Columbia University
in the City of New York

**A THOMSON SCATTERING STUDY OF
COLLISIONAL-SHOCK HEATED PLASMAS**

by

Donald H. McNeill

1973

Plasma Laboratory

School of Engineering and Applied Science

Columbia University

New York, N.Y. 10027

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School of Engineering and Applied Science

Columbia University

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ABSTRACT

A THOMSON SCATTERING STUDY OF COLLISIONAL-SHOCK HEATED PLASMAS

Donald Hubert McNeill

Thomson scattering was used to obtain electron temperature and density profiles in deuterium plasmas produced by strong collisional shock waves. The transverse shock speeds were 3 to 25 cm/ μ sec in 50 to 100 mTorr deuterium. Ruby laser light was incident at a right angle to the axis in the midpoint of the annular gap of an electromagnetically driven coaxial shock tube and scattered light was collected at 90 degrees. Plasma created by reflecting a strong shock from a dielectric wall was also studied. It is found that the plasma electron temperature (10-30 eV) is much less than the ion temperature (100-500 eV) as is predicted by theory for the early life of the plasma. Electron-ion Coulomb collisional temperature relaxation does not occur during the times of observation (3-4 μ sec). Electron heating in the plasma is caused primarily by Ohmic dissipation of currents in the shock structure and piston. Plasma densities ($0.5-3.5 \times 10^{16}$ cm⁻³) agree well with shock theory predictions. For gas ionizing shocks, good agreement with collisional theory is found for the electron temperature.

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CHAPTER I INTRODUCTION

This thesis presents the results of laser Thomson scattering measurements of electron temperature and density in the plasma produced by very strong collisional shock waves. Shock speeds of up to 10^8 cm/sec and a large volume ($\sim 3 \times 10^3$ cm³) of hot (ion temperature up to 10^7 °K) and dense ($n \sim 10^{16}$ cm⁻³) plasma have been obtained in the electromagnetically driven coaxial shock tube at Columbia University ^{1,2}. Shock formation, heating, and compression of plasmas is one of the major methods being considered for production of a controlled thermonuclear fusion reactor plasma ^{3,4}. A considerable amount of theoretical and experimental work has been carried out on shock waves in plasma physics ⁵. Because of the very small cross section for light scattering from free electrons, the use of light scattering as a plasma diagnostic became possible only after the development of high power ruby lasers ⁶. Laser scattering from laboratory plasmas is technically difficult because of the high laser light intensity compared to scattered light intensity but permits highly localized non-perturbing measurements of plasma properties.

The theory of collisional shock waves predicts preferential heating of ions as opposed to electrons in an ionizing shock wave. Following the initial shock heating,

equilibration of ion and electron temperatures should occur due to energy exchange between ions and electrons in Coulomb collisions. The experiment described in this thesis studies the electron heating in a shock heated plasma by making Thomson scattering measurements of electron temperature and density, scanning the entire lifetime of the plasma in the shock tube. A reflecting wall has also been mounted in the shock tube at a point 5 cm downstream of the laser scan point. This permits study of plasma which has been brought to rest by the wall and further heated by the wall-reflected shock wave. A schematic drawing of the shock tube is shown in figure 1.1, together with an x-t diagram of a shock wave, showing its position as a function of time inside the shock tube.

This experiment has shown that for fast transverse ionizing shocks the resulting plasma has relatively cold electrons ($T_e \ll T_i$) and that the electrons are initially heated by resistive dissipation due to currents at the shock front and at the magnetic piston. For gas shocks (no magnetic field) excellent agreement is obtained with theory, and the electron temperature relaxes to the expected values in the times and distances calculated.

The theory and the corresponding experimental results for collisional shock waves in plasma physics are discussed briefly in chapter II. The related aspects of Thomson scattering from laboratory plasmas are described.

The apparatus and associated components for this laser scattering experiment are discussed in chapter III. The operation of the shock tube is described and the overall scheme of the Thomson scattering experiment is outlined. Mechanical and electrical design of this experiment are discussed in detail, and the determination of the response of the polychromator and detectors is explained. In order to determine plasma electron density from Thomson scattering, it is necessary, as an initial calibration technique, to be able to carry out Rayleigh scattering on the system. Considerable effort in this experiment was devoted to attaining good Rayleigh scattering signals (compared to stray light levels) from a room temperature test gas in the shock tube. Chapter IV describes the analysis by which electron temperature, T_e , and number density, n_e , values are obtained from the raw data and discusses the limitations imposed on the use of Thomson scattering as a diagnostic in this experiment. These limitations are due to the high plasma background light levels obtained during shock tube operation. Tables 4.1 and 4.2 and figures 4.3 through 4.7 present the major results of this thesis for the T_e and n_e measurements obtained by means of Thomson scattering.

Chapter V interprets the results of the experimental measurements of T_e and n_e for this plasma. For the strong transverse ionizing shocks studied, electron temperatures are well below ion temperatures immediately

following the shock front and throughout the post shock plasma. On the other hand, for the slower gas shocks (zero magnetic field) and for the slow speed case of initial magnetic field of opposite polarity to the drive magnetic field, the theory and the experimental data agree quite well. Resistive (Joule) heating by the currents in the incident shock structure and in the magnetic piston is found to provide excellent agreement with observed electron temperatures in the case of transverse ionizing shocks. Joule heating in the shock structure is the primary means by which electrons are heated in this plasma since electron-ion temperature equilibration does not occur during the time of observation. Relationships between T_e and n_e and between magnetic field compression and plasma density compression are discussed. Finally, some special features of a few of the observed spectral profiles of Thomson scattered light in this experiment are examined. These special features may be due to omission of certain physical phenomena in the usual theoretical calculations of spectral distributions, or to correlation (ion scattering) effects, or to a possible non-Maxwellian (i.e., two electron temperature) velocity distribution in the plasma. The results of this experiment are summarized in the final section together with suggestions for some possible further investigations.

CHAPTER II THEORY AND PREVIOUS EXPERIMENTS

The theory of shock waves in plasma physics is extensively developed and several reviews have been published which describe shock wave theory and the corresponding experimental results ^{5,7,8}. This chapter presents the results of this theory, showing the expected properties of the shocked plasma and the reflected shock heated plasma. Some considerations of shock structure are included because structure related effects are found in some of the results of this experiment. In particular, collisional heating and ion-electron energy equilibration are considered. Particular emphasis is placed on effects which can be studied by Thomson scattering, from which values of T_e and n_e may be obtained. Since the development of the ruby laser, Thomson scattering has been used as a plasma diagnostic in a variety of laboratory conditions. It is not, however, a routine technique because of considerable experimental difficulties. Several reviews of laser scattering diagnostics for plasma studies have been written ^{6,9,10,11}. Only results are presented from the theory of Thomson scattering which are of possible use in studying the properties of plasmas produced in the coaxial shock tube.

