine Characteristics

The theta-pinch and CO_2 laser characteristics have been previously reported.^{2,47} An abbreviated description of the experimental apparatus is given here.

The linear Scylla I-C theta-pinch had a maximum energy storage of 175 kJ at 60 kV main bank voltage. Fifty-four 1.8 μf capacitors energized the 100 cm long, 10.5 cm diameter, single turn compression coil. The present set of experiments was conducted at 40 kV main bank voltage. A vacuum E_0 of .26 kV/cm was generated at the inner wall of the 3.6 cm i.d. quartz discharge tube with a peak crow-barred compression field of 33 kG at $\sim 1.8 \mu\text{s}$ after discharge initiation. The fill pressure ranged from 100 to 500 mtorr D_2 . Refraction of the axially directed CO_2 laser beam out of the plasma was observed when the plasma column had an on-axis density maximum.⁴⁷ A forward bias magnetic field of ~ 1.8 kG was employed which upon compression during main bank discharge resulted in an

on-axis density minimum. This density minimum refractively contained the laser beam in the plasma column.

The cold-cathode, gain switched, CO_2 laser oscillator used in the experiment was a modified A-3 amplifier from LASL L Division.⁴⁸ The active laser volume was 1 m in length with an 8 cm diameter. A confocal unstable resonator cavity was used. Optimum laser operation was obtained with a 3:1/4:1::He:N₂:CO₂ gas mixture and a pumping chamber stress level of $S = 4.3 \text{ kV/cm-atm}$ (63kV, 1400 torr). The laser had a multimode output in both transverse and longitudinal modes. The emission spectrum was dominated by the P-20 line. The measured small signal gain (R-16 line) was $.041 \text{ cm}^{-1}$ for an active length of 100 cm. The laser reliably produced $\sim 150 \text{ J}$ with 40-50% of the energy in the first 60 ns FWHM spike and the remainder in a $1.5 \mu\text{s}$ tail. The 8 cm diameter beam was axially focused to 2 mm diameter (focal length $\sim 3.2 \text{ m}$) at 10 cm into the meter-long, theta-pinch deuterium discharge. A decrease in longevity of the programmed density minimum at higher fill pressures restricted the CO_2 laser timing to a nominal $1.2 \mu\text{s}$ after main field initiation.

2. Diagnostic Arrangement

A variety of diagnostic systems were developed and employed to investigate the independent plasma and beam parameters as well as to study the laser-plasma interaction. A schematic layout of the experimental and diagnostic arrangement is shown in fig. 3. The laser parameters were monitored with:

- 1) calorimeter- incident energy
- 2) photon drag detector- incident pulse shape

Figure 3. Schematic diagram of experimental and diagnostic arrangement.

- 3) polychromometer with pyroelectric vidicon detector-
incident spectrum.

The plasma parameters were studied with:

- 1) ruby laser holographic interferometer- electron density
- 2) two monochromators for plasma spectroscopy-
electron temperature.

In addition to these plasma diagnostics, extra equipment was used to further study the interaction:

- 1) calorimeter- transmitted beam energy
- 2) monochromator with photovoltaic detectors- backscattered
beam spectrum.

Other diagnostic studies included:

- 1) exit beam burn patterns
- 2) streak photography
- 3) excluded flux measurement
- 4) reflected beam energy calorimetry.

The detailed diagnostic arrangements, with corresponding data and analysis, will be presented in the chapters concerned with the respective diagnostic measurements. For convenience and reference a listing of experimental and calculated parameter values is given in Table II. The parameters are given for 1.2 μ s time delay and for fill pressures of 100, 250 and 400 mtorr. Temperature measurements were made at two axial locations: 14 and 43 cm into the coil. Parameters at both axial locations are given.

TABLE II

ON-AXIS EXPERIMENTAL PARAMETERS FOR $z = 14$ AND 43 cm

<u>Parameters</u>	<u>100 mtorr</u>	<u>250 mtorr</u>	<u>400 mtorr</u>
(on-axis, $t = 1.2 \mu s$)	$z = 14./43.$	$14./43.$	$14./43.$
n_e (10^{17} cm^{-3})	.63	1.4	2.1
T_e (eV)	22./32.	13./17.	10./10.
T_e ($t = 1.4 \mu s$) (eV)	—	34./28.	16./12.
nkT ($J \cdot \text{cm}^{-3}$)	.22/.32	.29/.38	.34/.34
$\bar{E} = nkTV$ (J)	21.	27.	27.
B_0 (kG)		31.5	
$B = (B_0^2 - 8\pi nkT)^{1/2}$ (kG)	30.6/30.2	30.3/29.9	30.1/30.1
$\beta = 8\pi nkT/B_0^2$	.06/.08	.07/.10	.09/.09
$v_e = (T_e/m_e)^{1/2}$ (10^8 cm/s)	2.0/2.4	1.5/1.7	1.3/1.3
$v_i = (T_i/m_i)^{1/2}$ (10^6 cm/s)	3.2/3.9	2.5/2.8	2.2/2.2
$v_A = B/(4\pi n_i m_i)^{1/2}$ (10^7 cm/s)	5.9	1.2	1.0
$\ln \Lambda_p = \ln(1.5 \times 10^{10} T_e^{3/2} n_e^{-1/2})$	8.7/9.3	7.5/7.9	6.9/6.9
$\tau_{ee} = 3.3 \times 10^5 T_e^{3/2} / n_e \ln \Lambda_p$ ($10^{-12} s$)	62./102.	15./21.	7.2/7.2
$\tau_{ii} = 1.4 \times 10^7 A^{1/2} T_i^{3/2} / Z^4 n_i \ln \Lambda_p$ ($10^{-9} s$)	3.7/6.1	.9/1.3	.4/.4
$\tau_{ei} = 3.5 \times 10^5 T_e^{3/2} / Z^2 n_i \ln \Lambda_p$ ($10^{-12} s$)	66./108.	16./22.	7.6/7.6
$\tau_{ei}^{eq} = 3.5 \times 10^5 T_e^{3/2} / Z^2 n_i \ln \Lambda_p$ ($10^{-9} s$)	241./395.	59./80.	28./28.
$\omega_e = eB/m_e c$ ($10^{11} s^{-1}$)	5.4/5.3	5.3/5.3	5.3/5.3
$\omega_i = eB/m_i c$ ($10^8 s^{-1}$)	1.5/1.4	1.5/1.4	1.4/1.4

TABLE II (continued)

λ	$= v_e \tau_{ei}$	(10^{-3}cm)	13./26.	2.4/3.7	1.0/1.0
r_e	$= v_e / \omega_e$	(10^{-4}cm)	3.7/4.5	2.8/3.2	2.5/2.5
r_i	$= v_i / \omega_i$	(10^{-2}cm)	2.1/2.8	1.7/2.0	1.6/1.6
$\omega_e \tau_{ei}$			36./57.	8.5/12.	4.0/4.0
$\omega_i \tau_{ii}$			.6/.9	.1/.2	.06/.06
ω_{pe}	$= (4\pi n e^2 / m_e)^{1/2}$	(10^{13}s^{-1})	1.4	2.1	2.6
ω_{pi}	$= (4\pi n e^2 / m_i)^{1/2}$	(10^{11}s^{-1})	2.3	3.5	4.3
λ_D	$= v_e / \omega_{pe}$	(10^{-5}cm)	1.4/1.7	.7/.8	.5/.5
λ_0	(μm)			10.5915	
k_0	(cm^{-1})			5.93x10 ³	
ω_0	(rad^{-1})			1.78x10 ¹⁴	
v_0	$= 2.56 \times 10^5 \lambda_0 P^{1/2}$	$(\text{cm/s}, P = 3 \times 10^{10} \text{ w/cm}^2)$		4.7x10 ⁷	
P	(w)			1x10 ⁹	
E	(J)			150.	
$\ln \Lambda_0$	$= \ln(3.0 \times 10^{14} T_e^{3/2} / \omega_0)$		5.2/5.7	4.4/4.7	3.9/3.9
ℓ_{abs}	$= 3.2 \times 10^6 \omega_0^2 T_e^{3/2} / Z n_e^2 \ln \Lambda_0$	(cm)	514./824.	56./78.	19./19.
γ_{em}	$= c/2 \ell_{\text{abs}}$	(10^8s^{-1})	.3/.2	2.7/1.9	7.9/7.9
k	$= 2k_0$	(cm^{-1})		1.2x10 ⁴	
$k \lambda_D$			.17/.20	.08/.09	.06/.06
c_s	$= (T_e / m_i)^{1/2}$	(10^6cm/s)	3.2/3.9	2.5/2.8	2.2/2.2
ω_{ia}	$= k c_s$	(10^{11}s^{-1})	3.8/4.7	3.0/3.4	2.6/2.6
v_p	$= 2.05 v_i$	(10^6cm/s)	6.6/8.0	5.1/5.7	4.5/4.5
ω_p	$= k v_p$	(10^{10}s^{-1})	7.9/9.6	6.2/6.9	5.4/5.4
γ_p	$= .85 k v_i$	(10^{10}s^{-1})	3.3/4.0	2.6/2.9	2.2/2.2

3. Summary of Experimental Results

For continuity, a brief sketch of the diagnostic methods and a summary of the experimental results are now given. The areas of study outlined in the introduction will be followed and discussed individually.

1) Anomalous Backscatter, SBS-

The backscatter diagnostic system used a HgCdTe photovoltaic detector situated at the exit slit of a 0.5 m scanning monochromator with a measured instrument resolution of 10° Å HWHM. The SBS backscatter signal consisted of a single, strong pulse occurring within ± 15 ns of the incident beam's peak power. Furthermore, the stimulated backscatter signal showed some erratic, shot-to-shot power fluctuation and a cutoff at an incident power of $\sim 1.2 \times 10^{10}$ W/cm².

The single pulse nature appears due to a time-varying, axial temperature gradient. An increased SBS threshold and gain length may also be present due to Faraday rotation of the incident and backscattered waves.

A spectral scan of the backscatter radiation, normalized to a wavelength integrated backscatter monitor, was compared to theory. The number of exponentiations of the backscattered signal from the noise level, $n_{ef} \geq 14$, required by the fit of the SBS spectrum, gave a lower limit for the interaction length of ≥ 13 cm. Consideration of the SBS power, 50 Mw, leads to an upper limit on the number of exponentiations, $n_{ef} \leq 24$ or an upper limit on the interaction length of ≤ 22 cm. The restricted interaction length of 13 - 22 cm is consistent with the focus region of ~ 50 cm.

SBS was shown to have a higher number of exponentiations than has previously been reported, which is consistent with the longer available interaction length of this experiment. Strong evidence of reduced SBS coupling, due to temperature inhomogeneity, is apparent in the prompt, single pulse time history and the fluctuation of backscattered power. The power reflectivity was found to be $\leq 5\%$. The energy reflectivity, however, was $< 0.1\%$ due to the short, single SBS pulse characteristic. Hence, SBS did not constitute a significant loss mechanism.

2) Propagation-

Ruby laser, end-on holographic interferometry provided a measurement of the radial electron density profile integrated over the length of the plasma column. This measurement was required to define the general plasma characteristics and to allow calculation of growth rates and thresholds pertinent to anomalous backscatter. Moreover, it provided a means to estimate the refraction of the CO_2 laser beam by the plasma column.

A simple linear interpolation of the radial density profile with an on-axis density minimum was made. A geometrical optics formulation for ray tracing then indicated that the laser beam was contained upon re-expansion past its focus point. Interferograms further showed that the density minimum had totally diffused away by $2.0 \mu\text{s}$ delay time. Thus, after $2.0 \mu\text{s}$ the laser beam is refracted out of the plasma column and lost. This refractive side loss represents $\sim 50\%$ of the total laser beam energy. These refractive

loss estimates are consistent with net absorption calculations and local heating measurements.

Exit beam burn patterns, ~ 2 m past the theta-pinch coil, provide additional qualitative information. At the lower fill pressures (100 to 250 mtorr), where less absorption occurs, the burn pattern consisted of a single, focused burn ~ 0.5 cm in diameter. This clearly indicates refractive trapping of the CO₂ laser beam by the density minimum. Within the burn spot, spatial modulation of the intensity pattern was often seen. The regular spacing of the intensity modulation suggests filamentation scattering. However, time restrictions did not allow a conclusive study of this burn feature.

Significant plasma expansion due to laser heating is not expected with a low β plasma ($\beta < 0.1$). Study of interferograms taken at 1.8 and 3.0 μ s delay time shows no expansion.

3) Absorption-

Calorimetry was employed to monitor and normalize the laser pulse. A measurement of the net absorption of the plasma column was obtained through the use of incident, reflected and transmitted energy calorimeters. It was found that a straightforward energy inventory using these measurements was inadequate. A complete accounting of the incident energy was possible only when allowance was made for refractive side losses and bleached transmission (time dependent absorptivity). The analysis for 250 mtorr fill pressure resulted in 46% absorption of the laser energy for the portion of the beam contained in the plasma column (refractive side losses

subtracted from incident energy) and 15% absorption of the total incident energy. This examination made clear that a major limitation on laser heating in this experiment was refractive loss of the beam over the particularly long geometry of the plasma column.

4) Heating-

Local electron temperature variations in the plasma column were studied through spectral line emission measurements. A double monochromator, observation system was employed at two axial locations: 14 and 43 cm into the theta-pinch coil. Carbon and nitrogen impurity ions were studied. In order to stabilize the rather low, fluctuating, naturally occurring impurity level, a small amount (0.5%) of CH_4 or N_2 was introduced during the static, D_2 filling of the discharge tube.

The use of impurity line intensity ratios for T_e measurement was hampered by the lack of equilibrium of ground state populations, the complicated equilibrium dependence of excited levels within an ion stage and by radial temperature variations after laser firing. For these reasons, analysis of line intensities for T_e measurement was abandoned and a technique which depended on the dynamic nature of n_e , T_e and the ion populations was employed. The electron temperature was inferred from an ionization rate code which simulated the measured appearance time of line radiation intensity peaks. Subsequent stages of the carbon and nitrogen impurity ions were studied. At 14 cm into the coil, these measurements resulted in an electron temperature of 13 eV at the time of laser firing and 34 eV ~ 200 ns later; ie, just after the first laser spike (250 mtorr fill

pressure). In the absence of laser firing the temperature increased from 13 eV to 14 eV over this time period. At 43 cm into the coil, with laser firing, the temperature increased from 17 eV to 28 eV during this time period.

To compare these temperatures with theoretical predictions, it is necessary to study in detail the absorption process and energy loss mechanisms. If only simple IBA was applicable ($\ell_{\text{abs}} \sim T_e(0)^{3/2}/n_e^2$), the resultant electron temperatures would be in the range of many hundreds of eV. The effect of bleaching (a time dependent absorption length, $\tilde{\ell}_{\text{abs}} = \ell_{\text{abs}}[T_e(x,t)/T_e(0)]^{3/2}$) reduces significantly the absorptivity of the plasma. The anticipated temperatures then fall in the neighborhood of 90 eV. Further consideration of electron-ion equipartition, thermal conduction, radiation and propagation geometry leads to predicted electron temperatures which are consistent with the measured values. This comparison provides a strong verification of the bleaching phenomena.

The general conclusion of this experiment is that classical effects dominated the laser-plasma interaction. The observed SBS was shown to be consistent with time dependent temperature gradients. The energy reflectivity of SBS was negligibly small. Channeling was demonstrated. The on-axis density minimum did not persist for the entire duration of the laser pulse. Hence, refractive loss of the beam occurred which reduced the net absorption efficiency. Lastly, heating was shown to be well described by bleached absorption.

It is difficult to apply these results in detail to other experimental efforts due to differing plasma parameters, time scales, etc. The result of this study is to demonstrate general validity of theoretical descriptions and to point out various areas where pre-conceived concerns were justified or inadequate. With regard to LHS reactor schemes, the present experiment in particular indicates that: 1) totally contained propagation along the entire plasma column is especially difficult for a long aspect plasma ($L/r \gtrsim 100$) without a sustained, on-axis density minimum, 2) bleached absorption is consistent with theory and 3) anomalous backscatter losses do not present limitations to LHS projections.

III. ANOMALOUS BACKSCATTER - Stimulated Brillouin Scatter

A. Theory

1. Physical Description

To illustrate the mechanism of SBS, a physical description will first be given. In fig. 4, a decay of an incident, "pump", em wave is depicted. The daughter waves consist of an es, forward propagating wave and an em, backward propagating wave. Each of these daughter waves couples with the pump to drive the other daughter wave. This process of wave growth, through the presence of another, phase related, simultaneously growing wave, is often referred to as stimulated growth or stimulated scatter. The fact that the process is dependent on the presence of a pump wave, its parameters and its modulation of the plasma medium, allows the description "parametric" decay.

The coupling of the daughter waves, through the pump wave, to produce wave growth is further illustrated in fig. 4. Any "noise", density fluctuations, will be harmonically jiggled by the E field of the pump wave. The jiggling of density fluctuations leads to currents which act as a source for the scattered em wave. Thus, the es wave and pump couple to drive the scattered em wave. The other part of the coupling lies in the scattered em wave beating with the pump wave. The time averaged force due to these waves is spatially periodic such as to enhance the density perturbations of the es wave. A more detailed view of the beating of the two em waves shows the resultant force to arise from the interaction of the jitter velocities produced by one wave and the magnetic field of the other.