2.1 SHOCK WAVES IN PLASMA PHYSICS

In nature, shock waves arise when a large amount of energy is suddenly released into a continuous medium. The disturbances involved are large. The small amplitude compression or sound wave of linearized gas dynamics becomes, for large amplitudes, a steep permanent gradient profile, or shock wave-- whose shape is determined by the balance of non-linear dispersion and dissipative effects. Shock waves occur in gases, liquids, and solids ¹², and analogs have been found in various electromagnetic phenomena ¹³.

In a gas, as the intensity of shock heating and compression increases, we go from sound waves (linear case), to gas dynamic shocks, to ionizing and plasma shocks, and finally to relativistic shocks. Gas dynamic shocks have been studied for over a century and the literature is still growing. The texts by Courant and Friedrichs ¹⁴ and by Zel'dovich and Raizer ¹² cover many of the aspects of gas dynamic shocks. Shock waves have become a topic of considerable interest in plasma physics as a means of producing and then heating and compressing a plasma to thermonuclear temperatures. Wide ranging reviews of shock waves in plasma physics have been published by Gross ^{7,8} and by Chu and Gross ⁵.

In the experiments in the Columbia University high energy electromagnetically driven coaxial shock tube, a

shock heated and compressed plasma is produced in deuterium. The mechanism of shock and driver (magnetic piston) in an electromagnetic coaxial shock tube is discussed by Meehan¹⁵. In this thesis we deal with ionizing shocks, so called because the preshock gas has electrical conductivity $\sigma_1 = 0$, and the post shock plasma has $\sigma_2 > 0$. This case is intermediate between a gas dynamic shock, for which $\sigma_1 = \sigma_2 = 0$, and a plasma (or MHD) shock, for which both σ_1 and $\sigma_2 > 0$. The ionizing shock propagates into an initial cold (300°K) low pressure (50-100 mTorr) gas. See figure 1.1 for a schematic illustration of the shock tube and the plasma in it. The ionizing shock may be considered as consisting of two stages. The first stage is a fast relaxation zone of a few mean free paths (in the initial gas) thickness, in which there is a sharp density jump and excitation of translational degrees of freedom of the gas. Viscosity and thermal conductivity are the dissipative processes at this stage. The second, much thicker stage, is the "slow" relaxation zone in which the excited translational motion relaxes to other degrees of freedom. The processes which occur in this zone, in order of increasing energy required to initiate them, are molecular rotation, oscillation, and dissociation, excitation of electronic levels, and ionization.

The structure of a plasma shock wave is determined basically by two features: the requirement of quasi-

neutrality and the great mass difference between electrons and ions. In the coordinate system fixed to a stationary shock front most of the directed motion of the incident gas is converted to thermal motion by viscous dissipation in the shock structure. Since the ion mass (M_1) is much greater than the electron mass (m_e) (For deuterium, $M_1/m_e \sim 3700$), a shock front of speed U (in the lab frame; in the shock rest frame the unshocked gas has speed U), such that U is much greater than the thermal speed of the initial gas, gives the ions a kinetic energy $E_1 \sim M_1 U^2$ and the electrons $E_e \sim m_e U^2$, where $E_1/E_e \sim M_1/m_e \gg 1$. Since temperature has the same mass and velocity dependence, $T_1 \gg T_e$ immediately after heating by the shock front. Electrons may be further heated in the shock by adiabatic compression, resistive heating, acceleration in the ions' electric field, or by elastic collisions with ions. Other mechanisms may heat the plasma, particularly in the non-equilibrium relaxation zone. Recently, considerable interest has developed in the theory of collisionless shocks, i.e., those whose thickness is much less than the mean free path of the preshock medium. Collisionless shock theory and experiments are described by several authors 5,16,17,18. Collisionless or collective field heating mechanisms generally heat electrons more strongly. However, some collisionless heating mechanisms, such as ion trapping in the potential troughs of magnetosonic waves may heat ions as well 19.

In a collisional plasma shock, Shafranov²⁰ distinguishes the cases of shocks in plasmas without a magnetic field and those with a magnetic field transverse to the direction of shock propagation. The shock without an initial magnetic field has the usual density jump and a region following the jump where equilibration of ion and electron temperatures occurs. At the leading edge of the shock structure there is a region of raised electron temperature due to electron thermal conductivity. For a strong transverse bias magnetic field the thermal conductivity is lowered by a factor $(\omega_{ce}\tau_e)^2$ (where ω_{ce} is the electron cyclotron frequency and τ_e is the collision period) which increases as B^2 . Thus, with the initial magnetic field the electron precursor is lost and only the post jump region remains.

In this experiment a precursor was indeed observed when no bias magnetic field was present. Photoionization of the initial gas by ultraviolet radiation from the shock heated gas or ionization by some earlier breakdown across the tube away from the launching wall are possible means by which an initial weakly ionized lower temperature plasma is produced in the shock tube. With a transverse bias field no precursor is observed. Equilibration of electron and ion temperatures should occur in a time of approximately $0.3\mu\text{sec}$ for the measured values of $T_1 \sim 300\text{ eV}$ ²¹ and early post shock values of $T_e \sim 15\text{ eV}$ (measured in this experiment), if Coulomb collisional

relaxation occurs under the post shock conditions. Electron-ion temperature equilibration is not observed during the time in which transverse ionizing shock plasmas were studied in this experiment (several microseconds). In chapter V the nature of the electron temperature and density profiles throughout the entire shock wave and post shock plasma is discussed.

At low shock wave speeds, i.e. $M \leq 5$ (where the Mach number $M \equiv U_s/a_1$, where U_s is the shock speed in the laboratory frame, and a_1 is the sound speed in the initial gas), the shock jump equations (conservation relations across the shock jump) give the post shock state (u_2, p_2, T_2, ρ_2) in terms of the initial state ($u_1 = U_s, p_1, T_1, \rho_1$) according to the Rankine-Hugoniot relations for a nonconducting gas ^{12,14}. Ionization occurs in hydrogen for $M \sim 20$ to 55. Above 55 ionization in hydrogen is complete for the range of initial pressures considered in this thesis ⁵. The shocks in this experiment have Mach numbers in the range 30 to 280 ($U_s \sim 2.7$ to 25×10^6 cm/sec in deuterium with initial pressures 50-100 mTorr). The Rankine-Hugoniot equations are then complicated by the addition of chemical (dissociation and ionization) effects and electromagnetic effects. The jump equations for a strong (ionizing) shock propagating into a stationary preshock gas with an initial magnetic field include equations for conservation of mass flux, momentum flux, and energy flux, an equation of state for pre-shock and post-

shock gases, Faraday's law and an equation for pressure balance across the magnetic piston. Appendix A lists these shock jump equations. Solution of these equations will predict the post shock plasma state in terms of initial parameters but these equations do not yield shock structure information. Because of the chemistry, for given conditions detailed numerical calculations must be carried out. Limiting values, as the shock speed becomes large, for the shock jump ratios of density, magnetic field, pressure, post shock gas speed, and temperature for a nonrelativistic ionizing shock are as follows (subscript 1 for incident gas; 2 for post shock plasma):