Figure 4. The physical mechanism of the SRS interaction.

Damping of both daughter waves occurs naturally. Hence, for net growth of these waves through the above, stimulated process, the growth rate must exceed the combined damping rates. This balance places a threshold on the required pump power, a distinct characteristic of stimulated scatter.

The phenomena of SBS has long received attention in the field of nonlinear optics⁴⁹ and more recently in the area of plasma physics. Thus, the theoretical descriptions available in the literature for linear and nonlinear growth, saturation and inhomogeneous mediums are extensive. In the following sections an endeavor is made to lay a foundation for the aspects of SBS most pertinent to the present experiment. The references provide a source for more formal and extensive theoretical treatments.

2. Coupled Equations

The work found in references 7, 8, 15, 16 and 18 represents the basis for theoretical interpretation of stimulated scatter in this experiment. A derivation, following Forslund's, et.al., fluid treatment, of the essential features of SBS is now presented. The standard two-fluid equations, combined with Maxwell's equations, lead to a generalized description of the excitation of a scattered em wave, ion wave and electron wave. From these coupled equations, a general dispersion relation is found for the SBS es wave.

The general analysis considered is the one-dimensional stability of a large amplitude light wave in an infinite homogeneous plasma. The fundamental equations, in the two-fluid formulation, are:

continuity-

$$\dot{n}_j + \nabla \cdot (n_j \underline{v}_j) = 0 \quad (\text{III-1})$$

momentum-

$$m_j n_j [\dot{\underline{v}}_j + (\underline{v}_j \cdot \nabla) \underline{v}_j] = q_j n_j [\underline{E} + (\underline{v}_j / c) \times \underline{B}] - \nabla P_j - m_j n_j \tau_{ei} \underline{v}_j \quad (\text{III-2})$$

equation of state-

$$P_j n_j^{-\delta} = \text{constant} \quad (\text{III-3})$$

Maxwell's equations-

$$\nabla \cdot \underline{E} = 4\pi \sum q_j n_j \quad \nabla \times \underline{E} = - \frac{1}{c} \dot{\underline{B}} \quad (\text{III-4a,4b})$$

$$\nabla \cdot \underline{B} = 0 \quad \nabla \times \underline{B} = \frac{4\pi}{c} \sum q_j n_j \underline{v}_j + \frac{1}{c} \dot{\underline{E}} \quad (\text{III-4c,4d})$$

The last term in the momentum equation is an approximate expression for the collisional transfer of momentum between species.⁵⁰ In eq. (III-3), δ is the ratio of specific heats. Two of Maxwell's equations are unnecessary since (III-4a) and (III-4c) can be derived from the divergences of (III-4d) and (III-4b), respectively. Thus, with the ideal gas law ($P_j = n_j T_j$), there are 18 scalar equations for 18 scalar unknowns: n_e , n_i , T_e , T_i , P_e , P_i , $\underline{v}_e$, $\underline{v}_i$, $\underline{E}$ and $\underline{B}$.

For convenience and conformity with the work of Forslund, et.al., these equations will be rearranged before a solution is sought. From (III-4b) and (III-4c), we may write $\underline{E}$ and $\underline{B}$ in terms of potential functions:

$$\underline{B} = \nabla \times \underline{A}$$

$$\underline{E} = -\underline{\nabla}\phi - \frac{1}{c} \dot{\underline{A}}.$$

Choosing the Coulomb or transverse gauge, $\underline{\nabla} \cdot \underline{A} = 0$, we may then replace the other Maxwell equations, (III-4a) and (III-4d), by

$$\nabla^2 \phi = -4\pi \sum q_j n_j \quad (\text{III-5})$$

$$\nabla^2 \underline{A} - \frac{1}{c^2} \ddot{\underline{A}} = -\frac{4\pi}{c} \sum n_j q_j \underline{v}_j + \frac{1}{c} \underline{\nabla} \dot{\phi} \quad (\text{III-6})$$

where the sources for the field potentials are on the right. Eq. (III-3) and the ideal gas law may be combined to yield a modified equation of state: $\underline{\nabla} p_j = \delta_j T_j \underline{\nabla} n_j$. Inserting this expression and the field potential expressions for $\underline{E}$ and $\underline{B}$ into (III-2) yields

$$\begin{aligned} \dot{\underline{v}}_j + (\underline{v}_j \cdot \underline{\nabla}) \underline{v}_j &= \frac{q_j}{m_j} \left(-\underline{\nabla} \phi - \frac{1}{c} \dot{\underline{A}} \right) + \frac{q_j}{m_j c} \underline{v}_j \times (\underline{\nabla} \times \underline{A}) \\ &- \frac{\delta_j T_j}{m_j n_j} \underline{\nabla} n_j - \nu_{ei} \underline{v}_j. \end{aligned} \quad (\text{III-7})$$

(III-1), (III-5), (III-6) and (III-7) constitute the model equations for describing stimulated scatter.

We now proceed to derive the three coupled wave equations characterizing the electromagnetic, ion and electron fluctuations. This will be accomplished using expansions of n_j and $\underline{v}_j$, keeping leading order terms. Wave propagation and spatial variations will be considered in only one dimension ($\hat{e}_z$ or $\hat{e}_\parallel$) to simplify the

derivation. A separation of the electromagnetic field potential into incident pump and scattered contributions will be introduced later.

The wave equation for the electromagnetic field potential is considered first. The choice of wave propagation in one direction in a homogeneous plasma allows us to write $\nabla = \hat{e}_z(\partial/\partial z) = \nabla_{||}$. Furthermore, since the Coulomb gauge was chosen one has

$$\nabla_{||}\dot{\phi} = 4\pi \sum_j q_j n_j v_{j||} \quad .51$$

Hence, the source in the wave equation for $\underline{A}$ can be expressed solely in terms of the perpendicular current; ie, $\nabla_{||}\phi = 0$ and $\underline{A} = \underline{A}_{\perp}$, or

$$\nabla^2 \underline{A}_{\perp} - \frac{1}{c^2} \ddot{\underline{A}}_{\perp} = - \frac{4\pi}{c} \sum_j n_j q_j \underline{v}_{j\perp} \quad (III-8)$$

This is equivalent to saying that the $\underline{E}$ and $\underline{B}$ fields of the em wave are functions only of $\underline{A}$, and the $\underline{E}$ fields associated with the electrostatic density fluctuations are functions only of ϕ . The separation of field potentials in the fluid formulation is similar to the separation of $v_{\perp}$ and $v_{||}$ motion in the single particle description of Sec. II-B. The transverse em wave couples directly to $v_{\perp}$. Whereas, the anomalous effects, contained in longitudinal fields are dependent on $v_{||}$.

Only $\underline{v}_{j\perp}$ needs to be specified for completion of the electromagnetic wave equation. Considering the $\perp$ component of eq. (III-7) and neglecting thermal and collisional terms, one readily finds to lowest order

$$\underline{v}_{j\perp} = - \frac{q_j \underline{A}_{\perp}}{m_j c} \quad (\text{III-9})$$

in analogy to eq. (II-6). The thermal term is negligible since variations in the plasma homogeneity are only allowed in the $\parallel$ direction.

The collisional term does not appear in the work of Forslund, et.al. Instead, a phenomenological damping term was introduced. Keeping the approximate collisional expression above, the damping term will appear as a natural consequence. The $\perp$ component of eq. (III-7) then yields

$$\dot{\underline{v}}_{e\perp} = \frac{e}{m_e c} \dot{\underline{A}}_{\perp} - \nu_{ei} \underline{v}_{e\perp}. \quad (\text{III-10})$$

Due to the large mass difference between electrons and ions, only $\underline{v}_{e\perp}$ need be kept for the current contribution in (III-8). (III-10) has the approximate solution:

$$\underline{v}_{e\perp} \approx \frac{e \underline{A}_{\perp}}{m_e c} + \frac{e}{m_e c} \frac{\nu_{ei}}{\omega_0^2} \dot{\underline{A}}_{\perp}, \quad (\text{III-11})$$

where the identification $\underline{A}(t) \sim \underline{A}_0 \sin(\omega_0 t)$ is made and it is assumed that damping is a small correction to (III-9); ie, $\nu_{ei}/\omega_0^2 \ll 1$ and eq. (III-11) satisfies (III-10) only to lowest order. Substitution of (III-11) into (III-8) then gives

$$\nabla^2 \underline{A}_\perp - \frac{1}{c^2} \ddot{\underline{A}}_\perp = \frac{4\pi e^2 n_e}{c^2 m_e} \underline{A}_\perp + \frac{4\pi e^2 n_e}{c^2 m_e} \frac{\nu_{ei}}{\omega_0^2} \dot{\underline{A}}_\perp.$$

Expanding about the mean density $n_e^{(0)}$, $n_e = n_e^{(0)} + n_e^{(1)}$, defining $\omega_{pe}^2 = 4\pi e^2 n_e^{(0)} / m_e$, and keeping leading order terms, one finds

$$\ddot{\underline{A}}_\perp + 2\gamma_{em}\dot{\underline{A}}_\perp - c^2\nabla^2 \underline{A}_\perp = -\omega_{pe}^2 \left(1 + \frac{n_e^{(1)}}{n_e^{(0)}}\right) \underline{A}_\perp, \quad (\text{III-12})$$

where $\gamma_{em} = \nu_{ei}\omega_{pe}^2/2\omega_0^2$. γ_{em} represents the ohmic dissipation of IBA, as seen in the approximate equality $\gamma_{em} \approx c/2\ell_{abs}$. Eq. (III-12) is the desired wave equation describing the propagation of em waves.

The wave equations for density fluctuations are obtained from eq. (III-7). We begin with the expansion of n_j as before. In addition, $\underline{v}_j$ is expanded according to $\underline{v}_j = \underline{v}_j^{(0)} + \underline{v}_j^{(1)}$, where $\underline{v}_j^{(0)} = \underline{v}_{j\perp}$ and $\underline{v}_j^{(1)} = \underline{v}_{j\parallel}$. With these expressions substituted into (III-7), the resulting first order equation is

$$\begin{aligned} \dot{\underline{v}}_{j\parallel}^{(1)} = & -\frac{q_j}{m_j} \nabla_\parallel \phi + \frac{q_j}{m_j c} \underline{v}_{j\perp}^{(0)} \times (\nabla_\parallel \times \underline{A}_\perp) \\ & - \frac{\delta_j T_j}{m_j n_j^{(0)}} \nabla_\parallel n_j^{(1)} - \nu_{ei} \underline{v}_{j\parallel}^{(1)}. \end{aligned}$$

With the lowest order expression for $\underline{v}_{j\perp}$, (III-9), and the vector identity $\underline{A}_\perp \times (\nabla_\parallel \times \underline{A}_\perp) = (1/2)\nabla_\parallel (\underline{A}_\perp \cdot \underline{A}_\perp)$, this becomes

$$\begin{aligned}
\dot{\underline{v}}_{j\parallel}^{(1)} &= \frac{-q_j}{m_j} \underline{\nabla}_{\parallel} - \frac{q_j^2}{2m_j^2 c^2} \underline{\nabla}_{\parallel} (\underline{A}_{\perp} \cdot \underline{A}_{\perp}) \\
&\quad - \frac{\delta_j T_j}{m_j n_j^{(0)}} \underline{\nabla}_{\parallel} n_j^{(1)} - \nu_{ei} \underline{v}_{j\parallel}^{(1)}.
\end{aligned}
\tag{III-13}$$

Expansion of n_j and $\underline{v}_j$ in (III-1) and (III-5) yields

$$\underline{\nabla}_{\parallel} \cdot \underline{v}_{j\parallel}^{(1)} = - \frac{\dot{n}_j^{(1)}}{n_j^{(0)}} \quad \text{and}
\tag{III-14}$$

$$\nabla_{\parallel}^2 \phi = - 4\pi \sum q_j n_j^{(1)}.
\tag{III-15}$$

Operating on (III-13) with $\underline{\nabla}_{\parallel} \cdot$ and substituting in (III-14) and (III-15) gives

$$\begin{aligned}
\ddot{n}_j^{(1)} + \nu_{ei} \dot{n}_j^{(1)} &- \frac{\delta_j T_j}{m_j} \nabla_{\parallel}^2 n_j^{(1)} \\
&= - \frac{4\pi q_j n_j^{(0)}}{m_j} \sum q_j n_j^{(1)} + \frac{q_j^2 n_j^{(0)}}{2m_j^2 c^2} \nabla^2 (\underline{A}_{\perp} \cdot \underline{A}_{\perp}).
\end{aligned}
\tag{III-16}$$

The terms on the right represent the anomalous effects of electric field inhomogeneities and magnetic field distortion of transverse motion, respectively. These terms are analogous to those on the RHS of eq. (II-19).

The completed electron and ion wave equations are contained in (III-16). The wave equation for electrons is

$$\begin{aligned} \ddot{n}_e^{(1)} + 2\gamma_p \dot{n}_e^{(1)} - \delta_e v_e^2 \nabla_{\parallel}^2 n_e^{(1)} \\ = -\omega_{pe}^2 (n_e^{(1)} - n_i^{(1)}) + \frac{e^2 n_e^{(0)}}{2m_e^2 c^2} \nabla_{\parallel}^2 (\underline{A}_{\perp} \cdot \underline{A}_{\perp}), \end{aligned} \quad (\text{III-17})$$

where $v_e = (T_e/m_e)^{1/2}$ is the electron thermal velocity and $\gamma_p = \nu_{ei}/2$ represents collisional damping. The fluid model is incapable of describing Landau damping. Hence, when Landau damping is dominant over collisional damping, it should be used in γ_p . Due to the large mass of the ions, direct damping of the ion motion and response of the ions to $\underline{A}$ is neglected in the ion wave equation. Thus, the wave equation for ions is

$$\begin{aligned} \ddot{n}_i^{(1)} - \delta_i v_i^2 \nabla_{\parallel}^2 n_i^{(1)} \\ = -\omega_{pi}^2 (n_i^{(1)} - n_e^{(1)}) \end{aligned} \quad (\text{III-18})$$

where $v_i = (T_i/m_i)^{1/2}$ is the ion thermal velocity.

Eqs. (III-12), (III-17), and (III-18) are the desired wave equations describing $\underline{A}$, n_e and n_i . The em field wave equation describes all em wave propagation. We now separate this equation into the pump and scattered components through expansion of $\underline{A}$:

$$\underline{A} = \underline{A}_0 + \underline{A}', \quad (\text{III-19})$$

where $\underline{A}_0$ is the incident pump wave and $\underline{A}'$ is the perturbing scattered

wave. Substitution of (III-19) into (III-12), and subsequent separation into 0th and 1st orders yields

$$-c^2 \nabla^2 \underline{A}_0 + \ddot{\underline{A}}_0 = -\omega_{pe}^2 \underline{A}_0 \quad (\text{III-20})$$

$$-c^2 \nabla^2 \underline{A}' + 2\gamma_{em} \dot{\underline{A}}' + \ddot{\underline{A}}' = -\omega_{pe}^2 \underline{A}' - \omega_{pe}^2 \frac{n_e^{(1)}}{n_e^{(0)}} \underline{A}_0, \quad (\text{III-21})$$

where the subscripts for perpendicular and parallel directions have been dropped. With $\underline{A}_0 = -\frac{cE_0}{\omega_0} \sin(\omega_0 t - k_0 z) \hat{e}_\perp$, (III-20) reduces to the standard dispersion relation for em waves in a plasma: $\omega_0^2 = \omega_{pe}^2 + k_0^2 c^2$. Inclusion of the terms containing γ_{em} and $n_e^{(1)} \underline{A}'$ in (III-20), would have provided for a higher order equation describing pump depletion. Eq. (III-21) describes growth of scattered waves, $\underline{A}'$. The second term on the right represents current sources due to es perturbations, $n_e^{(1)}$, driven by the pump, $\underline{A}_0$. The phase matching relationship between em pump, em scattered and es waves is defined by this term.

To constructively drive the scattered wave, the phase of $n_e^{(1)} \underline{A}_0$ must be close to the phase of $\underline{A}'$ in the first term on the RHS of eq. (III-21). The first term is the driven current or resonant source. One may formally obtain the phase matching relationship through Fourier analysis of (III-21). Equivalently, one may consider $n_e^{(1)}$ or $\underline{A}'$ at some frequency and examine (III-21) for the matching frequency of the other component. Considering a harmonic density perturbation:

$$n_e^{(1)} = \bar{n}_e^{(1)} \cos(\omega t - kx),$$

the strongest coupling occurs with scattered waves given by

$$\underline{A}'_{\pm} = \bar{A}'_{\pm} \sin[(\omega_0 \pm \omega)t - (k_0 \pm k)x] \hat{e}_1;$$

ie, $\omega_0 = \omega \pm \omega'$ and $k_0 = k \pm k'$. $\underline{A}'_{-}$ and $\underline{A}'_{+}$ are the Stokes and anti-Stokes components, respectively.