$$\rho_2/\rho_1 \rightarrow 4$$

$$B_2/B_1 \rightarrow 4$$

$$P_2/P_1 \rightarrow M^2$$

$$u_2/u_1 \rightarrow \frac{1}{4} \quad (u_1 = U_s \text{ and } u_2 \text{ are the velocities of incident and post shock gases in the shock front rest frame.})$$

$$T_2/T_1 \rightarrow \frac{1}{4}M^2.$$

The limiting ratios of reflected plasma quantities to incident post shock quantities for very strong incident shocks with no initial magnetic field are as follows⁸ (subscript 5 for reflected post shock values; 2 for incident post shock):

$$\rho_5/\rho_2 \rightarrow 2.50$$

$$P_5/P_2 \rightarrow 6$$

$$U_r/U_s \rightarrow 0.50 \quad (U_s \text{ is the incident shock speed})$$

$$T_5/T_2 \rightarrow 2.40.$$

speed; U_r is the
reflected shock speed.)

The simplest model of coaxial shock tube behavior is the "snowplow" model ²². This model assumes an impermeable magnetic piston pushing a hot plasma at the front of which is a strong shock wave. The rate of change of momentum of the plasma is equated to the magnetic piston pressure, and this equation gives the piston and essentially the shock front velocity as a function of time. In most of the experiments carried out for this thesis, the shock produced plasma strikes a reflecting wall which brings the plasma to rest, while in turn the shock front is reflected back into the plasma, heating and compressing it further. Analyses of plasma heating by reflected shock waves have been done by Taussig ²³ and a computer simulation study of curved wall shock focussing has been done by Lui ²⁴. An earlier simulation of coaxial shock tube operation, including dissipative effects in the jump equations, was carried out by Schneider ^{25,26,27}. Figure 2.1 shows single fluid temperature ($T_e + T_i$) as a function of shock speed in deuterium calculated for several shock tube fill pressures and initial magnetic fields. These curves were calculated by R.T. Taussig in connection with the experimental shock tube research at Columbia. The calculation uses a set of equations for one-dimensional plasma motion (does not include radial effects) and does not include shock structure. The graphs of figure 2.1

Figure 2.1 Single fluid temperature calculated for transverse ionizing shocks in deuterium using the "snowplow" model of shock propagation.

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also show the expected plasma temperature resulting from plasma heating by reflecting a shock wave from a wall.

The present experiment is part of a number of investigations of plasma shock waves which have been made at Columbia University in the past ten years. Transverse ionizing shock experiments on the present coaxial shock tube are reported in the literature ^{1,2,21,28-35}. Detailed experiments on a smaller coaxial shock tube were done by R. Meehan ¹⁵. He used spectroscopic methods to measure the plasma temperature and density integrated over a line of sight in the plasma. He obtained electron temperatures in the 1 to 10 eV range for much smaller shock speeds and plasma volumes than are available in the present apparatus.

2.2 PLASMA MEASUREMENTS BY THOMSON SCATTERING

2.2.a) Theory. The theoretical basis for scattering of electromagnetic waves from an ionized gas was developed around 1960³⁶⁻⁴¹. Before the development of the ruby laser no sufficiently intense light source existed to do scattering in the optical frequency range. Numerous experiments have now been carried out using the technique of scattering laser light from plasmas; however, the technique remains difficult. Scattering measurements have the advantage over probe (magnetic, electrostatic) measurements in that they do not significantly perturb the plasma. They are also essentially a point measurement (in contrast to line of sight) and thereby can provide spatial resolution of a plasma. Use of a pulsed laser provides time resolution. Several very informative reviews of laser light scattering in laboratory plasmas have been published^{6,9-11}.

The differential cross section for scattering from electron density fluctuations, per unit angular frequency interval $d\omega$ per unit solid angle $d\Omega$, is given by

$$\frac{d^2\sigma}{d\omega d\Omega} = \sigma_T S(\omega, \bar{k}), \text{ where } \sigma_T \equiv \frac{8\pi}{3} r_e^2 = \frac{8\pi}{3} \left(\frac{e^2}{mc^2} \right)^2 = 6.65 \times 10^{-25} \text{ cm}^2$$

is the Thomson cross section and $S(\omega, \bar{k})$ is the frequency spectrum of electron fluctuations with wave vector $\bar{k} = \bar{k}_1 - \bar{k}_s$ ($\bar{k}_1$ is the incident wave vector; $\bar{k}_s$, the scattered wave vector.). $S(\omega, \bar{k})$ has been calculated for a thermal

equilibrium plasma by Salpeter ³⁶ and for non-thermal plasmas by Rosenbluth and Rostoker ³⁹. The general form of $S(\omega, \bar{K})$ as derived by Salpeter ³⁶ is given in appendix B. The frequency spectrum depends on the parameter $\alpha = 1/K\lambda_D$, where $K = |\bar{K}_1 - \bar{K}_s|$ and λ_D is the plasma Debye length ($\lambda_D = 7.44 \times 10^2 \sqrt{\frac{T_e(\text{eV})}{n_e(\text{cm}^{-3})}} \text{ cm}$). α is roughly the ratio of incident wavelength to the plasma Debye length, or more precisely $\alpha = \frac{1}{K\lambda_D} = \frac{\lambda_i}{\lambda_D (4\pi \sin \theta/2)}$, where λ_i is the incident laser wavelength and θ is the scattering angle. We are chiefly interested in the domain $\alpha < 1$, which corresponds to uncorrelated scattering from the free electrons; that is, the laser wavelength is less than the Debye length and the ion thermal motion contribution to the scattered light spectrum is small.

The Thomson scattering cross section for a plane polarized wave from a single stationary electron is $\frac{d\sigma}{d\Omega} = \frac{3\sigma_T}{8\pi} \sin^2 \psi$, where ψ is the angle between the incident electric field vector and the scattered wave vector and σ_T is the Thomson cross section. The wavelength shift from the incident wavelength due to scattering by an electron with velocity component v parallel to the differential wave vector $\bar{K}_1 - \bar{K}_s$ is given by $\Delta\lambda = \frac{\lambda v}{c} \sin \frac{\theta}{2}$. This is a combined Doppler shift from the electron's seeing the incident light shifted due to its own motion and from the observer's seeing the scattered light emitted by a moving electron. Thomson scattering measurements of temperature evaluate the component of

thermal velocity parallel to $\bar{K} = \bar{K}_1 - \bar{K}_s$.

One might expect that if there were NV stationary electrons randomly distributed in a scattering volume V that the scattered power would be $NVS_1\sigma_T$, where S_1 is the incident energy flux. In fact, the scattered power from these motionless electrons would be zero. This is clear if one takes two slabs of electrons oriented perpendicular to the differential wave vector $\bar{K}_1 - \bar{K}_s$ and separated by a distance $\pi/|\bar{K}_1 - \bar{K}_s|$. Each slab would contribute (if the slabs have equal numbers of electrons) the same intensity of scattered light, but 180 degrees out of phase. Thus no scattered light would leave the volume. That scattering occurs at all is due to density fluctuations which produce an unbalance in the number of electrons in each slab⁴². The calculated spectral distribution of light scattered by a plasma in fact relates the scattered light spectrum to the spectrum of density fluctuations of electrons in the plasma. The total intensity of scattering from fluctuations of thermal electrons is proportional to the electron density, and it is this fact which permits laser scattering techniques to be employed to measure the electron density.

The relative contribution of electron and ion density fluctuations to the scattered spectrum is determined by $\alpha = 1/\kappa\lambda_D$, which for a ruby laser scattering at 90° to the incident beam is $\alpha \sim 1.08 \times 10^{-8} \sqrt{\frac{n_e(\text{cm}^{-3})}{T_e(\text{eV})}}$. Electron densities from $.2-4.0 \times 10^{16} \text{ cm}^{-3}$ and electron temperatures

from 2 to 50 eV were measured verifying that indeed

$\alpha < 1$. The scattered spectrum is thus Gaussian and is due to electron density thermal fluctuations. The expected half width of the scattered spectrum is $\Delta\lambda_{\frac{1}{2}} =$

$$4 \sin \theta/2 \sqrt{\frac{2kT_e}{m_e c^2} \ln 2} \lambda; \text{ which becomes}$$

$$\Delta\lambda_{\frac{1}{2}} = 32.35 \sqrt{T_e (\text{eV})} \text{ \AA}$$

for 90° scattering of ruby laser light.

Several phenomena can occur in plasmas which may affect the observed spectral distribution of Thomson scattered light. We deal with these briefly and consider their relation to our case.

(1) Laser heating of plasma. For $\tau_{ee} < t_L < \tau_{ie}$, where t_L is the laser pulse duration and τ_{ee} and τ_{ie} are the electron-electron and ion-electron collision times, the absorbed photon energy remains among the electrons alone during the time of the laser pulse and the increment in T_e due to inverse bremsstrahlung absorption of laser light is given by the formula ⁶

$$\frac{\Delta(kT_e)}{kT_e} = 1.42 \times 10^{25} \frac{n_i Z^2}{(kT_e)^{3/2}} \frac{g_{ff}}{\nu^3} \left\{ 1 - e^{-h\nu/kT} \right\} \frac{\int_0^{t_L} L dt}{g},$$

where kT_e is in eV, $g_{ff} \sim 1$, n_i is the ion density in cm^{-3} , ν in sec^{-1} , and $\frac{1}{g} \int_0^{t_L} L dt$ J/cm² is the laser beam energy per unit area. For this experiment, the beam energy is 600 J/cm², $n_i \sim 2 \times 10^{16}$, and $kT_e \sim 10 \text{ eV}$; therefore,

$$\frac{\Delta(kT_e)}{kT_e} \sim .01. \text{ Thus, the relative temperature increase}$$

due to absorption of laser light is negligible for these experiments.

(ii) Relativistic effects. Thomson scattering becomes Compton scattering and the cross section is described by the Klein-Nishina formula for incident photon energies comparable to the electron rest mass energy. In this case $(h\nu)_{inc} \sim 1.8$ eV, while $m_e c^2 \sim .51$ MeV.

We also note that the usual derivations of scattered profiles assume that the electron thermal velocities are such that $v_e/c \ll 1$. Sheffield⁴³ has noted that even for temperatures of 900 eV an error of about 100 eV is incurred by assuming $v_e/c = 0$. At $T_e = 100$ eV, the error is about 5 percent for 90° scattering. The error arises from a shift of the scattered profile to the blue. In section 5.4 we discuss the possibility that this shift has been observed in this experiment.

(iii) Effect of impurities. The presence of electrons from highly ionized impurities enhances the total amount of scattering by the added impurity electron density. The effect of the higher Z positive charge centers changes the shape of the ion contribution but no change in $S(\omega)$ occurs due to these electrons⁴⁴.

(iv) Magnetic field. A magnetic field affects the shape of the scattered spectrum, but not the total intensity. For $\vec{k} \perp \vec{B}$, the spectrum for $\alpha \ll 1$ consists of a series of lines of width $k v_e$ and spacing $\omega_{ce} = eB/mc$ with a Gaussian envelope.⁹ The spacing for a 10 kGauss field is less than $0.5 \text{ \AA}$ and is therefore unresolvable

by this kind of optical system. Thomson scattering has been proposed as a means to determine the direction of the magnetic field in a Tokamak plasma ⁴⁵.

(v) Coulomb collisions will affect the appearance of the spectrum whenever the width of the fine structure being studied is comparable to the collision period. In our case the collision frequencies are much higher than any observed fine structure; thus, no collision effects appear in our data.⁹

(vi) Non-Maxwellian distributions of velocity cause a change in the spectral shape for both ion and electron components ^{46,9}. We consider two such deviations in section 5.4: (1) an electron drift velocity relative to the ions, or a bulk plasma drift of both ions and electrons, and (2) a two electron temperature plasma.

2.2.b) Thomson scattering experiments. The first Thomson scattering of laser light from electrons was observed by Fiocco and Thompson ⁴⁷, who scattered ruby laser light from an electron beam. They observed a shift of scattered light with respect to the incident wavelength of 259 Å ($n_e \sim 5 \times 10^9 \text{ cm}^{-3}$ and electron energy 2 keV) and counted about 4 photoelectrons per pulse. Following this proof that light scattered by electrons could be detected, many experiments have been carried out ⁹⁻¹¹.

Thomson scattering studies have been made of normal ionizing shocks (initial magnetic field parallel to the direction of shock propagation) by Niland and others ^{48,49}. They have made T_e and n_e profiles of the region behind relatively slow normal ionizing shock waves in hydrogen and helium (U_s up to 10 cm/ μ sec).

The only Thomson scattering measurements carried out before the present experiment on a coaxial electromagnetically driven shock tube of the type at Columbia are those of Patrick ⁵⁰. He used a 30-100 MW pulsed ruby laser and measured electron densities of 8 to $16 \times 10^{15} \text{ cm}^{-3}$ for filling pressures of 50 to 100 mTorr of hydrogen. No initial magnetic field was applied. Single point per shock speed electron temperatures were determined using a two channel detector with rotatable interference filters.

A plot of T_e as a function of shock

speed was made. T_e values of from 30 to 80 eV were found for shock speeds from 31 to 47 cm/ μ sec in hydrogen. Profiles of T_e and n_e across the shock structure were not made, and the errors in T_e values were large.