A similar expansion of $\underline{A}$ in the electron wave equation, (III-17), yields

$$\begin{aligned} \ddot{n}_e^{(1)} + 2\gamma_p \dot{n}_e^{(1)} - \delta_e v_e^2 \nabla^2 n_e^{(1)} \\ = -\omega_{pe}^2 (n_e^{(1)} - n_1^{(1)}) + \frac{e^2 n_e^{(0)}}{2m_e^2 c^2} \nabla^2 (\underline{A}_0 \cdot \underline{A}'), \end{aligned} \quad (\text{III-22})$$

where again the driving terms on the right establish the same phase matching conditions above. The second term on the right represents a beating of the pump and scattered em waves, which reinforces the density bunching of the es wave. This term originated from the $\underline{v} \times \underline{B}$ contribution to the momentum equation. Hence, an alternative view of this term is the interaction of the jitter velocity, v_0 , with the scattered field, B' , and the jitter velocity, v' , with the pump field, B_0 . The cross product, $\underline{v} \times \underline{B}$, produces a spatially harmonic, longitudinal force which enhances the es density distribution.

Before pursuing the derivation of the dispersion relation further, it is useful to limit out attention to only one phase relationship above. The Stokes component, $\underline{A}'_{-}$, is chosen; ie, $\omega_0 = \omega + \omega'$ and $k_0 = k + k'$. Quite often in the literature this choice is made only as a matter of convenience. At other times, it

is pointed out, through a heuristic argument, that only the Stokes component is compatible with the natural resonances of the plasma. This is because the plasma resonances are generally much less than the pump frequency. Hence, in the case of the Anti-Stokes relation, $\omega = \omega_0 + \omega'$ and the driven plasma frequencies are far above the pump frequency. In the case of the Stokes component, $\omega = \omega_0 - \omega'$ and the natural plasma modes may be resonantly driven with an appropriate choice of ω' . Similarly, when expanding $\underline{A}$ in (III-22), the $\underline{A}_0 \cdot \underline{A}_0$ term would have a driving frequency of $2\omega_0$. Again, this is nonresonant with the plasma modes as compared to the retained cross term, $\underline{A}_0 \cdot \underline{A}'$. The resonance approximation is justified as long as $\omega \ll \omega_0, \omega'$.

To observe the natural resonances, we perform a Fourier analysis of the electron and ion wave equations (implying the replacements $\partial/\partial t \rightarrow i\omega$ and $\partial/\partial z \rightarrow -ik$):

$$(-\omega^2 - 2i\gamma_p\omega + \delta_e v_e^2 k^2 + \omega_{pe}^2)n_e - \omega_{pe}^2 n_i = \frac{e^2 n_e^{(0)}}{m_e^2 c^2} k^2 A_0 A' \quad (\text{III-23})$$

$$(-\omega^2 + \delta_i v_i^2 k^2 + \omega_{pi}^2)n_i - \omega_{pi}^2 n_e = 0. \quad (\text{III-24})$$

By examination of (III-23), one readily recognizes from the structure that the electrons have a resonance given by $\omega^2 = \omega_{pe}^2 + \delta_e v_e^2 k^2$. (The ion term on the left side of the equation is small at the high frequency of this resonance and is neglected in arriving at this expression.) This resonance relation is just the dispersion relation for electron plasma waves. The ion resonance is less visible in the equations but results from the terms $\delta_e v_e^2 k^2 n_e$ in (III-23) and $\omega_{pi}^2 n_e$

in (III-24). Solving for n_e in (III-23) and substituting into (III-24) gives

$$\begin{aligned} \omega^4 - \delta_e v_e^2 k^2 \omega^2 - \omega_{pe}^2 \omega^2 - \delta_i v_i^2 k^2 \omega^2 \\ + \delta_e \delta_i v_e^2 v_i^2 k^4 + \delta_i v_i^2 k^2 \omega_{pe}^2 - \omega_{pi}^2 \omega^2 \\ + \delta_e v_e^2 k^2 \omega_{pi}^2 = 0, \end{aligned} \quad (\text{III-25})$$

where damping and driving terms are dropped. Taking the low frequency limit of (III-25) yields

$$\omega_{pe}^2 \omega^2 + \delta_e v_e^2 k^2 (\omega_{pi}^2 - \omega^2) = 0,$$

which upon rearrangement gives

$$\omega^2 = \frac{\delta_e k^2 c_s^2}{1 + \delta_e k^2 v_e^2 / \omega_{pe}^2} \approx k^2 c_s^2,$$

where c_s is the ion acoustic velocity, $c_s = (T_e/m_i)^{1/2}$. Thus, as might be expected, the ion-acoustic wave is the other plasma resonance. These electron and ion resonances correspond to SRS and SBS, respectively, in the present stimulated scatter examination. Even though SRS appears in this derivation, the focus of our attention will be on the low frequency plasma response of SBS.

The phase relationships, $\omega_0 = \omega + \omega'$ and $\underline{k}_0 = \underline{k} + \underline{k}'$, characterize the stimulated scatter mechanism. The origin of $\Delta\omega$, the resonant frequency of the es wave, can be seen as a result of the phase relationships and the em dispersion relation $\omega'^2 = k'^2 c^2$:

$$\begin{aligned}
(\omega_0 - \omega)^2 &= \omega'^2 = c^2 \underline{k}' \cdot \underline{k}' \\
&= c^2 (\underline{k}_0 - \underline{k}) \cdot (\underline{k}_0 - \underline{k}), \text{ or} \\
-2\omega_0 \omega &\approx c^2 (-2\underline{k}' \cdot \underline{k}_0 + k^2).
\end{aligned}$$

With $\Delta\omega$ defined as ω for the case of exact phase matching, the last expression implies

$$\Delta\omega = c^2 (\underline{k}' \cdot \underline{k}_0 - k^2/2) / \omega_0.$$

Since $\omega_0, \omega' \gg \Delta\omega$ (especially true for SBS, $\Delta\omega \sim \omega_{ia}$),

$$\underline{k}' \cdot \underline{k}_0 \approx k^2/2, \text{ or}$$

$$\underline{k} \approx 2\underline{k}_0$$

as assumed above. From this relation and $\underline{k}_0 = \underline{k} + \underline{k}'$, one finds

$$\underline{k}' \approx -\underline{k}_0,$$

which implies that dominant scatter occurs in the backscatter direction.

3. Dispersion Relation

The desired dispersion relation, describing the propagation of the SBS scattered em wave, is readily obtained from the wave equations above. The Fourier transformed es wave equations have been given in (III-23) and (III-24). Fourier analyzing the scattered em wave equation (III-21) gives

$$(-\omega'^2 - 2i\gamma_{em}\omega' + c^2k'^2 + \omega_{pe}^2)A' = -\frac{\omega_{pe}^2 A_0 n_e}{n_e^{(0)}}. \quad (\text{III-26})$$

Eqs. (III-23), (III-24) and (III-26) may be combined to yield a dispersion relation for ω . Solving for n_i and A' in (III-24) and (III-26) respectively, then substituting these expressions into (III-23) yields

$$\begin{aligned} & -\omega^2 - 2i\gamma_p\omega + \delta_e v_e^2 k^2 + \omega_{pe}^2 - \omega_{pe}^2 \frac{\omega_{pi}^2}{-\omega^2 + \delta_i v_i^2 k^2 + \omega_{pi}^2} \\ & = \frac{e^2 k^2 A_0^2}{m_e^2 c^2} \frac{-\omega_{pe}^2}{-\omega'^2 - 2i\gamma_{em}\omega' + c^2 k'^2 + \omega_{pe}^2}. \end{aligned}$$

Recalling $v_0 = eA_0/m_e c$, and slight rearrangement gives

$$\begin{aligned} & \frac{-\omega^2 - 2i\gamma_p\omega + \delta_e v_e^2 k^2 + \omega_{pe}^2}{\omega_{pe}^2} - \frac{\omega_{pi}^2}{-\omega^2 + \delta_i v_i^2 k^2 + \omega_{pi}^2} \\ & = \frac{k^2 v_0^2}{-\omega'^2 - 2i\gamma_{em}\omega' + c^2 k'^2 + \omega_{pe}^2}. \quad (\text{III-27}) \end{aligned}$$

Identifying the fluid susceptibilities as

$$\chi_e = \frac{-\omega_{pe}^2}{(\omega - i\gamma_p)^2 - \delta_e k^2 v_e^2}$$

$$\chi_i = \frac{-\omega_{pi}^2}{(\omega - i\gamma_p)^2 - \delta_i k^2 v_i^2},$$

we then have, approximately,

$$\chi_e^{-1} + 1 - (\chi_i^{-1} + 1)^{-1} = -k^2 v_0^2 [(\omega' + i\gamma_{em})^2 - \omega_{pe}^2 - c^2 k^2]^{-1}.$$

Rearranging the left side yields

$$\frac{1 + \chi_e + \chi_i}{\chi_e(1 + \chi_i)} = -k^2 v_0^2 [(\omega' + i\gamma_{em})^2 - \omega_{pe}^2 - c^2 k^2]^{-1}, \quad (\text{III-28})$$

as found in ref. 8. Alternatively, expanding the right side of (III-27) to lowest order (with $\omega' = \omega_0 - \omega$ and $k' = k_0 - k$) gives

$$\begin{aligned} \frac{1}{\chi_e} + \frac{1}{1 + \chi_i} &= \frac{k^2 v_0^2}{2\omega\omega_0 - 2c^2 k k_0 + c^2 k^2} \\ &= \frac{2k_0^2 v_0^2}{\omega_0(\omega - \Delta\omega)}, \end{aligned} \quad (\text{III-29})$$

as found in ref. 7, where $k = 2k_0$.

Rearranging (III-29) gives

$$= \Delta\omega + \frac{2k_0^2 v_0^2}{\omega_0} \frac{\chi_e(1 + \chi_i)}{1 + \chi_e + \chi_i}. \quad (\text{III-30})$$

Eq. (III-30) is the form of the dispersion relationship used later to calculate the es wave spectrum, and hence, the scattered em wave spectrum.

Before seeking specific expressions for growth rates and thresholds, it is instructive to qualitatively examine the dispersion relation further. The excellent discussion of dispersion relations by Briggs⁵² forms the basis for much of this review. Our main interest is in properly interpreting the dispersion relation given in (III-30) under the constraints of the model used to derive it. Hence, we will insure the proper usage of the growth rates and thresholds resulting from the dispersion relation.

The formal definition of a dispersion relation is given by the equation $D(\omega, k) = 0$, where the function D determines permissible relations between ω and k for solutions of wave propagation of the form $\Psi(z, t) = \psi_{\omega, k} \exp(i\omega t - ikz)$. The criterion of stability used by most authors is whether or not the dispersion relation of an infinite homogeneous system admits negative imaginary values of the frequency for some real values of the wave number ($\omega_i = f(k_r)$). However, two distinct types of instabilities can be identified physically. A spatial pulse can propagate along the system so that it grows only in its own reference frame but decays with time at a fixed point in space (convective instability); or, it can increase with time at every point in space (absolute instability). A related problem is the interpretation of dispersion relations for which imaginary values of the wave number are found for real values of the frequency ($k_i = f(\omega_r)$). In this case one considers the steady-state solution

of wave propagation away from some constant source. Again, two physically distinct solutions may occur. The waves may propagate and grow (amplifying waves) or be purely damped away from the source (evanescent waves).

The comparison of the two problems above, $\omega_i(k_r)$ and $k_i(\omega_r)$, has been performed and a dualism between the solutions found.⁵³ Basically, the expressions "convective instability" and "amplifying wave" refer to the same phenomenon. The difference in these concepts lies only in the excitation being considered (spatial pulse or localized, sinusoidal source) and not in the properties of the medium. The spatial amplification rate relates to the convective growth rate through $k_i = \omega_i^{\max}/v_g$, where v_g is the propagation velocity of the pulse. v_g corresponds to a pulse disturbance growing at the maximum temporal growth rate of any unstable wave ($\max \omega_i(k_r)$). This velocity is not necessarily the same as the group velocity. The difference is due to the fact that in a medium supporting unstable waves, the k spectrum of the pulse will be dominated by regions where complex ω exist rather than by the initial peak in the k spectrum.

The absolute instability and evanescent wave complete the dualism between the solutions, $\omega_i(k_r)$ and $k_i(\omega_r)$. The relationship may be stated in words: If the dispersion relation admits only real values of k for real values of ω , but allows imaginary ω for real k , then any instability is absolute; whereas, if the dispersion relation admits only real values of ω for real values of k , but allows imaginary k for real ω , then any spatial wave is evanescent. One may see the dualism in these solutions through plots of the dispersion

relations in ω, k space. The three, generic solutions to the dispersion relation are depicted schematically in fig. 5.

$k_i(\omega_r)$ solutions deal with steady state responses and therefore cannot describe an absolute instability. Conversely, $\omega_i(k_r)$ solutions are unable to describe evanescent waves. This asymmetry in the descriptions given by solutions $\omega_i(k_r)$ and $k_i(\omega_r)$ is due to the multi-dimensional nature of ω, k space. Convective instabilities and amplifying waves constitute solutions of the dispersion relation which involve all four variables, ω_r , ω_i , k_r , and k_i . Hence, they allow description by both $\omega_i(k_r)$ and $k_i(\omega_r)$ functions. Absolute instabilities and evanescent waves constitute solutions involving different subsets of the variables. Absolute instabilities correspond to $\omega_i(k_r)$ solutions ($k_i(\omega_r) = 0$) and evanescent waves correspond to $k_i(\omega_r)$ solutions ($\omega_i(k_r) = 0$).

With these solutions of dispersion relations in mind, we must now distinguish what class of instabilities is described by the infinite homogeneous model, as embodied in (III-30). To explicitly examine the possible solutions of a dispersion relation requires a detailed analysis in the complex plane. The relation of interest, (III-30), is so involved as to preclude such an analysis. However, physical arguments can be made which lead to a consistent description of the solutions found in the infinite homogeneous model. The description is substantiated by the results of others.

In the infinite homogeneous model only one instability threshold will be found. Is this a convective or absolute threshold? Does the "infinite homogeneous growth rate", γ_0 , above threshold correspond to convective or absolute growth? One soon realizes that an infinite

Figure 5. The generic solutions of dispersion relations.

medium which supports convective growth, will only display absolute growth. An infinite distributed source of pulses all growing at γ_0 will continually arrive at any observation point such that the growth at that fixed point will be at the rate γ_0 . Thus, convective growth in an "infinite" medium appears as absolute growth. "True" absolute growth (not depending on the medium being infinite) may occur at some higher threshold. However, the absolute growth rate, γ_a , is typically much less than the convective rate and would never be seen in an infinite medium. It has been shown in fact, that the infinite homogeneous instability threshold is the convective threshold²⁸ and that the infinite homogeneous growth rate, γ_0 , is the growth rate of a propagating pulse in its own reference frame.⁵⁴

To distinguish between convective (amplifying) and absolute growth requires a finite system. In this case, growth at the convective rate at a fixed observation point, will end when a pulse from the boundary arrives. Thereafter, growth will cease if the system is only convectively unstable or growth will continue at the slower absolute rate if the system is absolutely unstable. To quantitatively establish these regimes in the finite case requires examining the dispersion equation for k_i as well as ω_i . This has been done by Forslund, et.al.⁸ Their solutions of the form $\omega_i(k_r)$ lead to a convective growth rate of γ_0 and an absolute growth rate of $\gamma_a = 2(c_s/c)^{1/2}\gamma_0$. Solving for $k_i(\omega_r)$, an amplification rate of $k_i = \gamma_0^2/\gamma_p c$ is found. Presumably the spatial amplification rate is related to the temporal convective rate by $k_i = \gamma_0/v_g$, where v_g is the propagation velocity of a pulse disturbance.

Thus, γ_0 is the proper growth rate in the finite, convective regime which this experiment operated in. Furthermore, the spectral dependence of γ_0 should correctly describe the stimulated backscatter observed in the experiment. The concepts discussed above are summarized in fig. 6.

A further result of the infinite homogeneous model, which bears mentioning, is the distinction of weak coupling and strong coupling. In the weak coupling limit, $\gamma_0 \ll \Delta\omega - \omega_{ia}$, the pump wave does not greatly alter the dispersion relation for the electrostatic wave, $\omega_{es} = \text{Re}(\omega) \approx \Delta\omega$, and the growth rate is given correctly by γ_0 . In the strong coupling limit, $\gamma_0 > \Delta\omega - \omega_{ia}$, the pump wave determines the properties of the electrostatic wave, giving rise to a modified electrostatic wave dispersion relation describing "quasi-modes". These quasi-modes were mentioned in the interaction summary of the second chapter. In this case, $\omega_{es} \gg \Delta\omega$, and the growth rate dependence on v_0 is weaker than that of γ_0 . Depending on the es wave damping rate, γ_p , the strong coupling limit may lie below or above the absolute threshold. For the present experiment, γ_0 lies in the weak coupling, convective regime, even though strong coupling occurs below the absolute threshold. The ordering of these regimes for the present experiment is displayed in fig. 7. The different regimes are delineated above the axis in fig. 7 with respect to the infinite homogeneous growth rate. Below the axis the regimes are shown as a function of incident power.