CHAPTER III APPARATUS AND TECHNIQUE

This chapter describes the experimental apparatus used to obtain Thomson scattering data on the shock heated and compressed plasma. A brief description of the operation of the shock tube is followed by an overall description of the Thomson scattering experiment. Subsequently, the mechanical, optical, and electronic design of the apparatus is described. Finally, the characteristics of the system as a whole are discussed. Of particular interest are the system response and discrimination and its behavior when Rayleigh scattering measurements were made from cold (room temperature) argon and nitrogen in the shock tube. This is a significant aspect of this work since in practice Thomson scattering is a rather difficult technique, particularly for determining the electron density of a plasma.

3.1 SHOCK TUBE OPERATION-- PRODUCTION OF SHOCK WAVES AND PLASMA

The plasma studied in this experiment is produced in a high energy coaxial electromagnetically driven shock tube. Details of the Columbia University High Energy Shock Tube have been described previously by several authors^{2,28-31}. The operational characteristics of this shock tube, such as shock speed and magnetic field behavior, were first described by P.R. Moriette^{28,29} and

E. Halmoy ^{30,31}. Very detailed observations and studies of the physical and electrical processes at work ~~in~~ an earlier smaller coaxial electromagnetically driven shock tube at Columbia University were made by R.J. Meehan ¹⁵. Other details on design and operation of electromagnetic shock tubes may be found in the paper by Gross and Miller ⁵² and in the book by Makarov and Chekalin ⁵³.

The length of the shock tube is approximately 3 meters. The diameter of the inner conductor is 13 cm and of the outer conductor, 23 cm, leaving an annular gap of 5 cm. By the time the shock has progressed to a position about 125 cm away from the launching end of the shock tube, it has reached a nearly steady speed and the wave has a well defined structure. In the experiments reported upon in this thesis the measurements were made 175 cm from the launching wall at a point roughly half-way between the inner and outer conductors. Most measurements of the plasma electron temperature and density were made when the device contained a boron nitride reflecting wall, located 180 cm from the launching wall, so that studies could be made of the incident and reflected shock heated plasma at the 175 cm point.

Since 1971, when Halmoy and Moriette completed their work on the shock tube, only a few changes have been made in the construction of the shock tube. A new switching system for the high voltage (main) condenser bank was installed in late 1971, tested thoroughly, and found to be

satisfactory. These switches were developed at Columbia by R.S. Post and Y.G. Chen⁵⁴.

The main shock producing condenser bank consists of ninety-eight condensers connected in parallel and divided into six modules, with a total capacitance of 52 microfarads. The half ringing period of this bank with the shock tube load is approximately 10 μ sec. These condensers can be charged to 40 kV without switch self-breakdown at the usual switch gap settings and with air filling in the gaps. For this lower voltage range the bank is usually charged to 35 kV. When the switches are filled with atmospheric pressure SF_6 and the master switch is filled with dry nitrogen, the bank may be charged to 90 kV or more. The usual charging voltage in this case is 75 kV with occasional tube firings at 85 kV. The bank is discharged through six (one for each of six modules) non-uniform field distortion switches which are triggered simultaneously by a master switch.⁵⁴

In addition, an initial magnetic field may be produced transverse to the direction of shock propagation. For this a bias bank consisting of forty-six 20 kV condensers of 15 microfarads apiece is used. The bias circuit, for which the typical firing voltage is 10 kV, has a quarter ringing period of 65 μ sec. From the time of bias bank firing, a delay until firing the main bank of 12 microseconds gives a bias field of 1 kGauss, and a delay

of 65 μ sec gives 3.6 kGauss. In this experiment initial magnetic fields of 0, 1, and 3.6 kGauss were used.

All of the plasma experiments reported upon in this thesis used deuterium as a fill gas. Fill pressures of 40 to 100 mTorr were used as measured on an alpha-ray emission ionization gauge (NRC alphasatron). Higher filling pressures than these result in very low shock speeds and consequently cooler and less interesting plasma, while lower pressures give diffuse shock structures indicating poor definition of shock and post shock regions.

3.2 THE THOMSON SCATTERING EXPERIMENT

Prior to discussing the detailed components of the laser scattering experiment, a general outline of the apparatus is presented. Figure 3.1 shows a sketch of the laser, beam stops, beam focussing lens, shock tube, entry and exit ports, and beam dump. This figure also indicates the relative sizes and positions of the main equipment used. Figure 3.2 shows photographs of the experimental apparatus on the shock tube and of the laser. Figure 3.3 shows a vertical cross sectional view of the shock tube, together with the scattered light collection port and beam entry and exit points on the shock tube. An insert in this figure shows a diagram of the wave vectors corresponding to the incident and scattered light and to their difference. The differential wave vector thus shown lies in the direction of the component of

LASER THOMSON SCATTERING EXPERIMENT ON SHOCK TUBE (NOT TO SCALE)

(Top) View of apparatus on shock tube. From left: beam stop, lens, entry window, monitor light pipe, beam entry tube, shock tube, collecting optics, and beam dump. In background: McPherson 1 meter spectrograph.

(Bottom) Holobeam 3081 laser.

Figure 3.2

CROSS SECTION OF SHOCK TUBE AT SCAN POSITION FOR THOMSON SCATTERING

Figure 3.3

thermal fluctuations which is observed in this experiment (see section 2.2). Figure 3.4 shows the timing sequence for triggering the various discharges during an experimental run. The laser flashlamps are fired first. There should be a delay of about $950\mu\text{sec}$ to get optimal laser pulse energy and duration. Approximately $880\mu\text{sec}$ after the flashlamp monitor pulse occurs (indicating that the laser flashlamps have begun to discharge), the bias bank is fired. Following this by 0, 12, or $65\mu\text{sec}$ the main (shock producing) bank fires. The azimuthal magnetic field probe at the 150 cm position of the shock tube emits a large, fast pulse when the magnetic front of the shock arrives at the 150 cm position. This pulse is used to trigger a pulse delay generator and to trigger oscilloscope sweeps. The pulse delay generator, after a specified delay, sends a pulse to the grid of the high voltage thyatron to activate the Pockels cell Q-switch of the laser. After a delay of about 300 nsec a ruby laser pulse of up to ~ 7 Joules and ~ 20 nsec full width half maximum enters the shock tube perpendicular to the direction of shock propagation. The 90° (upward) scattered light from this beam is collected and sent to a six channel polychromator and photomultipliers. The resulting photomultiplier signals are recorded on oscilloscopes in a screen room.

TIMING SEQUENCE FOR THOMSON SCATTERING EXPERIMENT

Figure 3.4 The total delay from laser flash-lamp triggering to laser output pulse is $950 \pm 5 \mu\text{sec}$.