Figure 6. Instability regimes.

Figure 7. Instability regime ordering for present experiment.

4. Growth Rate

The first growth rate of importance occurs near threshold; ie, where damping is important. Setting this growth rate equal to zero will establish the instability threshold. The other growth rate of importance occurs well above threshold; ie, where damping is negligible. This growth rate, the "infinite homogeneous growth rate", is the fundamental rate describing SBS.

To derive the near-threshold growth rate requires retention of damping terms. Hence, we begin with eq. (III-28):

$$\frac{1}{\chi_e} + \frac{1}{1 + \chi_i} = \frac{-k^2 v_o^2}{(\omega + i\gamma_{em})^2 - \omega_{pe}^2 - c^2 k'^2}.$$

Expanding the right side as before, but keeping damping terms, yields

$$\frac{1}{\chi_e} + \frac{1}{1 + \chi_i} = \frac{-k^2 v_o^2}{2i\omega_o(\gamma_{em} - \gamma) + (2\gamma_{em}\gamma + \omega_{es}^2 - \gamma_{em}^2)},$$

where $\omega = \omega_{es} + i\gamma$, and the ordering employed is $\gamma < \gamma_{em}$, γ_p , ω_{es} , $kv_i \ll kv_e$, $\omega_{pi} \ll \omega_{pe} \ll \omega_o$. Considering the ion, low frequency response, results in susceptibilities of the form

$$\chi_e = \frac{\omega_{pe}^2}{\delta_e k^2 v_e^2}$$

$$\chi_i = \frac{-\omega_{pi}^2}{2i\omega_{es}(\gamma - \gamma_p) + (2\gamma\gamma_p - \gamma_p^2 - \delta_i k^2 v_i^2 + \omega_{es}^2)}.$$

Thus, we have

$$\begin{aligned}
& \frac{-k^2 v_0^2}{2i\omega_0(\gamma_{em}-\gamma) + (2\gamma_{em}\gamma + \omega_{es}^2 - \gamma_{em}^2)} \\
&= \frac{1}{\chi_e} + \frac{1}{1 + \chi_i} \\
&= \frac{\delta_e k^2 v_e^2}{\omega_{pe}^2} \\
&+ \frac{2i\omega_{es}(\gamma - \gamma_p) + (2\gamma\gamma_p - \gamma_p^2 - \delta_i k^2 v_i^2 + \omega_{es}^2)}{2i\omega_{es}(\gamma - \gamma_p) + (2\gamma\gamma_p - \gamma_p^2 - \omega_{pi}^2)} \\
&= \frac{\delta_e k^2 v_e^2}{\omega_{pe}^2} + 1 + \frac{\omega_{pi}^2 - \delta_i k^2 v_i^2 + \omega_{es}^2}{2i\omega_{es}(\gamma - \gamma_p) + (2\gamma\gamma_p - \gamma_p^2 - \omega_{pi}^2)} \\
&\approx 1 + \frac{\omega_{pi}^2}{2i\omega_{es}(\gamma - \gamma_p) - \omega_{pi}^2} \\
&= \frac{2i\omega_{es}(\gamma - \gamma_p)}{2i\omega_{es}(\gamma - \gamma_p) - \omega_{pi}^2}. \tag{III-31}
\end{aligned}$$

Expanding (III-31) and considering the real part yields

$$\begin{aligned}
-4\omega_{es}\omega_0(\gamma - \gamma_p)(\gamma_{em} - \gamma) &= k^2 v_0^2 \omega_{pi}^2, \text{ or} \\
-4\omega_{es}\omega_0[-\gamma_p\gamma_{em} + \gamma(\gamma_{em} + \gamma_p)] &\approx k^2 v_0^2 \omega_{pi}^2.
\end{aligned}$$

Solving for γ gives

$$\gamma = \left(\frac{-k^2 v_o^2 \omega_{pi}^2}{4\omega_{es}\omega_o} + \gamma_p \gamma_{em} \right) / (\gamma_{em} + \gamma_p), \quad (\text{III-32})$$

as found in ref. 8.

To find γ_o , the infinite homogeneous growth rate, we begin with (III-30):

$$i\gamma \approx \omega - \Delta\omega = \frac{-2k_o^2 v_o^2}{\omega_o} \frac{\chi_e(1+\chi_i)}{1 + \chi_e + \chi_i}. \quad (\text{III-33})$$

Neglecting damping reduces the susceptibilities to

$$\chi_e = \frac{\omega_{pe}^2}{\delta_e k^2 v_e^2}$$

$$\chi_i = \frac{\omega_{pi}^2}{-\omega^2 + \delta_i k^2 v_i^2}.$$

Thus, we have

$$\begin{aligned} i\gamma &= \frac{-2k_o^2 v_o^2}{\omega_o} \frac{\omega_{pe}^2 \omega_{pi}^2}{\omega_{pe}^2 (-\omega^2 + \delta_i k^2 v_i^2) + \omega_{pi}^2 \delta_e k^2 v_e^2} \\ &= \frac{-2k_o^2 v_o^2}{\omega_o} \frac{\omega_{pi}^2}{-\omega^2 + \delta_i k^2 v_i^2 + \omega_{es}^2} \end{aligned}$$

$$= \frac{-2k_0^2 v_0^2}{\omega_0} \frac{\omega_{pi}^2}{(-\omega_{es}^2 - 2i\omega_{es}\gamma + \gamma^2) + \delta_1 k^2 v_1^2 + \omega_{es}^2}$$

$$= \frac{-2k_0^2 v_0^2}{\omega_0} \frac{\omega_{pi}^2}{-2i\omega_{es}\gamma + \gamma^2 + \delta_1 k^2 v_1^2}.$$

Solving for γ yields,

$$\gamma_0 = \frac{k_0 v_0 \omega_{pi}}{(\omega_0 \omega_{es})^{1/2}}, \quad (\text{III-34})$$

as found in the literature.

Solving for the high frequency, electron response (SRS) would have given

$$\gamma = \left(\frac{-k^2 v_0^2 \omega_{pe}}{4\omega_0} + \gamma_p \gamma_{em} \right) / (\gamma_p + \gamma_{em}), \quad (\text{III-35})$$

just above threshold.⁸ The infinite homogeneous growth rate becomes

$$\gamma_0 \approx (v_0/c)(\omega_0 \omega_{pe})^{1/2}. \quad (\text{III-36})$$

5. Threshold

The thresholds for SBS and SRS are readily obtained from the respective growth rates near threshold. Setting γ equal to zero in (III-32) and (III-35) yields the thresholds:

$$v_{OT}^2 = 4 \frac{\gamma_p \gamma_{em}}{\omega_{pi}^2} \frac{\omega_{es} \omega_o}{k^2}$$

$$= \frac{2\gamma_p \gamma_{em}}{\omega_{pi}^2} c c_s \quad \text{SBS}$$

$$v_{OT}^2 = \frac{\gamma_p \gamma_{em}}{\omega_{pe} \omega_o} c^2 \quad \text{SRS.}$$

It should be noted that these thresholds may be rewritten in a general form for convective instability thresholds; ie,

$$\gamma_o > (\gamma_p \gamma_{em})^{1/2} = \gamma_T^{\text{conv}}.$$

The absolute instability thresholds are much higher:

$$\gamma_o > \frac{1}{2} (c/v_p)^{1/2} \gamma_p = \gamma_T^{\text{abs}},$$

where v_p is the plasma wave velocity (c_s or v_e for SBS and SRS respectively). In Table III, γ_o and v_{OT} , as well as γ_T^{conv} and γ_T^{abs} , are tabulated for the parameter regime of this experiment. The damping rates for the ions are taken from the kinetic treatment of Fried and Gould (γ_p in Table II).^{55,56} Whereas, the electron damping rates are calculated collisional damping rates.

TABLE III
GROWTH RATES AND THRESHOLDS FOR STIMULATED SCATTER

<u>Parameters</u>	<u>100 mtorr</u>	<u>250 mtorr</u>	<u>400 mtorr</u>
(on-axis, $t = 1.2 \mu s$, $z = 14. \text{ cm}$)			
<u>SBS</u>			
$\gamma_o = (v_o/c)(\omega_o/2\omega_{ia})^{1/2}\omega_{pi} \quad (10^{10}s^{-1})$	1.2	2.1	2.7
$\gamma_T^{conv} = (\gamma_p \gamma_{em})^{1/2} \quad (10^9s^{-1})$	1.0	2.6	4.2
$\gamma_T^{abs} = (c/c_s)^{1/2}\gamma_p/2 \quad (10^{12}s^{-1})$	1.1	1.0	.9
$v_o = 2.56 \times 10^5 \lambda_o p_{max}^{1/2} \quad (10^7 \text{ cm/s})$		4.7	
$v_{oT} = (2\gamma_p \gamma_{em}/\omega_{pi}^2)^{1/2}(\omega_{ia}/\omega_o)^{1/2}c \quad (10^6 \text{ cm/s})$	3.9	6.0	7.2
$P_T = 2.7 \times 10^{-5}(\gamma_p \gamma_{em}/\omega_{pi}^2)(\omega_{ia}/\omega_o)c \quad (10^8 \text{ W/cm}^2)$	2.1	4.9	7.1
$\ell_{conv} = (\gamma_p c/2\gamma_o^2) \quad (\text{cm})$	3.4	.9	.5
<u>SRS</u>			
$\gamma_o = (v_o/c)(\omega_o \omega_{pe}/2)^{1/2} \quad (10^{10}s^{-1})$	5.5	6.8	7.5
$\gamma_T^{conv} = (\gamma_p \gamma_{em})^{1/2} \quad (10^9s^{-1})$	.7	4.1	10.
$\gamma_T^{abs} = (c/v_e)^{1/2}\gamma_p/2 \quad (10^{11}s^{-1})$	.9	4.5	9.9
$v_o = 2.56 \times 10^5 \lambda_o p_{max}^{1/2} \quad (10^7 \text{ cm/s})$		4.7	
$v_{oT} = (2\gamma_p \gamma_{em}/\omega_o \omega_{pe})^{1/2}c \quad (10^6 \text{ cm/s})$	.6	2.9	6.3
$P_T = 2.7 \times 10^{-5}(\gamma_p \gamma_{em}/\omega_o \omega_{pe})c^2 \quad (10^8 \text{ W/cm}^2)$	.04	1.1	5.4
$\ell_{conv} = (\gamma_p c/2\gamma_o^2) \quad (\text{cm})$	.07	.2	.3

6. Saturation

The theory, described above, retained only the leading order terms. This resulted in a linear growth rate for S'S daughter waves. It is clear from the viewpoint of energy conservation that these waves cannot grow indefinitely; ie, a saturation mechanism must appear at some point in the waves growth. At such time the pump and daughter waves would experience amplitude oscillations as a balance between saturation and growth is approached. A nonlinear temporal solution is required to describe saturation. Only a cursory review of saturation is presented since more restrictive effects on SBS are found due to plasma inhomogeneities.

In parametric scattering, the daughter wave amplitudes obey certain relations first formulated by Manley and Rowe:⁵⁷

$$\frac{\partial n_{es}}{\partial t} = \frac{\partial n'}{\partial t} = - \frac{\partial n_0}{\partial t}. \quad (\text{III-37})$$

The scattering process may be viewed as the annihilation of a photon with frequency ω_0 and the simultaneous generation of a photon at ω' and plasmon at ω_{es} . Eq. (III-37) says that the increase in scattered plasmon and photon density is equal to the corresponding decrease in pump photon density. This is just a statement of conservation of quanta.⁵⁸ The relation may be rewritten as

$$\frac{\epsilon_{es}}{\hbar\omega_{es}} = \frac{\epsilon'}{\hbar\omega'} = \frac{\epsilon_0}{\hbar\omega_0}, \text{ or}$$

$$\frac{\epsilon'}{\epsilon_0} = \frac{\omega'}{\omega_0} \approx 1, \text{ for SBS.}$$

Thus, the energy of the backscattered em wave can approach that of the incident pump wave. This corresponds to maximum pump depletion. This level of saturation may never be approached however due to a finite medium or nonlinear effects.

Somewhat lower saturation levels have been discussed.⁸ Ion wave steepening may occur due to higher order terms in the original continuity and momentum equations. This wave steepening can result in particle trapping or wave breaking with subsequent strong damping. However, such effects appear to occur in the vicinity of maximum pump depletion. Hence, relatively large levels of backscatter could still be expected. Inhomogeneities of the medium result in phase mismatching of the waves. As seen in the next section, this leads to greater restrictions on the backscattered power level.

7. Inhomogeneous Effects

A "homogeneous" plasma typically refers to plasma parameters being constant over the volume of the plasma. A "finite" plasma is simply a statement on the volume of the plasma. However, a finite plasma constitutes an inhomogeneous plasma in its lack of translational symmetry. Before considering inhomogeneities due to parameter variation, the consequences of finite plasmas will be studied. Through this study, various types of thresholds will be delineated. This classification will be useful when interpreting "thresholds" related to parameter variations.

In previous sections, "threshold" has been denoted by imaginary ω being negative; ie, $\gamma_+ t_+ > 0$, where t_+ is the time of gain and $\gamma_+ t_+$ signifies the number of e-foldings, n_{ef} , above the initial, signal

noise level. In the convective regime, it has been noted that gain is limited by the size of the plasma; ie, $t_+ = l_+/v_g$. Thus, the criterion above may not be sufficient. To insure significant gain, one requires $\gamma_+ t_+ = \gamma_+ l_+/v_g > 1$. In an actual experiment, an even more stringent criterion must be adopted. There is inherent noise in any detection arrangement. Hence, one requires $\gamma_+ t_+ > \ln(S_{DN}/S_N)$, where S_N is the initial signal level and S_{DN} is the diagnostic noise level.

A further limitation possibly encountered in a finite plasma, is regions of net damping. The decrease in signal level in these loss regions must be accounted for in any gain calculation; ie, $n_{ef} = \gamma_+ t_+ - \gamma_- t_-$. Summarizing these various threshold criteria, we have

- 1) positive gain: $\gamma_+ t_+ - \gamma_- t_- > 0$
- 2) significant gain: $\gamma_+ t_+ - \gamma_- t_- > 1$
- 3) detectable gain: $\gamma_+ t_+ - \gamma_- t_- > \ln(S_{DN}/S_N)$.

In an inhomogeneous plasma where parameters vary over space, the plasma wave will have a different wavelength in each part of the plasma. Hence, the phase matching conditions, $\omega_0 = \omega_{es} + \omega'$ and $\underline{k}_0 = \underline{k}_{es} + \underline{k}'$, can be satisfied only over limited regions. The mismatched phase may grow so large as to destroy the coupling which allowed growth of the daughter waves. This situation has been considered,⁵⁹ and results in the "significant gain" threshold of

$$\frac{\pi \gamma_0^2}{cc_s \kappa'} > 1,$$

where $\kappa' = \partial k / \partial z$ represents the rate of phase mismatching. In SRS, the electron plasma wave dispersion relation is strongly dependent on density. Hence, κ' is determined by the density gradient length. In SBS, the ion wave dispersion relation depends strongly on temperature and the temperature gradient length determines κ' .

Growth will cease beyond the point where the phase is exactly mismatched, $\Delta k l_+ = \pi$, where $\Delta k = \kappa' l_+$.²⁸ Thus, we have $l_+^2 \approx \pi / \kappa'$. For an ion-acoustic wave, $\omega \sim k T^{1/2}$. Hence, at the mismatch point we have

$$k = k_0 \frac{T_0^{1/2}}{T^{1/2}}$$

$$\kappa' = - (1/2) k_0 T_0^{1/2} T^{-3/2} (\partial T / \partial z)$$

$$l_+^2 \approx \frac{\pi}{\kappa'} = \frac{2\pi T^{3/2}}{k_0 T_0^{1/2} (\partial T / \partial z)}$$

$$l_+ \approx \left[\frac{2\pi T}{k (\partial T / \partial z)} \right]^{1/2} \approx .3 \text{ cm},$$

where the parameters for 250 mtorr fill pressure at 1.2 μ s are used. This coupling length is shorter than the gain length. Therefore, significant SBS is not expected. Temperature measurements at 1.4 μ s indicate that the temperature gradient was reduced greatly near peak power and was again prohibitive just after peak power. This

temperature history and corresponding SRS coupling is in agreement with the observed SRS. The coupling length due to temperature inhomogeneity will be discussed further in the context of the experimental results.