3.3 APPARATUS-- MECHANICAL AND ELECTRICAL DESIGN

3.3.a) Laser. The necessity for a pulsed laser for Thomson scattering arises from the need for a monochromatic, very intense source of short time duration. Observation of plasma light levels through ports in the side of the shock tube indicated a need for a laser of greater than 200 MW full power in order to achieve acceptable signal to plasma light levels. The laser chosen for this purpose is a Pockels cell Q-switched oscillator-amplifier combination ruby laser, the model 3081 manufactured by Holobeam. This laser has, at its peak power of 500 MW, a minimum energy of 7 Joules and a pulse width (full width half maximum) of 15 nsec, with a beam divergence of 3 to 5 milliradians. In practice, the laser was operated at slightly less than full power, and it had a pulse width of less than 25 nsec. Laser power measurements indicate 400 MW at the flashlamp bank voltages generally used in this work. Burn pattern measurements (on developed polaroid film) indicate a beam divergence of 2 mrad. The laser flashlamps are cooled and the cavity held at a constant temperature by a circulating water bath maintained at a constant temperature of 35° C. The oscillator (maximum power 100 MW) has a helical flashlamp (3.00 to 3.60 kV) which pumps a ruby rod of 3/8 inch diameter and 3 inches length. The output from the oscillator is horizontally polarized and hence the

amplifier output is horizontally polarized. The oscillator output enters a 1:2 upcollimator which is between the oscillator and amplifier sections. The beam then enters the amplifier rod (8 inches long, $3/4$ inch diameter) which is also pumped by a helical flashlamp (7.0 to 8.5 kV). The entry face of the amplifier rod is flat, perpendicular to the rod axis and is antireflection coated. The exit end of the amplifier ruby rod is cut at a Brewster angle to prevent front surface reflections from the amplifier going back into the oscillator. The beam thus leaves the amplifier in a direction about 30° off the laser axis. The beam has an elliptical cross section ($3/4$ inch major diameter, $3/8$ inch minor diameter). This off axis character prevents precise alignment of the ruby laser by means of a small cw helium-neon laser because of the difference in index of refraction of the ruby rod for the He-Ne laser's $6328 \text{ \AA}$ wavelength and for the ruby laser's $6943 \text{ \AA}$. Final alignment of beam direction requires successive firings of the laser to ensure correct positioning. For this reason the laser optical rail was positioned by means of four vertically adjustable screws, and once aimed properly, it was clamped firmly in place. Further alignment of the focal spot of the laser beam in the shock tube was carried out as described later.

In order to vary the time of the laser pulse to

correspond with the arrival of a particular element of plasma at the 175 cm position of the shock tube, it is necessary to use a remote delay which senses some earlier aspect of the shock tube plasma behavior. A reasonably precise method is to use the direct output of a magnetic field probe upstream of the point where measurements are made. This allows for internal electronic delays and permits scanning of all parts of the shock wave and post shock plasma. The magnetic field probe gives a sharp spike of 100-200 nsec rise time and 50-100 V, well above electrical noise levels and adequate to trigger a pulse delay generator. The pulse delay generator provides a 0-100 μ sec delayed pulse (in steps of .01 μ sec) of 50-600 V and 10 nsec rise time. A 350 V pulse from this pulse delay generator was sent to the grid of the KU-27 thyatron in the Pockels cell pulser, bypassing all but the power supply function of the Pockels cell remote control unit. The other delays involved, from flashlamp triggering to bias and main bank firings, were fixed in accordance with the delay chosen to produce the desired bias field. A variation of $\pm 50 \mu$ sec from the optimal delay of 950 μ sec for Q-switching causes a reduction of ~ 5 percent in laser energy output. The only variable part of this timing sequence is caused by changes in the shock propagation speed from shot to shot. The shot to shot variation in the total time from flashlamp firing to laser pulsing is not much greater than 5 μ sec. The

reproducibility of, laser pulse energy and time duration is excellent, with less than 5% variation observed.

All laser functions, such as charging the oscillator and amplifier condenser banks, were arranged so as to be controlled from the shock tube control panel. The laser power supplies and controls were well grounded and electrical noise from the shock tube's firing never produced difficulty with early triggering of the Pockels cell.

3.3.b) Beam entry and exit. Because of the very high power of the incident laser beam compared to the Thomson scattered signal, it is of vital importance to reduce as much as possible the amount of light of the laser wavelength which reaches the detectors of scattered light but is not Thomson scattered light. Such unwanted light, which originates from reflection off plasma vessel walls or from scattering in lenses and windows, is called stray light or parasitic scatter. The reduction of this light to low levels is among the most important design considerations for this type of experiment. The relative magnitude of incident energy to detectable signal for this experiment is about 10^{12} to 1. A typical incident beam energy is 6 Joules and a maximal detected signal in one channel is 10^{-14} J. Or, we can say that out of the total of 2×10^{19} incident photons, the scattering volume (viewed by the collecting lens-- about 5 mm^3) will scatter a total of 4×10^{10} photons

over all directions, so that the $2\pi \times 10^{-3}$ sr collection cone of the detection system will pick up about one in 10^{12} of the incident photons. Further loss of detectable signal occurs in the succeeding optics but the stray light is attenuated by an equal factor.

Difficulties in reducing stray light in this system were enhanced by the initial choice of small ($\frac{1}{4}$ inch) beam entry and exit ports in the shock tube (so as to avoid disruption of current patterns by large holes in the walls of the tube) and by the fact that the shock tube walls and structure are of metal. That center line (laser wavelength) Thomson scattering signals could be detected at all is a good achievement. In all data taken when Thomson scattering was observed, the center channel Thomson scattering signal was detectable with a magnitude between $1/20$ and 1 times the stray light signal. As a consequence it was possible to determine plasma electron densities in this experiment.

The natural beam divergence of the laser (2 mrad) can cause light to be incident on the focussing lens at a significant distance from the beam axis, so that this light will not focus properly or may strike containment vessel walls. To limit such off axis light two optical stops were made ($3/4 \times 3/8$ inch elliptical holes in Al sheet) and placed at distances of 50 cm and 250 cm in front of the laser output. These reduced the apparent

beam divergence to $\sim .5$ mrad at the lens. The beam focusing lens is a planoconvex lens of (measured) focal length 62 cm at 6943 Å. This focal length was chosen so as to have a beam major diameter of approximately $1/8$ inch at the point where the beam enters the shock tube wall (small enough to avoid grazing the tube walls and large enough to avoid air breakdown by having the beam too narrowly focussed outside the shock tube). The beam focussing lens is mounted on an X-Y positioner whose normal is almost, but not quite, parallel to the laser beam, thereby avoiding reflecting the beam back to the laser. As it has a long focal length the focal point position of the lens is not strongly sensitive to the X-Y position. Hence a fine adjustment enabled very precise alignment. 14 cm beyond the lens is a polished quartz entry window ($1/4$ inch thick) mounted at the Brewster angle (55°) to the incoming beam so that the horizontally polarized laser beam is almost completely transmitted. One end of a light pipe was located facing the entry window so as to send a light signal upon laser firing into the screen room where it was fed into a Monsanto MD-1 silicon PIN photodiode (response time less than 1 nsec). The output pulse of the photodiode serves as a laser timing and relative power monitor.

Upon passing through the entry window the laser beam enters a 1.6 cm diameter cylindrical entry tube which is at shock tube pressure. A $1/4$ inch diameter

opening into the shock tube is at the other end of the entry tube. This entry tube has a retractable cover flap used to shield the window from sputtered particles. It also has a vacuum seal valve so that the window may be removed and cleaned while the shock tube is held at vacuum. The beam then enters the shock tube annular plasma region (see figure 3.3) crossing 16.5 cm to the $\frac{1}{4}$ inch diameter beam exit hole in the shock tube wall. The focal spot is at the midpoint of this transit and has diameter less than 1.3 mm.

The beam leaves the shock tube section through a cylindrical tube similar to but shorter than the entrance tube and goes into a beam dump. Since the main power of the beam has not been absorbed or scattered, it must be disposed of outside the scattering area if it is not to be reflected back into the plasma vessel and thereby swamp all plasma scattered signal by its subsequent reflections off tube walls. A considerable amount of time was devoted to this problem and a fairly simple optical dump was developed. It was found to be sufficient to terminate the beam through a 1 inch O.D. glass tube, 4 inches long, flared at the end to $1\frac{1}{4}$ inches I.D. and cut at a Brewster angle to the axis, and attached with epoxy to a piece of one half inch thick black absorbing glass. The glass tube and black glass were then wrapped loosely in several layers of black rubber backed cloth to absorb reflected light. Use of this angled black glass

absorber as a beam dump was sufficiently effective to permit use of Rayleigh scattering calibration of the optical apparatus for evaluation of the plasma electron density.

A simple but effective alignment procedure for the laser beam inside the shock tube was developed. As mentioned earlier, the laser beam was aimed so its center portion passed through the tube entry and exit ports. A simple procedure that was used to obtain final alignment was to remove the beam dump, put in a clear glass exit window, and then look by eye back from that point to the laser amplifier rod face. Without the beam focussing lens in this optical path the ruby crystal shows clearly through the two aperture stops. With the lens inserted, any off positioning that results could be compensated for by adjusting the X-Y positioner on which the lens is mounted, until the crystal face is again centered. Once the crystal face was again in the center of the field of view, a test firing of the laser was made to ensure that the exit beam burn pattern was a small ellipse. It was found that a clearcut burn pattern always resulted from this alignment-- there were no streaks or "stars" in the burn pattern as might result if the beam glanced off a metal surface. The beam was thus focussed in the middle of the scattered light viewing port field of vision.

3.3.c) Scattered Light Collection, Dispersion, and Detection. This experiment measured the Thomson Scattered light at 90° to the incident beam. To collect the scattered photons a hole was made in the top of the outer conductor of the shock tube at the 175 cm position. See figure 3.3 for details. A brass fitting with a quarter inch hole in its center was fitted into this hole in the shock tube. It served both as a collimator and as the frame for a quartz covering window, the collecting lens, and the light pipe to the polychromator. To prevent the quartz window acquiring a thin coating of vapor deposited metal when experiments other than the laser scattering were being performed, a ceramic shutter was installed in front of the quartz window. From the top of the quartz window to the center of the focussed beam the distance is 7.4 cm. Above the window a 1.75 cm focal length double convex lens was fixed so as to focus an image of the scattering volume approximately 2 cm above the lens. At this image position was placed a pinhole (diameter 0.6 mm) in a black painted aluminum sheet. Without a pinhole no Thomson scattering signals could be observed, even though preliminary calculations (see appendix C) had indicated that they should be observable. The plasma volume viewed by the full quarter inch aperture of the collimator was large enough that its light emission swamped the Thomson scattered light. The pinhole sheet, together with the circular

end of a light pipe (This 5 foot long light pipe, with a $\frac{1}{4}$ inch diameter on this end, leads to the polychromator entry port where the light pipe end "fishtails" to a rectangle of dimensions 1.5x20.5 mm) was mounted on an X-Y positioner which was attached to the brass collimator fitting. Alignment of the horizontal position of the pinhole was necessary whenever laser beam re-alignment was performed. To aid in this alignment of the collecting optics, a quarter inch rod with an angle cut face painted black was placed in the laser beam path so that the angled face was in view of the collecting lens. The laser was then fired, and the burnt face was then illuminated by a light aimed along the laser beam path. With the focal spot burn mark visible it is only necessary to move the pinhole so that the burn mark is visible through it. The light pipe is then replaced, and optical alignment is complete.

The fishtail end of the five foot long light pipe was inserted into the entrance slit of a 1 meter McPherson spectrograph, used as a polychromator for dispersing the Thomson scattered light from the plasma. This spectrograph has a plane reflecting grating of 600 lines/mm and 10x10 cm² area in a Czerny-Turner arrangement with 1 meter focal length concave mirrors between the entrance and exit slits and the grating. The linear dispersion of this instrument is 15 Å/mm.

The entrance F number of the spectrograph is $f/10$, while the effective F number of a standard flat polished light pipe was measured to be about $f/2.5$. Hence, there is a loss of considerable light in this arrangement. However, tests showed that plasma light levels alone were sufficient to ensure easily detectable signals despite this inefficient optical coupling. The exit of the McPherson spectrograph is a slot 2 cm high by 21 cm long with constant dispersion over this length. The detector chassis containing six photomultiplier tubes is attached to the inside of the spectrograph cabinet facing this exit slot.

The dispersed light from the spectrograph was sent into a six channel detector mounted inside the spectrograph cabinet. This detector consisted of six $1/16$ inch thick plexiglass strips 5 to 7 inches long and $3/4$ inch wide, separated with butyl rubber inserts of suitable thickness between channels at the exit slot of the spectrograph. The plexiglass strips lead to six RCA type C31025C photomultipliers. Figure 3.5 shows a front and back view of the chassis upon which the light guides and photomultipliers are mounted. The plexiglass strips were cut to size and the edges polished with emery and jeweler's rouge to a smooth reflecting surface. Additional treatment of these light guides was really unnecessary, since painting them with white

Front and back views of photomultiplier chassis.
Baffles between the plexiglass strips are not
shown.

Figure 3.5

reflecting paint or further polishing of edges was found to give little transmission improvement. The nine dynode photomultipliers were each wired with a $10 \times 100 \text{ k}\Omega$ divider chain⁵⁵, and all six were supplied in parallel by a 0-1500 V, 0-10 mA regulated DC power supply set at 1250 V. See figure 3.6 for the divider chain wiring diagram.