The phase matched, SRS coupling length may be estimated in a similar manner. In this case density inhomogeneity results in phase mismatching. The dispersion relation for electron plasma waves is $\omega^2 = \omega_{pe}^2 + 3k^2 v_e^2$. Hence, we have

$$k = \left(\frac{\omega^2 - \omega_{pe}^2}{3v_e^2} \right)^{1/2} \approx \left(\frac{\omega_{pe}^2}{3v_e^2} \right)^{1/2}$$

$$\kappa' \approx \frac{\omega_{pe}}{2(3)^{1/2} v_e} \frac{\partial n / \partial z}{n}$$

$$l_+ \approx \left(\frac{2\pi 3^{1/2} v_e}{\omega_{pe}} \frac{n}{\partial n / \partial z} \right)^{1/2} \approx 0.03 \text{ cm},$$

where $\partial n / \partial z \approx 10^{16}$ along the axis between the focus region and the front end of the coil. The coupling length is significantly shorter than the gain length for SRS. In contrast to the axial temperature gradient, this density gradient remains prohibitive and strong SRS coupling is never expected.

8. Faraday Rotation

Due to Faraday rotation, the incident and backscattered em waves will have their plane of polarizations rotated. The length for a 90° , relative rotation is given by⁶⁰

$$l_f = \frac{m_e^2 c^2 (\pi/2) \omega_0^2}{2\pi e^3 B n_e}$$

$$\approx 8.4 \times 10^{-7} \frac{\omega_0^2}{B n_e} \approx 5 \text{ cm.}$$

This relative rotation of the waves' polarization will result in a spatial modulation of the effective pump strength, since the beating of these waves provides the ponderomotive driving force for the SBS interaction.²⁸ The ponderomotive force dependence may be written as

$$|E_{NL}| \propto |\langle \underline{v}_0 \times \underline{B}' \rangle| \propto |\langle \underline{v}_0 \cdot \underline{E}' \rangle|$$

$$\propto v_0 E' |\cos \theta|$$

$$= \hat{v}_0 E',$$

where θ is the relative angle between the waves planes of polarization and $\hat{v}_0 = v_0 |\cos \theta|$ is the effective pump strength. In regions where the effective pump power falls below the convective instability threshold, net damping would occur. Balancing the gain and damping over a cycle gives

$$\int_{l_+} (\gamma_0^2 / \gamma_p) dz \geq \int_{l_-} \gamma_{em} dz$$

$$(\gamma_0^2 / \gamma_p) \int_{l_+} \cos^2(az) dz \geq \gamma_{em} l_-$$

$$(\gamma_0^2 / \gamma_p a) [a l_+ / 2 + \sin(2a l_+) / 4] \geq \gamma_{em} l_-, \text{ or}$$

$$\gamma_0^2 \geq \gamma_p \gamma_{em} l_- / [l_+ + \sin(2a l_+) / 4a], \quad (\text{III-38})$$

where $a = \pi / 2l_f$, $l_+ = (1/a) \cos^{-1}(v_{0T}/v_0)$ and $l_- = l_f - l_+$. When

l_+/l_- becomes small, the usual threshold ($\gamma_0^2 > \gamma_p \gamma_{em}$) is expected to be higher by the small factor in (III-38). Whenever $i_f < l_{conv}$, the periodic variation in γ_0 results in a factor of 2 increase in the gain length.

These calculations assume linearly polarized backscatter. Since the normal em modes of a magnetized plasma are circularly polarized, the backscatter may not be linearly polarized. The degree of elliptical polarization would depend on the noise source polarization. Elliptically polarized backscatter should result in greater gain than linearly polarized backscatter (in the presence of Faraday rotation). Regardless of polarization, a magnetized plasma exhibiting strong Faraday rotation is expected to have reduced SBS gain.

B. Backscatter Diagnostic System

1. Arrangement

In fig. 8, the details of the backscatter detection system are shown. The viewed backscatter was collected by the CO_2 laser, focusing mirror and a beam splitter. A HgCdTe photovoltaic detector (S.A.T., class A4), S_1 , was situated at the exit slit of a .5m, Ebert, scanning monochromator (Jarrel-Ash). The monochromator used a 5.2 cm x 5.2 cm, 150 grooves/mm grating (Bausch & Lomb) blazed at 37° for $8 \mu m$ radiation in the first order. The instrument had a calculated linear dispersion of $88 \text{ \AA/mm}$ and resolving power of 7800:

$$\Delta\lambda/l = d\cos\theta/fn = 88 \text{ \AA/mm}$$

$$\lambda/\Delta\lambda = Wn/d = 7800,$$

Figure 8. The backscatter detection arrangement.

where d is the groove spacing, f is the focal length, n is the order, and W is the grating width.⁶¹ The measured instrument resolution with 200 μm slits was 10 $\text{\AA}$ HWHM. The instrument function, $S_1(\lambda)$, is displayed in fig. 9. Calibration was performed with a tunable CW, CO_2 laser (Korad).

In addition to the wavelength-resolved detector, S_1 , a second HgCdTe photovoltaic detector (S.A.T., class A2), S_2 , was used to monitor the wavelength-integrated backscattered signal. To maintain S_2 below saturation, a layer of 2.4 mm thick, polyethylene was used as an attenuator. The detectors, S_1 and S_2 , had a cutoff frequency of 10 MHz. Hence, the actual time history of the very fast SBS signal (risetime ≤ 1 ns) is not detectable and an integrated signal, limited by the oscilloscopes 5 ns risetime, was observed.

Two detection geometries were used: 1) axial geometry - scattering volume viewed through the CO_2 laser focusing mirror ($.0^\circ - .4^\circ$ viewing angle) and 2) annular geometry - scattered radiation collected by CO_2 laser focusing mirror ($.4^\circ - 2.4^\circ$ viewing angle). It was found that both signals, S_1 and S_2 , were significantly stronger (by a factor ≥ 3) in the narrower field of view. However, the wavelength shifted signal (SBS) retained similar spectral and temporal characteristics in both fields of view. Thus, the more convenient annular geometry was predominantly used and produced the data discussed here.

Figure 9. The measured instrument function, $S_I(\lambda)$.

2. Data Analysis

The stimulated backscatter signal consisted of a single, strong pulse occurring within ± 15 ns of the incident beam's peak power. The signal appeared at the expected, shifted wavelength (red shifted by an amount corresponding to ω_{ia}). In fig. 10, typical data (400 mtorr fill pressure) are displayed for two successive shots in which stimulated backscatter occurred on the second but not the first shot (verified by an increase in both S_1 and S_2). Some stray laser reflectance is observed in addition, with the integrated-wavelength detector S_2 . This background backscatter ($\leq 0.4\%$) can easily be discerned by its time profile and is subtracted out in the analysis.

The two shots in fig. 10 illustrate the main difficulty in the backscatter study; an intermittent shifted backscatter signal. With the CO_2 laser operated at its peak output of ~ 200 J, the shifted backscatter signal occurred at an optimum percentage of shots of 80%. In this operating regime a normalized, spectral scan of the shifted backscatter radiation was made at 250 mtorr fill pressure with scanning steps of $5.2 \text{ \AA}$; ie,

$$S_D(\Delta\lambda) = A \frac{S_1(\Delta\lambda)}{S_2}, \quad (\text{III-39})$$

where A is a constant function of detection geometry and detector sensitivity. The result of the scan is shown in fig. 11. The peak in the scan falls at $\sim +51 \text{ \AA}$ with respect to the incident P-20 laser line. The $10 \text{ \AA}$ HWHM of the spectrum appears limited by the

Figure 10. Related stimulated Brillouin backscatter data. The top trace is the spectrally integrated backscatter signal (S_2) on the middle display and the spectrally resolved signal (S_1 , +54 A) on the lower display.

B_z

Photon Drag
(.5 μ s/div)

Spectrally
Integrated
Backscatter

Photon Drag
(.2 μ s/div)

Spectrally
Resolved
Backscatter
(P-20 + 54 Å)

Photon Drag
(.1 μ s/div)

Without
Stimulated
Backscatter

With
Stimulated
Backscatter

Figure 11. Comparison of normalized theoretical and measured spectral profiles of stimulated backscatter with $T_i = T_e = 36$ eV and $n_{ef} = 14$. The dashed and solid lines represent S_B and S_T , respectively. The points represent the data S_D . Error bars indicate rms deviations of several shots.

instrument resolution of $10 \text{ \AA}$. The shift in frequency compares favorably with the corresponding ion acoustic frequency.

An attempt to measure the dependence of backscatter power on incident power was made. The incident power was obtained from normalizing the incident, time resolved, photon drag signal to the incident energy calorimeter signal. A large amount of scatter in the backscatter power dependence was observed and attributed to operation near threshold. The intermittent occurrence of stimulated backscatter may be considered a manifestation of this power scatter. Though no correlated dependence of SBS power with incident power was obtained, a lower power cutoff or threshold was measured. These experimental thresholds were $1.2 \times 10^{10} \text{ w/cm}^2$ for 250 mtorr fill pressure and $0.7 \times 10^{10} \text{ w/cm}^2$ for 400 mtorr fill pressure. These are cutoffs for the peak incident power. SBS was only observed near the peak power and never during the rest of the laser pulse.

The stray reflectance, mentioned above provided a means of estimating the absolute power in the stimulated backscatter signal. Calorimeter measurements of backscatter show the total backscattered energy (in annular geometry) to be $\sim .6 \text{ J}$; ie, $\sim .4\%$ reflectivity. This energy represents only the stray reflectance, since the stimulated signal was extremely short and could contribute little backscattered energy. The $.4\%$ stray reflectivity normalizes S_2 to the incident power. With this normalization, a typical stimulated backscatter signal had a peak power of $\leq 1 \times 10^6 \text{ w}$. Allowing for the slow response of the oscilloscope in comparison to the signal rise time ($\sim .1 \text{ ns}$) results in an absolute power estimate of $\leq 5 \times 10^7 \text{ w}$ for the stimulated backscatter signal.

C. Comparison with Theory

1. Power Dependence

a. Power Threshold and Signal Gain

The fact that SBS occurred temporally near the peak incident power strongly suggests that the laser power only marginally exceeded threshold. However, the theoretical convective instability threshold, $P_T \approx 6 \times 10^8 \text{ w/cm}^2$, was significantly below the laser power capability of $\sim 3.5 \times 10^{10} \text{ w/cm}^2$. This disagreement in homogeneous theory and experimental observations is consistent with strong SBS coupling constrained to a narrow time window. The time window for allowed gain corresponds to a transient period of reduced axial temperature gradient.

As previously discussed, plasma phase variation due to a longitudinal temperature gradient seems quite strong. Before and after the first laser spike ($t = 1.2$ and $1.4 \mu\text{s}$, respectively), the electron temperature is 13 and 34 eV at $z = 14 \text{ cm}$ (250 mtorr fill pressure). The corresponding temperatures at $z = 43 \text{ cm}$ are 17 and 28 eV (see fig. 28). The pump and daughter waves phase matching exists for only .3 cm before and after the first laser spike, in contrast to the gain length of 1 cm. This difference in lengths should prevent SBS from occurring. Since the axial gradient has opposite sign before and after the first laser spike, a period of reduced gradient and stronger SBS coupling occurs between these two times. During this period the laser reaches peak power and SBS is observed. Significant SBS gain constrained to this time period is consistent with the above-mentioned, single-pulse characteristic, the

sometimes erratic power fluctuations, and SBS occurring only near peak power which is far above threshold.

Other mechanisms giving rise to single pulse SBS have been proposed in previous experimental discussions: laser phase variation¹⁰ and plasma motion.⁶² The predicted growth rate from random phase theory is very similar to that from coherent phase theory. Though laser phase variation seems likely on the time scale of SBS growth, the observed single pulse characteristic could not be explained by this process. Plasma motion in the present experiment is insignificant and can also be discarded as an explanation for the restricted SBS signal.

b. Intermittency and SBS Power Variation

The intermittency in SBS occurrence is believed to be a simple manifestation of the measured, shot-to-shot, variation in SBS power. It has been shown that significant SBS gain was limited to a narrow time window near peak incident power. Hence, it is expected that small shot-to-shot variations in experimental conditions could modify this time period such that SBS does not occur. Two conditions are possible which might prevent SBS on an arbitrary plasma discharge. Either SBS gain is always prohibited by an inhomogeneous temperature or reduced temperature gradients exist only when the SBS threshold is not exceeded. Critical variations which might cause shot-to-shot changes in the temperature history during laser heating are: 1) temperature and density variations resulting in a fluctuating absorptivity and 2) density profile variations in the plasma column resulting in fluctuating containment of the beam in the plasma.

c. n_{ef}^{max}

Through a measurement of stray reflectance, the absolute SBS power was estimated earlier at $\leq 5 \times 10^7$ w. An upper limit for the number of exponentiations of the SBS signal and hence, the interaction length, can be calculated through consideration of this SBS power and the SBS noise source. Thermal scattering is anticipated as the most likely noise source since the plasma is produced independently from the laser-plasma interaction. This differs from laser produced plasma experiments where turbulent scattering appears as the most likely noise source. The thermal scatter is given by eq. (III-11):

$$S_N = \frac{|S_0|}{2\pi} \sigma_T(\theta) n V S(\underline{k}', \omega') d\Omega d\omega, \text{ or}$$

$$S_N = \frac{P_0 r_0^2}{2\pi A} n V S(\Delta\omega) d\Omega d\omega,$$

where P_0 is the incident power, A is the focused beam area, r_0 is the classical electron radius and backscatter is assumed (polarization perpendicular to the detection direction). From the Salpeter approximation,^{63,64} the functional form of the spectral density function, $S(\Delta\omega)$, is given by

$$S(\underline{k}, \omega) = \frac{2\pi^{1/2}}{kv_e} \Gamma(x_e) + \frac{2\pi^{1/2}}{kv_i} Z\alpha^2 \Gamma(x_i)$$

$$\approx .15 \times 10^{-10} \text{ sec}$$

$$\Gamma(x) = \frac{\exp(-x^2)}{[1 + \operatorname{Re}[\chi(x)]]^2 + [-\operatorname{Im}[\chi(x)]]^2}$$

$$x_j = \omega/kv_j, \quad \alpha = 1/k\lambda_D,$$

where $T_e = T_i$, $\alpha \ll 1$ and χ_e and χ_i are given in the next section.

The other parameters used in calculating S_N are

$$P_0/A = 2 \times 10^{10} \text{ w/cm}^2$$

$$r_0^2 = 8.0 \times 10^{-26} \text{ cm}^2$$

$$n = 1.4 \times 10^{17} \text{ cm}^{-3}$$

$$V = .45 \text{ cm}^3$$

$$d\Omega = 5 \times 10^{-4}$$

$$d\omega = 1.7 \times 10^{10} \text{ rad/s.}$$

In fig. 12 $S(\Delta\lambda)$ is displayed for the case of 250 mtorr fill pressure. The approximate value of S_N over the wavelength region of interest is $S_N \approx 2 \times 10^{-3} \text{ w.}$ This noise level implies $n_{ef}^{\max} \leq 24$ for the SBS power of $\leq 5 \times 10^7 \text{ w.}$

2. Wavelength Dependence

a. Spectral Profile

The stimulated backscatter intensity $S_B(\Delta\lambda)$, growing exponentially from the noise intensity $S_N(\Delta\lambda)$ (presumably due to thermal scattering of the incident pump wave), is given theoretically⁸ by

Figure 12. The spectral density function, $S(\Delta\lambda)$.

$$\begin{aligned}
S_B(\Delta\lambda) &= S_N(\Delta\lambda) \exp[2\gamma(\Delta\lambda)\ell_{\text{int}}/v_g] \\
&= S_N(\Delta\lambda) \exp[n_{\text{ef}}\gamma(\Delta\lambda)/\gamma_{\text{max}}],
\end{aligned}$$

where $\gamma(\Delta\lambda)$ is the SBS growth rate and n_{ef} is the maximum number of e-foldings at the wavelength of maximum growth. $\gamma(\Delta\lambda)$ can be obtained from either coherent^{7,8} or random^{10,15} interaction theory. However, in the present work, wavelength dependence was nearly the same for the two theories; the coherent theory, as described in the theory section, is selected. In the present weak coupling regime, $\gamma(\Delta\lambda)$ is given by eq. (III-30):

$$\begin{aligned}
\gamma(\Delta\lambda) &= \text{Im}(\omega) \\
&= \Delta\omega + \frac{2k_0^2 v_0^2}{\omega_0} \frac{\chi_e(1+\chi_i)}{1 + \chi_e + \chi_i}.
\end{aligned} \tag{III-40}$$

Rather than the fluid susceptibilities defined earlier, kinetic susceptibilities are used:⁶⁵

$$\begin{aligned}
\chi_e &= \alpha^2 [1 - 2\chi_e \exp(-x_e^2) \int_0^{x_e} \exp(p^2) dp - i\pi^{1/2} \chi_e \exp(-x_e^2)] \\
\chi_i &= \alpha^2 Z(T_e/T_i) [1 - 2\chi_i \exp(-x_i^2) \int_0^{x_i} \exp(p^2) dp - i\pi^{1/2} \chi_i \exp(-x_i^2)] \\
x_e &= \omega/kv_e \quad v_e = (2T_e/m_e)^{1/2} \\
x_i &= \omega/kv_i \quad v_i = (2T_i/m_i)^{1/2} \\
\alpha &= 1/k\lambda_D.
\end{aligned}$$

With the monochromator's instrument function $S_I(\lambda)$, shown in fig. 9,

convoluted with $S_B(\Delta\lambda)$, a theoretical profile $S_T(\Delta\lambda)$ was calculated for comparison with the data scan (S_D , eq. (III-39)):

$$S_T(\Delta\lambda) = A \frac{\int S_B(\Delta\lambda - \lambda') S_I(\lambda') d\lambda'}{\int S_B(\Delta\lambda) d\lambda}.$$

$S_T(\Delta\lambda)$ was calculated for the case of 250 mtorr fill pressure. The density used was the measured value of $1.4 \times 10^{17} \text{ cm}^{-3}$. Electron-ion equipartition was $\sim 50 \text{ ns}$. Hence, $T_e \approx T_i$ was assumed. With these input values, the optimized free parameters were found to be $T_e \approx 36 \text{ eV}$ and $n_{ef} \geq 14$ for a fit of the experimental line shape. The profile of $S_T(\Delta\lambda)$ is compared with the data, $S_D(\Delta\lambda)$, in fig. 11. It is noteworthy that $S_N(\Delta\lambda)$ is slowly varying compared to $\exp(n_{ef}\gamma(\Delta\lambda)/\gamma_{\max})$ in the wavelength region of interest and hence, the normalized spectrum is only strongly sensitive to $\gamma(\Delta\lambda)$.

The electron temperature required for a theoretical fit of the SBS spectrum, 36 eV, is consistent with the 34 eV inferred from impurity ion spectroscopy measurements (see Ch. VI). The spectroscopy temperature value of 34 eV was measured just after the first laser spike; ie, $\sim 50\text{-}100 \text{ ns}$ after the peak laser power and time of SBS. This time difference in the two measurements is not considered significant enough to prevent comparison of the two values. The spectroscopy measurement required interpolation over time, of the temperature history during the laser spike. Hence, the detailed temperature history during the spike is not available. However, the actual temperature at time of SBS occurrence would be

expected to be close to that just afterwards, since increased heat conduction would inhibit any further rise in temperature.

b. Frequency Shift

The SBS profile is dependent in a somewhat complicated manner on T and n_{ef} . It is interesting to make a simpler theoretical comparison of just the wavelength shift from the incident wavelength to the peak of the SBS spectrum. This red shift should correspond to the ion-acoustic frequency. In the fluid⁶⁶ and kinetic⁵⁶ limit the ion-acoustic frequency is given by

$$\text{fluid} - \omega = k \left(\frac{T_e + 3T_i}{m_i} \right)^{1/2} \rightarrow \Delta\lambda = 32-57 \text{ \AA}$$

$$\text{kinetic} - \omega = 2.05kv_i \rightarrow \Delta\lambda = 33-59 \text{ \AA},$$

where the wavelength shifts correspond to the temperature values at the time of laser firing and just after the first laser spike, 13. eV and 34. eV respectively. It can be seen that the two calculations give very similar wavelength shifts. Furthermore, the shift at the late-time temperature is in good agreement with the experimental shift of $\sim 51 \text{ \AA}$. This is in accordance with the temperature from the profile fit, 36 eV.

c. $n_{ef}^{\min}$

Since the experimental spectrum was limited by the instrumental resolution, only a lower limit for the number of signal exponentiations is obtainable from the profile fit; ie, $n_{ef} \geq 14$.

Higher n_{ef} values would narrow the actual SBS spectrum, S_B . But the observed SBS spectrum, S_T could not become narrower than the instrument function, S_I . Hence, only a lower limit on n_{ef} is determined by the requirement that S_T be narrow enough to fit the experimental error bars.

3. Summary

a. Temporal and Spatial Constraints on Interaction

The most distinctive features of the SBS interaction were the single pulse signal and variation in power from shot-to-shot. These characteristics appear due to a restricted time period of significant SBS gain. Since this time period is limited to the time near peak power, the occurrence of single pulse SBS signals near peak power is expected. Furthermore, shot-to-shot variations in this time window could occasionally prevent SBS occurrence and give rise to fluctuations in SBS power.

Spectral analysis and SBS power considerations, for the case of 250 mtorr fill pressure, lead to a required number of exponentiations of the backscattered signal, n_{ef} , and to a corresponding interaction length, l_{int} . The bounds on the number of exponentiations are $14 \leq n_{ef} \leq 24$. The interaction length is calculated to be $13 \text{ cm} \leq l_{int} \leq 22 \text{ cm}$, from the expression $l_{int} = n_{ef} l_{gain}$. A maximum factor of two increase in this length is expected due to Faraday rotation. The length of the focus region for the incident pump was $\sim 50 \text{ cm}$. Beyond this length the peak power would be reduced below the calculated threshold. The restricted interaction length of

13-22 cm (or a maximum of 26-44 cm with Faraday rotation) is consistent with the focus region.

b. Comparison to Previous Experiments

The most similar previous experiments are found in the laser-produced plasma experiments of Massey⁹ and Offenberger¹⁰ and the Z-pinch experiments of Turechek.⁶⁷ Direct comparison of the experiments is difficult due to the large number of parameters associated with the interaction and their variation with different experiments. The significant differences of the present experiment compared with those above are

- 1) different plasma production method, geometry and parameters
- 2) convective-inhomogeneous regime (axial temperature gradient)
as compared to prior convective-homogeneous regimes
- 3) larger n_{ef} and corresponding longer ℓ_{int}
- 4) different determination of n_{ef} .

Massey and Offenberger observed prompt, single pulse SBS signals and attributed the pulsed characteristic to plasma motion and laser-phase variation, respectively. It should be noted that to obtain a prompt SBS signal required that their laser focal length be relatively short ($\sim f/5$ as compared to $\sim f/35$ in the present experiment). Turechek observed a delayed SBS signal and attributed the delay to necessary electron heating to lower the threshold below the incident power level.

The higher n_{ef} found in the present experiment is consistent with the longer plasma volume. The measurement of n_{ef} , however, is quite different than that employed before. Massey and Offenberger obtained n_{ef} from the SBS power dependence on the incident power. Turechek determined n_{ef} strictly from an absolute power measurement. In the present work a fit of the SBS spectrum provided the lower bound to n_{ef} . Massey and Offenberger both determined SBS spectrums and were able to imply the T_e/T_i ratio needed at their higher densities. Turechek measured spectral shapes but only inferred temperatures from the red shift of the peaks.

Graphical comparison between the results from the different experiments and theory is desirable. However, as mentioned, experimental differences make this comparison difficult. Possible graphical comparisons and their respective difficulties include

- 1) Threshold: a plot of P_T vs n_e for families of curves

determined by T_e - lack of data from experiments and

differences in plasma homogeneity prevent comparison

- 2) Power: a plot of SBS power vs the product nI for families of

curves determined by n_{ef} - different ions and large power

scatter in present experiment present difficulties

in interpretation

- 3) Spectral Profile: a plot of $\Delta\lambda_{HWHM}$ vs n_{ef} for families of

curves determined by T_e and T_e/T_i - n_{ef} was calculated from

this information in the present experiment and is

unavailable as an independent parameter

- 4) Frequency Shift: a plot of ω_{peak} vs T_e for families of curves determined by T_e/T_i and m_i - possible to perform but a simplistic comparison; all experiments verified agreement between spectral shift and ω_{ia} .

c. Discussion and Conclusions

In summary, prompt, single pulse stimulated Brillouin backscatter has been observed. The characteristics of the signal confirming the SBS origin are:

- 1) marked enhancement of backscattered power above stray scatter
- 2) single pulse correlation with peak incident power
- 3) red shifted signal spectrum
- 4) wavelength shift of signal peak consistent with ion-acoustic frequency
- 5) spectral fit requiring a temperature consistent with an independent temperature measurement
- 6) an experimental interaction length consistent with focus region.

The incentive and initial interest for studying SBS was its possible reduction of laser heating of plasmas. Although, SBS was observed and studied closely, it only constituted a small fraction of the incident laser energy. The SBS power reflectivity was $\leq 5\%$, whereas, the SBS energy reflectivity was $\leq 0.1\%$. These reflectivity results should be considered strictly in the context of the present experiment. The independently produced plasma results in an SBS noise level which is anticipated to be considerably less than that of a laser-produced plasma. The low density, $\sim 10^{17} \text{ cm}^{-3}$ (compared to $\sim 10^{18} \text{ cm}^{-3}$ of Massey and Offenberger), as well as strong damping due to $T_e/T_i \sim 1$, contributed to a relatively long gain length. A low initial signal level and long gain length severely limited the possible net gain. Furthermore, due to axial temperature gradients, there was only a limited time period of significant SBS gain. The inhibited SBS gain was seen in the measurements of fluctuating SBS power, high threshold and low SBS reflectivity. Thus, in conclusion, SBS did not constitute a significant loss mechanism under the constraints of this experiment.

IV. PROPAGATION

A. Density Measurement

Propagation of the CO₂ laser beam through the plasma column depends critically on the density profile of the plasma. To determine the plasma density distribution, and hence, the plasma refractive effects, end-on holographic interferometry was employed. Density measurements were also important for SFS calculations.

1. Holographic Interferometry

Holographic interferometry is based on the technique of photographically recording the amplitude and phase of an optical field on a hologram plane.^{68,69,70} With this storage capability, comparison of wavefronts originally occurring at different times is possible. In particular, one can record a wavefront in the absence of plasma and compare it with a wavefront traversing the same path in the presence of plasma. Since the two beams follow the same optical path, the usual interferometric requirements of 1) precise alignment for exact beam recombination and 2) high quality optical components for prevention of extraneous phase variations, are greatly reduced. The interferometric requirement of good temporal and spatial coherence is necessary in order to record the desired phase information. Furthermore, since the holographic reference and scene beams are angularly separated by several degrees, high resolution, photographic plates are required for recording the fine fringe structure (≥ 200 ℓ/mm). A Q-switched ruby laser was used for the

light source and Agfa-Gevaert 10E75 plates were used for the hologram recording.

The recorded fringe change, due to the refracting plasma present during one of the exposures, is given by

$$N = \int \frac{\mu - 1}{\lambda} dl,$$

where the integral is along the path length of a ray, λ is the wavelength of the laser radiation and μ is the refractive index of the plasma. The dominant, electron refractivity given in (II-16) is

$$\mu = (1 - \omega_{pe}^2 / \omega^2)^{1/2}.$$

Expanding μ for $\omega \gg \omega_{pe}$, and writing out the functional form of ω_{pe} yields

$$\begin{aligned} N &= \frac{-1}{2\pi} \frac{e^2 \lambda}{mc^2} \int n_e dl \\ &= \frac{-1}{3.2 \times 10^{17}} \int n_e dl, \end{aligned} \tag{IV-1}$$

for the ruby laser wavelength, $\lambda = 6943 \text{ \AA}$. In the present experiment, the theta-pinch plasma was viewed longitudinally. The measurement given by (IV-1) is the line integral along the plasma axis of the radial density profile.

2. Diagnostic Arrangement

In fig. 13, the holographic interferometric arrangement is displayed. A Mach-Zehnder interferometer layout was used. A substantial part of the interferometric studies was performed with a double-pulsed ruby laser.⁷¹ The double-pulsed laser allowed recording of time-differential interferograms. It was originally thought that time-differential measurements; ie, recording of relative density changes during a single shot, would be particularly useful for study of plasma variation due to laser heating. However, high densities and steep density gradients limited use of the double-pulse system in the time-differential mode.

The double-pulse system used a unique method of introducing background, reference fringes. Between exposures a planar phase shift is typically introduced in the reference beam for calibration of partial fringe shifts and for unambiguous determination of plasma gradient signs. In a conventional arrangement, a mirror is tilted slightly in the reference beam. In the present system, use is made of the orthogonal polarizations of the two pulses. A Rochon prism provides the necessary deflection of the reference beam between pulses. Since no mechanical manipulation is required between the two exposures, the time separation was reduced between the two laser pulses to $\sim 20 \mu\text{s}$. The first pulse occurred before plasma production and the second was adjusted to the time of interest during the plasma discharge. The closely spaced laser pulses avoided technical problems often encountered due to room vibrations, thermal effects, etc. Several other holographic interferometric techniques were used

Figure 13. A schematic diagram of the holographic interferometric arrangement.

in an attempt to improve the measurements of the present high plasma densities.

3. Full Field-of-View--without CO₂ Laser

Single-pass holographic interferograms were initially taken of the forward-biased plasma column over the range of 100-500 mtorr fill pressure. The CO₂ laser focusing mirror was removed allowing a full 3.6 cm i.d. discharge tube, field-of-view. The desired on-axis density minimum (for refractive trapping of the CO₂ laser beam) was observed over the above range of fill pressures. A decrease in longevity of the density minimum at higher fill pressures restricted the CO₂ laser timing to a nominal 1.2 μ s after main field initiation. Representative interferograms at this time are displayed in fig. 14. High fringe density and refractive shadowing generally made data reduction difficult. Typically one chord was available which showed clear fringes and from which peak and on-axis densities could be measured. On-axis density values as a function of time and fill pressure are shown in fig. 15. The dashed lines represent interpolated histories for use in the ionization rate code for temperature evaluation. The solid line for 250 mtorr represents an estimated history from interpolation between the 200 and 300 mtorr data sets.

Radial profiles were evaluated in detail for the case of 100 mtorr fill pressure. At this fill pressure fringe patterns were clear enough to allow complete analysis during the entire plasma discharge. Furthermore, the density gradients were small enough that the fractional error in the optical path due to refractive bending

Figure 14. End-on interferograms of plasma column with bias field.

Field-of-view is the inside diameter (3.6 cm) of the discharge tube.

100 mTORR

200 mTORR

300 mTORR

Figure 15. On-axis density as a function of time for 100, 200, 300, and 400 mtorr fill pressure. Density values for 250 mtorr are interpolated from these data.

was negligible.¹² Hence, a faithful reconstruction of the density profile for the fringe pattern is obtainable. In fig. 16 the radial plasma profile is shown as a function of time for 100 mtorr fill pressure. As can be seen in the figure, the radius of the minimum at $\sim 1.2 \mu\text{s}$ is $\sim .25 \text{ cm}$, whereas the change in density over this radial distance is $\sim .4 \times 10^{16} \text{ cm}^{-3}$. Similar values are found for 250 and 400 mtorr fill pressure cases in the next section. These values will be used in estimating the refractive containment of the CO_2 laser by the plasma density minimum.

4. Restricted Field-of-View--With CO_2 Laser

The CO_2 laser focusing mirror restricted the ruby interferometry beam to an elliptical field-of-view with a 2 cm major axis. Interferograms with this field-of-view were taken at 100, 250 and 400 mtorr fill pressure.

Initial attempts at obtaining interferograms for plasma discharges with the CO_2 laser were severely hampered by additional optical components required by the CO_2 laser beam. In particular the ruby beam suffered aberration at the theta-pinch exit end from: 1) a wire grating used to reflect the transmitted CO_2 laser beam into a calorimeter, 2) a NaCl window to pass the CO_2 laser beam, and 3) two Ge windows at Brewsters angles to separate the ruby and CO_2 laser beams. Successful interferograms were obtained with these three components removed. A quartz vacuum window replaced the salt window. CO_2 laser beam damage to the quartz window was limited to the small energies of the plasma-attenuated beam. Typical interferograms at $1.2 \mu\text{s}$ are shown in fig. 17.

Figure 16. Radial density profile as a function of time for 100 mtorr fill pressure.

100 mtorr

250 mtorr

400 mtorr

Figure 17. Restricted field-of-view interferograms at time of laser firing, 1.2 μ s. Vertical field of view is 2.0 cm.

Figure 16. Radial density profile as a function of time for 100 mtorr fill pressure.

100 mtorr

250 mtorr

400 mtorr

Figure 17. Restricted field-of-view interferograms at time of laser firing, $1.2 \mu\text{s}$. Vertical field of view is 2.0 cm.

Since the discharge tube wall is not visible in this field-of-view, an absolute density measurement is not possible. However, it is possible to make relative measurements of the density with respect to a standard radius. This allows comparison of the relative central density with and without laser firing. Moreover, the depth and radius of the on-axis density minimum are still measured absolutely. In fig. 18 the relative density on axis (with respect to a .93 cm radius) and absolute depth of the density minimum are displayed for 250 mtorr fill pressure. In addition, the absolute radius of the density minimum is also shown. The loss of a density minimum is seen to occur at $\sim 2.0 \mu\text{s}$ regardless of laser firing. The similarity of the relative density on axis in fig. 18 to the inferred density history for 250 mtorr fill pressure in fig. 15, supports the interpolation employed in fig. 15.

B. Refractive Containment

The refractive containment of the CO_2 laser beam by the on-axis density minimum was calculated from the measured, density profile values. Elaborate, ray-tracing solutions have been found for various analytic forms of density profiles.⁷³ However, a simplified calculation for a bounded, off-axis density maximum was desired. In particular, an estimate is made of the critical, incident angle, θ_c . Radiation with an incident angle from the axis greater than θ_c will escape the density minimum and leave the plasma column as sidescattered radiation. It is found that θ_c is greater than the maximum incident angle. Thus, the CO_2 laser beam was fully contained in the plasma column during the presence of the density minimum.

Figure 18. Relative density on axis and absolute depth of density minimum for 250 mtorr fill pressure. Radius of density minimum is also shown.

The geometry of the analysis is shown in fig. 19. A linear density profile is used from the axis, $r = 0$, to the density maximum, $r = r_{\min}$. The profile outside the density maximum is not considered since radiation reaching $r > r_{\min}$ is assumed lost from the system. The ray paths are considered from the focus point of the CO_2 laser (~ 10 cm into the theta-pinch). The incident angle of a ray is θ_0 . After some axial distance, z , a ray will have been deflected towards the axis by an angular amount $\Delta\theta$ or radial distance of Δr . The ray which has $\Delta\theta$ equal to its incidence angle at $r = r_{\min}$ defines θ_c .

The Eikonal equation for ray trajectory is⁷⁴

$$\frac{d}{ds} \left(\mu \frac{dr}{ds} \right) = \nabla \mu, \quad (\text{IV-2})$$

where s is along the ray and $\underline{r}$ describes a point on the ray from an arbitrary origin. Assuming cylindrical symmetry, (IV-2) becomes

$$\frac{d^2 r}{dz^2} = \frac{1}{\mu} \frac{d\mu}{dr}, \quad (\text{IV-3})$$

where only radial dependence for μ is considered; ie, $\mu = \mu(r)$. With $d\mu/dr = \text{constant}$, integration of (IV-3) yields

$$r = \frac{1}{2} \frac{d\mu}{dr} z^2 + z \tan \theta_0 + r_0 \quad (\text{IV-4})$$

$$\Delta r \approx \frac{1}{2} \frac{d\mu}{dr} z^2$$

where it is assumed $\mu \approx 1$ and $r_0 = 0$. $\Delta\theta$ is given by

Figure 19. The geometry of the simplified refractive containment calculation. The origin is located at the laser focus. Solid and dashed lines represent ray paths with and without refraction, respectively.

$$\begin{aligned}
\Delta\theta &= \theta_0 - \theta \\
&= \theta_0 - \tan^{-1} \frac{dr}{dz} \\
&\approx \theta_0 - \tan^{-1} \left(\frac{d\mu}{dr} z + \tan \theta_0 \right).
\end{aligned} \tag{IV-5}$$

Setting $\Delta\theta = \theta_0$ for $r = r_{\min}$ defines θ_c :

$$\begin{aligned}
\Delta\theta = \theta_0 &= \theta_0 - \tan^{-1} \left(\frac{d\mu}{dr} z + \tan \theta_0 \right) \\
\frac{d\mu}{dr} z &= -\tan \theta_0 = -\tan \left(\frac{r - \Delta r}{z} \right) \\
&= -\tan \left(\frac{r_{\min}}{z} - \frac{1}{2} \frac{d\mu}{dr} z \right).
\end{aligned} \tag{IV-6}$$

Solving (IV-6) for z , the position of maximum expansion, requires an iterative solution. θ_c is then given by (IV-5).

From (II-16), μ is given as

$$\begin{aligned}
\mu &= (1 - \omega_{pe}^2 / \omega^2)^{1/2} \\
&= (1 - n/n_c)^{1/2} \\
&\approx 1 - n/2n_c \\
&= 1. - 5 \times 10^{-20} n,
\end{aligned} \tag{IV-7}$$

where it is assumed $\omega \gg \omega_{pe}$ and $\lambda = 10.6 \mu\text{m}$. With μ given by (IV-7), the required parameters for calculation of θ_c are $r_{\min}$ and dn/dr . The mean values for these parameters are estimated from the holographic interferometry data. For $r_{\min} = 0.3 \text{ cm}$ and

$dn/dr = 3.3 \times 10^{16} \text{ cm}^{-4}$ (250 mtorr fill pressure), the solution of (IV-6) is $z = 29 \text{ cm}$ (from the front end of the coil) and $\theta_c = 1.8^\circ$. Since the maximum incident angle is only $\sim 0.9^\circ$, the CO_2 laser beam is fully contained in the density minimum.

As mentioned in the discussion of the interferometry data, the density minimum only persists for $\sim 1/2$ the laser pulse. Eq. (IV-4) was used with a density gradient of $6.7 \times 10^{16} \text{ cm}^{-4}$ in a calculation of the refractivity for an on-axis density maximum. It was found that the beam was refracted out to $r = 0.5 \text{ cm}$ at $z \approx 19 \text{ cm}$. Thus, the laser radiation is presumed totally lost as sidescatter during the second half of the laser pulse when an on-axis density maximum is present.

C. Plasma Expansion due to Heating

The data displayed in fig. 18 for 250 mtorr fill pressure shows no significant change in the density profile due to laser heating of the plasma. In particular, attention should be directed to the cluster of data points near $1.8 \mu\text{s}$ and $2.9 \mu\text{s}$ delay time. Regardless of laser firing, the density minimum has diffused away completely by $2.0 \mu\text{s}$, and the relative on-axis density at $\sim 3.0 \mu\text{s}$ is unchanged. The lack of plasma expansion is attributed to the low β of the plasma. β is defined as the ratio of local kinetic energy density to the external magnetic field energy density, $8\pi nT/B_0^2$. A brief derivation of density dependence on T follows which illustrates the effect of low β .

We begin with pressure balance, flux conservation and particle conservation:

$$B_G^2 = B^2 + 8\pi T$$

$$BA = \text{constant}$$

$$nA = \text{constant},$$

where B is the internal magnetic field, A is the area of the plasma's cross section and expansion is assumed only in the transverse direction. Elimination of B and A in these equations results in

$$\frac{B_G^2}{n^2} = 8\pi nT = \text{constant}, \quad \frac{B_G^2}{n^2 T^2} = 8\pi n^* T^2,$$

where the primed quantities represent the plasma after expansion. Using the definition of β and defining $X = T^*/T$ and $Y = n^*/n$, the last equation becomes

$$Y^2 = \frac{1/\beta + XY}{1/\beta + 1}.$$

Rearrangement of this equation yields

$$Y = \frac{1}{\beta X + (1-\beta)Y}.$$

For high β , $Y \approx 1/X$; i.e., the change in density is inversely proportional to the change in temperature. For low β , $Y \approx 1/Y$; i.e.,

$n^2 \approx n$ and no expansion occurs with heating. Experimental measurements give $X \approx 3$ and $\beta \approx 0.1$. For these values, eq. (IV-8) implies $Y \approx 0.90$. Thus, significant plasma expansion due to laser heating is not expected. The interferometric measurements are in agreement with this conclusion.

D. Exit Burn Patterns and Filamentation

Further information about the propagation of the laser beam through the plasma was provided by exit burn patterns. Exposed Polaroid 100 film was positioned 2 m past the theta-pinch coil. An intense, focused burn pattern was observed at 100 and 250 mtorr F111 pressure. The burn pattern occurred consistently at 100 mtorr. At 250 mtorr the pattern occurred intermittently. The burn pattern was localized to an area $\approx .5$ cm in diameter. The correspondence of the burn size to the size of the density minimum indicated refractive trapping or channeling of the laser beam in the density minimum. In fig. 20, examples of burn patterns are displayed for 100 and 250 mtorr F111 pressure.

Within the burn spot a regular, spatial modulation of the intensity pattern was often seen. The regular spacing of the intensity modulation strongly suggests filamentation scattering. The spacing in the modulation pattern ranged downward from 1 mm to $\approx .2$ mm. The power threshold for filamentation was given in Ch. 11. The required power for this range of filamentation spacing is 4×10^8 to 2×10^{10} w/cm². These powers are consistent with the peak power at the laser focus of $\approx 3 \times 10^{10}$ w/cm². Furthermore, the calculated growth rate for filamentation, $\approx 2.7 \times 10^8$ s⁻¹, was sufficiently short

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250 0101

V. ABSORPTION

A. Energy Inventory

1. Calorimetry

Three energy calorimeters (Scientech, Inc., # 36-0401, 4" diameter disc calorimeters) were used to measure the incident, reflected and transmitted laser energies. From these measurements estimates of the absorption by the plasma column are made. The position of the calorimeters to measure incident and transmitted energies are indicated in fig. 3. The reflected energy calorimeter, when in use, was positioned in front of the backscatter spectroscopy, diagnostic system. Seventy cm from the coil exit end, a 16.7 cm diameter wire screen was placed to reflect part of the transmitted CO_2 laser radiation into a calorimeter. The wire screen consisted of .254 mm, stainless steel wires placed vertically with 1.8 mm spacing. The reflectance of the screen was calibrated in the absence of plasma with a known incident laser energy. The transmitted energy measurement results for 100, 250 and 400 mtorr fill pressure are shown in fig. 21 for an incident energy of 150 J. Backscattered energy measurements gave ~ .7 J for the same incident energy. This backscattered energy can be attributed to stray machine reflection. The stimulated backscatter energy was only a small fraction of this amount.

Figure 21. Transmitted energy calorimetry measurements for an incident laser energy of 150 J.

2. Calculated Absorption

The net absorption of the plasma was evaluated through a simple energy inventory in the various energy channels the incident laser energy may follow. The energy inventory requires interpretation within the context of propagation and heating effects. These improvements in the evaluation of the absorption are considered in the next section.

In the upper part of Table IV, the energies in reflected, absorbed, transmitted and sidescattered radiation are tabulated for an incident energy of 150 J. For the simple evaluation of this section, the absorbed energy, E_{abs} , is calculated from IBA for the initial plasma parameters; ie, a constant absorption length is assumed. The energy lost to sidescatter radiation is neglected. With these approximations it is seen from Table IV that the energy inventory is incomplete; ie, the summed energies in reflection, absorption, etc., do not equal the incident energy.

B. Calculation Improvements

1. Sidescatter and Bleaching

The dominant mechanisms believed responsible for the discrepancies in energy inventory are sidescatter losses and bleached absorption (time dependent absorption length, $\bar{l}_{\text{abs}}$). Sidescatter losses affect the absorption evaluation most at low fill pressures while bleached absorption is important at high fill pressures.

Sidescatter and bleached absorption are simultaneously incorporated into a revised absorption calculation. Two-thirds of the incident energy is assumed lost to sidescatter. This sidescatter

TABLE IV
ENERGY INVENTORY AND ABSORPTIVITY

Energy Inventory

<u>P_{fill}</u>	<u>E_o</u>	<u>E_{cont}</u>	<u>E_{refl}</u>	<u>E_{tran}</u>	<u>E_{side}</u>	<u>E_{abs}</u>	<u>E_{tot}</u>
100	150.	150.	.7	40.	.0	27.	~68.
250	150.	150.	.8	14.	.0	125.	~140.
400	150.	150.	.9	6.	.0	149.	~156.

Revised Energy Inventory

<u>P_{fill}</u>	<u>E_o</u>	<u>E_{cont}</u>	<u>E_{refl}</u>	<u>E_{tran}</u>	<u>E_{side}</u>	<u>E_{abs}</u>	<u>E_{tot}</u>
100	150.	50.	.7	40.	100.	6.	~147.
250	150.	50.	.8	14.	100.	23.	~138.
400	150.	50.	.9	6.	100.	37.	~144.

Absorptivity

<u>P_{fill}</u>	<u>E_{abs}/E_o</u>	<u>E_{abs}/E_{cont}</u>
100	.04	.12
250	.15	.46
400	.25	.74

occurs during the latter 1/2 of the laser pulse when the on-axis density minimum has diffused away. The resulting "contained" energy, E_{cont} , is used directly in calculating a revised E_{abs} in a bleached absorption calculation. The formulation for bleached absorption is considered in detail in the next chapter. The essential result for the present discussion is⁷⁵

$$\left\{ \frac{2}{3} (1+R)^{3/5} + 2(1+R)^{1/5} + \ln \left[\frac{(1+R)^{1/5} - 1}{(1+R)^{1/5} + 1} \right] \right\} \frac{R(z)}{R_0} = -\kappa_0 z,$$

where $R(z)/R_0 = E(z)/E_0$, E_0 is the incident energy per unit area, $E(z)$ is the remaining laser energy per unit area for some z and κ_0 is the initial IBA absorption coefficient (ℓ_{abs}^{-1}).

2. Revised Absorptivity and Summary

A revised energy inventory incorporating sidescatter losses and bleached absorption is presented in the middle part of Table IV. Improved agreement between the summed energies and the incident energy is evident. Net fractional absorption is also shown for both the total incident energy and just the refractively contained energy.

The striking feature of Table IV is the high sidescatter loss. Though the estimate of sidescattered radiation is rough, it illustrates the critical need for good refractive containment of the laser beam over the entire plasma path. Refractive containment has typically not been a difficulty in other devices.⁷⁶⁻⁷⁹ These other experiments operated at higher β , allowing laser heating to sustain an on-axis density minimum. Furthermore, the Scylla I-C plasma

column is significantly longer. The laser beam traverses only a small amount of the plasma column before it is refractively lost due to an unfavorable density profile. One concludes that refractive containment is a greater problem with long, low β , plasma columns.

VI. HEATING

A. Spectroscopic Temperature Measurement

Spectroscopic analysis of impurity radiation for temperature measurement was chosen due to equipment availability and ease of diagnostics installation. A general review of spectroscopic diagnostic techniques is available by Elton.⁸⁰ The particular spectroscopic method emphasized in the present study was rate equation analysis of line emission histories.

The double monochrometer observation system employed is depicted in fig. 22. Impurity line emission was studied at two axial locations, 14 and 43 cm into the the coil. The monitored plasma light was imaged at a field aperture with dimensions 4 mm in the plasmas radial direction and 6 mm in the plasmas axial direction. The plasma light was then directed to two monochrometers with 50 μ slits and focal lengths of 0.3 m and 0.5 m. In order to stabilize the rather low, fluctuating, naturally occurring impurity level, a small amount ($\sim 0.5\%$) of CH_4 or N_2 was introduced during the static, D_2 filling of the discharge tube. The set of lines studied are listed in Table V. Carbon and nitrogen line histories, both with and without CO_2 laser firing, are displayed in fig. 23 and 24 ($z = 14$ cm). The effect of the CO_2 laser is seen as the ion stages are burned through with a resulting decrease in line intensity. The histories without CO_2 laser firing have the times of the peak line intensities indicated during the burn-through phase. Nitrogen data show similar histories with somewhat later peak intensity times due to higher ionization potentials. These peak intensity times are used

Figure 22. Spectroscopic diagnostic arrangement.

TABLE V
TIMES OF PEAK EMISSION FOR STUDIED IMPURITY LINES

<u>Ion</u>	<u>λ</u>	<u>E_{exc}</u>	<u>E</u>	<u>time at 100 mtorr</u>	<u>time at 250 mtorr</u>	<u>time at 400 mtorr</u>
				z = 14./43.	14./43.	14./43.
	($\text{\AA}$)	(eV)	(eV)	(μs)	(μs)	(μs)
C _{II}	6580.5	16.3	24.4	.2/.2	.45/.45	1.0/.9
C _{III}	4069.0	43.0	47.9	.4/—	.7/.6	1.2/—
C _{IV}	5801.3	39.7	64.5	.6/—	.95/—	
N _{III}	4379.1	42.5	47.4	.4/.4	.7/.6	1.1/1.1
N _{IV}	4057.8	53.2	77.5	.65/.5	1.4/1.0	
N _V	4603.7	59.2	97.9	1.3/.9		
with CO ₂ laser (at 1.2 μs)						
C _{IV}	5801.3	39.7	64.5	"/"	"/"	1.3/1.5
N _{IV}	4057.8	53.2	77.5	"/"	1.3/"	1.2/
N _V	4603.7	59.2	97.9	"/"	1.4/1.4	1.8/

Figure 23. C_{II} , C_{III} and C_{IV} line intensity traces for 250 mtorr fill pressure without and with CO_2 laser firing. Times of peak intensities, as used in the rate code, are indicated.

Figure 24. N_{III} , N_{IV} and N_V line intensity traces for 250 mtorr fill pressure without and with CO_2 laser firing. Times of peak intensities, as used in the rate code, are indicated.

in the rate equation analysis. A summary of peak intensity times is included with the list of studied lines in Table V.

The first attempt to spectroscopically measure an electron temperature used the method of relative line intensity measurements from subsequent ionization stages.⁸¹ The following discussion indicates the difficulties with this method and the reasons for pursuing a rate equation analysis of line emission histories. The results of the rate equation analysis are presented in the next section.

Local thermodynamic equilibrium (LTE) and Coronal equilibrium models provide contractible, analytical methods of calculating state populations in two extreme cases. LTE requires detailed balance of the dominating reactions. In optically thin, laboratory plasmas, detailed balance can only be expected in collisional reactions. Hence, LTE implies electron collisional de-excitation rates exceed radiative decay rates. Coronal equilibrium refers to the other extreme, wherein radiative decay dominates de-excitation processes. Between these regimes, corrected models (semi-coronal) or explicit rate calculations must be used. Fig. 25 shows these different equilibrium regimes for an ion with ionization potential of 100 eV.⁸² The parameter regimes of the present experiment are also shown in fig. 25. It can be seen that no analytical equilibrium model exists for the entire ion stage. However, as discussed by Wilson and others,⁸²⁻⁸⁴ a thermal limit in excitation energy exists within each ion stage. Below this limit radiative decay dominates depletion of the state. Above this limit, collisional transitions to the next higher ionization stage dominate depletion of the state.

Figure 25. Domains of validity for various equilibrium models for an ion with ionization potential of 100 eV. The parameter regimes for different experimental fill pressures are also shown.

Knowledge of the position of the thermal limit with respect to the studied lines may still allow an analytical description of relative line intensities using LTE and Coronal calculations.⁸⁵

Temporal and spatial variations in the electron temperature created further complications in the analysis of relative line intensity measurements. Temporal variations implied lack of equilibrium between successive ion stage ground state populations. Hence, relative emission measurements must be made from lines dependent on the same ground state population. Radial variations (during CO₂ laser firing) resulted in integrated emission measurements over different temperature regions. Further technical problems included a poor signal-to-noise ratio after laser firing and the need for careful intensity calibration of the detection system sensitivity. Due to these many difficulties, relative line emission measurements were superseded by numerical simulations of the measured appearance time of line radiation intensity peaks.

B. Rate Code for Spectroscopic Data Analysis

Due to radial temperature variations after laser firing and lack of equilibrium of ion ground state populations, a general solution of the ionization, recombination and excitation rate equations was used to find the population history of radiating levels. The ionization rate code used for this general solution, simulated the measured appearance time of line radiation intensity peaks for a given temperature history. Iterative corrections to the input temperature history matched the simulated and measured times of radiation peaks. Convergence of these iterative corrections results in an effective

measurement of T_e . A general solution of the rate equations requires neither complete LTE nor Coronal equilibrium to be valid and avoids the need for exact intensity calibrations. However, this analysis method is sensitive to calculated atomic rates and the number of measured intensity peaks.

Rather than calculating and summing all discrete transitions which affect the population of a particular excited level,⁸³ total rates for the population of a given ion stage are evaluated.^{86,87} The emission history of a particular line excited from the ground state is then calculated. Use of total rates is justifiable when state populations are dominated by transitions from the ground state. The general rate equations used are:

$$Q_z = \frac{n_{i,z}}{n_i}, \quad n_i(t) \propto n_e(t) \quad \text{ion density definition}$$

$$\sum_z Q_z = 1 \quad \text{normalization}$$

$$\frac{1}{n_e} \frac{dQ_z}{dt} = Q_{z-1}S_{z-1} - Q_z(S_z + \alpha_z) + Q_{z+1}\alpha_{z+1}$$

population balance (VI-1)

$$Q_z(t) = Q_z(t-dt) + dQ_z(dt) \quad \text{population increment}$$

$$I_{z,p}(t) \propto n_{i,z}(t)n_e(t)X_p(t) \propto Q_z(t)n_e^2(t)X_p(t)$$

emission dependence

$$X_p(t) \propto T_e(t)^{-1/2} \exp(-E_p/T_e(t)) \quad \text{excitation dependence,}$$

where Q_z is the relative ion density of stage z , S_z and α_z are the total ionization and recombination rates, respectively, from species

z , $I_{z,p}$ is the emission intensity from the excited level p , X_p is the excitation rate to state p from the ground state and E_p is the excitation energy to state p . Use of the relative density, Q_z , avoids a source term in the rate equations due to plasma compression. The excitation rate is from Seaton.⁸⁸

Interpretation of the equations for a volume integrated intensity may be made through closer examination of the relative emission intensity, $I_{z,p}$:

$$I_{z,p}(t) = \int_V I_{z,p}(\underline{x}, t) dV \\ \propto X_p(t) \left[\int_{\text{core}} Q_z n_e^2 dV + \int_{\text{wing}} Q_z n_e^2 dV \right].$$

The term labeled "core", representing the volume integral over the axial region, is weighted by the n_e^2 of the high density in this region. In other words, the strongest intensity is on-axis. Hence, the actual line intensity history is well approximated by considering only the on-axis density:

$$I_{z,p}(t) \propto X_p(t) Q_z(t, \text{axis}) n_e^2(t, \text{axis}).$$

The ionization rate is assumed due solely to collisional ionization as given by the Exchange Classical Impact Parameter (ECIP) model of Burgess and Summers:^{89,90}

$$S_z = 2.5 \times 10^8 \left(\frac{E_H}{T_e} \right)^{1/2} \sum_i \exp(-E^{(i)}/T_e) J^{(i)} \sum_{j=1}^N E_j^{(i)} C(E_j^{(i)}) W_j,$$

where,

i - shell of ground state of ion to be ionized

$E(i)$ - ionization energy of electron in i^{th} shell

$J(i)$ - number of equivalent electrons in i^{th} shell

C - ionization cross section

$E_j(i)$, W_j - normalized nodes and weights of an N -point

Gauss-Laguerre quadrature.

Comparison of ionization rate coefficient models with measured rates favor the ECIP model.⁹¹ Hence, ECIP was chosen over semi-empirical models due to Lotz or Seaton.

The recombination rate is calculated for two processes: radiative recombination and three body recombination. Dielectronic recombination is not considered important in the present parameter regimes (as supported by graphs in reference 89). Moreover, ionization rates dominate the ion populations in the transient, heated plasma described here. The radiative recombination rate is given by Seaton:^{87,92}

$$\alpha_z = 5.2 \times 10^{-14} z (E/T_e)^{1/2} [.429 + .5 \ln(E/T_e) + .469 (T_e/E)^{1/3}],$$

where z is the ion charge before recombination and E is the ionization energy after recombination. The three body recombination rate is given by Griem:⁹³

$$\alpha_z = 1.4 \times 10^{-31} n_e n_{th}^6 z^{-6} (E/T_e)^2 \exp[E/T_e (n_{th} + 1)^2],$$

where,

$$n_{th} = 1.26 \times 10^{22} z^{14/17} n_e^{-2/17} (E/T_e)^{-1/17} \exp[4E/17n_{th}^3 T_e]$$

is the primary quantum number representing the thermal limit within the ion stage.⁸²

A flow chart for the rate code is shown in fig. 26. The density history is taken from the full field-of-view interferometric data. Temperature guesses at the times of intensity peaks (typically 4-5 intensity peaks available) are input. The density and temperature histories are linearly interpolated from these input values for the number of time steps chosen in the code. Subroutines for ionization and recombination rates are called at each time step as the population and intensity histories are calculated for carbon, nitrogen and oxygen. The temperatures at the calculated times of intensity peaks are taken as improved temperature guesses for the input, measured times. Iterative convergence to a unique, temperature history giving the correct times of intensity peaks, supplies the analyzed temperature.

Use of the rate code with times of peak line intensities from Table V results in the temperature history shown in fig. 27 for 250 mtorr. The temperature history was extrapolated, where necessary, from the last data point prior to laser firing to the time of laser firing. The initial temperature at $t = 0$ was 1 eV. Fig. 28 displays the temperature at the two axial positions at times corresponding to before and after the first laser spike. A strong temperature gradient and reversal of gradient is apparent. Though some error exists in the temperature analysis, the temperature

Figure 26. Flow chart for rate code.

Figure 27. Temperature history inferred from rate code and times of peak intensities from Table 5 (250 mTorr).

Figure 28. Axial temperatures before and after first laser spike.

Points are analyzed temperatures. Lines indicate temperature gradient.

gradient reversal is established definitively through the line intensity histories (see fig. 27). The late N_{IV} peak at $z = 14$ cm for the case of no laser firing, indicates the end of the plasma column is colder than the middle. With a lower temperature at $z = 14$ cm established, the prompt N_V peak requires a higher final temperature at the end of the plasma column in the case of laser firing.

C. Heating Model

A simple heating model was developed to calculate expected temperature changes due to laser heating. This model was used to discern which mechanisms dominate local heating. The various processes considered in the model include IBA, bleaching, electron-ion equipartition, beam focusing, heat conduction and radiation.

The foundation of the heating model is simple inverse bremsstrahlung absorption, IBA. Assuming the absorption to be independent of time variations in the temperature requires a solution of the equations:

$$\frac{dI}{dz} = - \kappa_0 I \quad \text{power depletion,}$$

$$\frac{dT}{dt} = \frac{2}{3n} \frac{dI}{dz} \quad \text{electron heating.}$$

The solution is simply

$$I(z) = I_0 \exp(-\kappa_0 z)$$

$$T(z) = \int_{\tau} dT = \frac{2\kappa_0}{3n} E_0 \exp(-\kappa_0 z),$$

where $E_0 = \int_{\tau} I_0 dt$, κ_0 is the IBA coefficient ($\kappa_0^{-1} = \ell_{\text{abs}} \sim 10^{34} T_0^{3/2}/n^2$), is the laser pulse duration and ΔT is the total change in temperature due to the laser pulse.

Refinement of this model requires a solution incorporating a time dependent absorption coefficient; ie,

$$\kappa = \left(\frac{T_0}{T_e(t)} \right)^{3/2} \kappa_0 = \tilde{\ell}_{\text{abs}}^{-1}.$$

Varying absorption, due to temperature changes resulting from laser heating, is often referred to as bleached absorption. The equations to be solved, incorporating bleached absorption, are⁷⁵

$$\frac{1}{c} \frac{\partial I}{\partial t} + \frac{\partial I}{\partial z} = -\kappa_0 \left(\frac{T_0}{T} \right)^{3/2} I \quad (\text{VI-2})$$

$$\frac{3}{2} n \frac{\partial T}{\partial t} = \kappa_0 \left(\frac{T_0}{T} \right)^{3/2} I. \quad (\text{VI-3})$$

The solution to these equations is

$$\left\{ \frac{2}{3} (1+R)^{3/5} + 2(1+R)^{1/5} + \ln \left[\frac{(1+R)^{1/5-1}}{(1+R)^{1/5+1}} \right] \right\}_{R_0}^R = -\kappa_0 z$$

$$T(z) = T_0 [1+R(z)]^{2/5},$$

where $R(z) = (5/3)(\kappa_0 E(z)/nT_0)$ and $E(z) = \int_{\tau} I(z) dt$.

Equipartition between ions and electrons, variations in the beam area and heat conduction can be included in the above solution through a single, fractional, heating coefficient, α , introduced on the right side of eq. (VI-3). The coefficient, α , appears in the solutions of eqs. (VI-2) and (VI-3) as:

$$\left\{ \frac{2}{3}(1+\alpha R)^{3/5} + 2(1+\alpha R)^{1/5} + \ln \left[\frac{(1+\alpha R)^{1/5-1}}{(1+\alpha R)^{1/5+1}} \right] \right\}_{R_0}^R = - \kappa_0 z$$

$$T(z) = T_0 [1 + \alpha R(z)]^{2/5}.$$

It will be shown that reasonable choices of α lead to agreement of predicted heating with observations.

Equipartition of energy between electrons and ions was typically much faster than the length of the first laser pulse (heating calculations are made only for the first pulse since temperature data is only available for this time period). Hence, a factor of $\alpha = 1/2$ was introduced. Characteristic radial and axial heat conduction times were longer than the laser pulse. Hence, no change in the analysis was made for heat conduction. Characteristic equipartition and heat conduction times are shown in Table VI.⁹⁴⁻⁹⁶ (Heat conduction times are taken from the transport equation: $(3/2)n(\partial T/\partial t) = \kappa \nabla^2 T$, where κ represents the heat conductivity. A scale length of 0.1 and 10 cm is taken for $\perp$ and $\parallel$ conduction, respectively.) Beam propagation through the plasma column was characterized by an axially varying beam size, $A(z)$. The analysis above was normalized for maximum laser intensity at the laser focus region (beam radius = 0.1 cm). The beam radius increased to ~ 0.2 cm near the axial spectroscopic station at

TABLE VI
CHARACTERISTIC ENERGY LOSS TIMES

P_{fill}	<u>Heat Conduction</u>						
	$\omega_e \tau_{ei}$	$\omega_i \tau_{ii}$	$\kappa_{\perp}^e$ (10^{21})	$\kappa_{\perp}^i$ (10^{21})	$\kappa_{\parallel}^e$ (10^{23})	$\tau_{\perp}^{con}$ (ns)	$\tau_{\parallel}^{con}$ (μs)
250	8.5/12.	.1/.2	2.8/2.0	2.9/5.2	1.6/2.9	368./292.	131./
400	4.0/4.0	.06/.06	4.8/4.8	1.5/1.5	.89/.89	500./500.	353./

	<u>Equipartition</u>	<u>Radiation</u>	
	τ_{ei}^{eq} (ns)	τ_{ff} (μs)	τ_{bb} (ns)
250	59./80.	~500.	~150.
400	28./28.	~300.	~100.

$z = 43$ cm. An additional factor of $1/4$ was included therefore in α for analysis at this station. Radiation losses were significant, as seen by the characteristic energy loss times in Table VI.^{97,98} However, they constituted a constant energy sink balanced by adiabatic heating due to the magnetic compression of the theta-pinch. Laser heating is considered a separate, uncoupled mechanism from radiation cooling and adiabatic heating in the simple model above.

D. Summary

A comparison of the model predictions with the experimental measurements is presented in Table VII. The calculated temperatures for 250 and 400 mtorr are shown with various model refinements. Simple IBA is inadequate for estimating the temperature increase. Bleaching is observed as a strong correction to simple IBA. The last refinement case, bleaching with equipartition and beam expansion, gives heated temperatures which are in order-of-magnitude agreement with the measurements of 34 and 27 eV at the respective axial positions. The mechanism appearing most dominant in the model, beyond simple IBA, is bleached absorption. The ability of the simple model to approximately reproduce the measurements provides strong verification of the bleaching phenomena.

TABLE VII
COMPARISON OF HEATING MODEL PREDICTIONS
AND EXPERIMENTAL MEASUREMENTS

<u>Heating Model</u>	<u>250 mtorr</u>
	$z = 14./43. \text{ (cm)}$
T_e (measured at $t = 1.2 \mu s$)	13./17. (eV)
T_e (measured at $t = 1.4 \mu s$)	34./28. (eV)
Simple IBA	593./315. (eV)
Bleached IBA	95./95. (eV)
Bleached IBA	
& Equipartition ($\alpha = .5/.5$)	72./72. (eV)
Bleached IBA,	
Equipartition	
& Beam Expansion ($\alpha = .5/.13$)	72./41. (eV)

VII. CONCLUSIONS

Conclusions have been given throughout the text which pertain to the particular experimental results being discussed. These results and conclusions are now summarized.

Stimulated Brillouin Scatter:

1. SBS was observed with a signal characterized by a single pulse near the time of peak power.
2. The single pulse occurred in a time window of reduced axial temperature gradient.
3. SBS displayed the expected red-shift in frequency corresponding to the ion-acoustic frequency.
4. A theoretical spectral fit of the backscattered line shape agrees well with the measured temperature.
5. The spectral fit of the data and an absolute power estimate gave the SBS exponential gain value, n_{ef} (number of e-foldings), to be 14-24 above the initial noise level.
6. This experimentally determined n_{ef} corresponds to an SBS interaction length consistent with the high power density, focus region.
7. The constraints on SBS gain due to temperature gradients can account for the high thresholds and intermittent behavior of the SBS signal.
8. The SBS power was limited by a long gain length, low plasma noise level and short laser focus region.
9. More detailed temperature measurements are needed to quanti-

- tatively study SBS dependence on the temperature gradient.
10. Single-shot spectral data with greater resolution are needed to study the theoretical spectral dependence in greater detail.

The most important feature of these results is the low SBS gain. The SBS power was limited by the long gain length and short region (with respect to the total plasma length) in which threshold was exceeded. The SBS energy was apparently limited by the short time during which reduced axial temperature gradients allowed gain over a significant length.

Laser Beam Propagation:

1. Refractive containment of the CO₂ laser beam was observed at early times during the laser pulse ($< 2 \mu s$ into plasma discharge).
2. Channeling of the laser beam in the programmed plasma density minimum was inferred from exit burn patterns.
3. Refractive containment was lost at late times ($> 2 \mu s$ into plasma discharge).
4. Laser heating did not prevent diffusive loss of the density minimum.
5. Further study is needed to substantiate filamentation structure in the exit burn patterns.

Propagation through a long plasma was limited. Improved containment of the laser beam requires a density minimum sustained through further magnetic field programming.

Laser Beam Absorption and Plasma Heating:

1. Local absorption efficiency was high and consistent with classical theory for the portion of the laser beam contained in the plasma.
2. Net absorption efficiency was low due to a lack of refractive containment of the laser beam at late times during the laser pulse.
3. Temperature measurements at two axial stations showed substantial heating due to the initial laser spike.
4. Classical modeling of the expected heating is consistent with the temperature measurements.
5. Bleached absorption; ie, an absorption coefficient incorporating the dynamically changing temperature, must be included in the analysis to account for the observed heating.

These results confirm that classical heating and refraction dominated the laser-plasma interaction, and SBS caused only a minor change in heating efficiency.

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