The RCA type C31025C (a developmental tube, shortly to appear as type 4832) is a GaAs photocathode photomultiplier. This photocathode has unusually high sensitivity in the range 5000 to 9000 Å and nearly flat absolute sensitivity from 2500 to 9000 Å. For this experiment it is desirable to use the same type tube over the entire range of measurement (6000 to 8000 Å) and thereby not have too large a range of sensitivity over this range (so no special calibration is required). It is also desirable to have sufficient sensitivity to overcome electrical noise. Also, a side-on photomultiplier fits much more easily into the McPherson spectrograph. The C31025C was first tested in a similar two-channel analyzer using the same spectrograph and found to be sufficiently sensitive. Some of the earlier data in this experiment were taken with three type C31025B's on three channels instead of C31025C's. This tube has a GaAsP photocathode, is less sensitive, and has decreasing sensitivity with wavelength, but was adequate prior to

Figure 3.6. Photomultiplier divider chain wiring diagram.

the availability of the C31025C's.

From the photomultiplier chassis six RG58/U cables with type N connectors were run to a connector on the screen room and then to oscilloscopes. The oscilloscopes used are Tektronix types 551, 555, and R556. Seven dual beam oscilloscopes were used to record traces of magnetic fields, shock tube current (measured with a Rogowsky coil at the high voltage cable connection to the shock tube), laser pulse monitor, and the six Thomson scattering channels. The oscilloscopes were regularly checked for gain and DC balance. Differences in relative frequency responses of the scopes (the laser pulse is essentially a 25 MHz signal) and relative photomultiplier sensitivities were taken into account when the data were analyzed.

3.4 APPARATUS-- SYSTEM CHARACTERISTICS

3.4.a) Bandwidth and Discrimination of Polychromator.

An essential ingredient of the measurements made of the scattered light is the disperser or polychromator. It is fundamental to know its characteristics (wavelength response, bandwidth, and discrimination) in order to correctly interpret the signals from it. The McPherson spectrograph is built so that one can easily vary the wavelength incident upon the center of the exit slot. It is therefore easy to test the bandwidth characteristics of the polychromator by "tuning" the spectrograph wavelength dial when monochromatic light enters the spectrograph. For this purpose a He-Ne gas laser (wavelength 6328 Å) was used. Its beam was directed onto a rough (sanded) finished block of plexiglass mounted between the lens and the pinhole aperture on the scattered light collection port on the shock tube. An adequate anode current was produced in the photomultipliers (up to 400 nA; the DC limit for the tubes is 10 μ A) when the wavelength dial was tuned to the center of one of the detector channels. The current was measured by means of a Keithley 910C electrometer (nanamp scales). Scanning the front of the detector light guide with the known wavelength and recording the McPherson wavelength dial reading versus relative response determined the instrumental profile, channel width, and interchannel cross-

talk. When all baffles were properly placed in the detector between the light guides, no problem with interchannel crosstalk was ever observed. The system (with $\sim 45 \text{ \AA}$ between channel centers) gives a discrimination of better than 1 part in 500 for the arrangements used in these scattering experiments. A plot of relative response versus wavelength for one of the detector channels is shown in figure 3.7. This response is Gaussian to a considerable degree of accuracy; if anything, on the far wings of the response curve it falls off slightly faster than Gaussian.

The relative sensitivities of the six detector channels was measured by use of the He-Ne laser. The response was also measured at $6943 \text{ \AA}$ (ruby laser wavelength) by turning the wavelength dial to the center of each channel and measuring the signal output from the stray light in the shock tube in each channel in succession. This permitted an in situ measurement of system sensitivity and relative response all the way from the laser focal point to oscilloscope output.

2

Figure 3.7. Relative response versus wavelength off maximum for polychromator channel No. 5. The full width half maximum is 25 Å.

3.4.b) Rayleigh Scattering Test of System. A demanding test of the optical system is the ability to detect Rayleigh scattering from a cold (room temperature) gas in the shock tube. For a cold gas filling density n (cm^{-3}) such that the Rayleigh scattering from a given volume of the gas is the same as the integrated Thomson scattering intensity from a density n_e (cm^{-3}) of electrons in the same volume, the following relation holds:

$$n_e = n \frac{\sigma_R}{\sigma_T}, \quad \text{where} \quad \sigma_T = \frac{8\pi}{3} \left(\frac{e^2}{m c^2} \right)^2 = 6.65 \times 10^{-25} \text{ cm}^2$$

is the Thomson scattering cross section for free electrons and σ_R is the Rayleigh scattering cross section for the cold test gas. This gives an estimate of signal amplitude compared to known plasma light levels for this experimental set-up. More importantly, the capability of doing Rayleigh scattering in this manner enables one to determine the electron density in a plasma by Thomson scattering. See section 4.1.c for details of the evaluation of electron density when the temperature (i.e., the half width of the Thomson scattered light spectral profile) has been determined and the magnitude of the center (laser wavelength) channel signal has been measured (The signal for this channel is given by the total detected laser signal minus the known stray light level.).

Rayleigh scattering was successfully carried out on this equipment before Thomson scattering was done. The results of one such Rayleigh scattering experiment on argon are shown

in figure 3.8. The resulting signal level was sufficiently above plasma light levels to allow Thomson scattering at the laser powers available. In the graph of figure 3.8 the stray light level is equivalent to Rayleigh scattering from 70 Torr of argon, and is equivalent to the (integrated over all wavelengths) Thomson scattering from 0.6×10^{16} free electrons per cm^3 in the same scattering volume. This is a highly satisfactory level for use of the centerline Thomson scattering signal to determine plasma electron density if Thomson scattering is detected. In this experiment the Thomson scattering signal ranged from 5 to 50% of the total channel signal due to the laser light (The total signal is given by the Thomson scattered light plus the stray laser light plus the plasma light--the plasma light is the part not due to laser emission.). Stray light levels varied from day to day, due to shifting of components or realignments, so stray light tests and Rayleigh scattering tests were carried out at least once each day that experimental runs were made. The stray light equivalent in Thomson scattering from free electrons (integrated over all wavelengths) was within the range 0.4 to 2.2×10^{16} electrons/ cm^3 in the scattering volume. These levels correspond to Rayleigh scattering from 50 Torr and 260 Torr argon, respectively.

Figure 3.8 Photomultiplier response to Rayleigh scattering of ruby laser light by argon in the shock tube. The Rayleigh scattering signal is equal to the stray light level when the fill pressure is 70 Torr Ar.

CHAPTER IV ANALYSIS AND PRESENTATION OF THOMSON SCATTERING DATA

In this chapter we discuss the treatment of data obtained from the Thomson scattering experiment carried out on the shock heated plasma. First, the analysis of raw data is discussed. Second, errors and limitations of this diagnostic as employed in this experiment are considered. Finally, the electron temperature and density values obtained from this experiment are presented in both tabular and graphical form. Detailed discussion of the significance of the results is reserved for the following chapter.

4.1 TREATMENT OF DATA

All necessary data for this experiment was recorded by a bank of seven dual beam oscilloscopes (Tektronix types 551, 555, and R556). Figure 4.1 shows the oscilloscope traces of the six photomultiplier outputs for one data shot and a trace of the laser monitor output added algebraically to the integrated output of a magnetic field probe located at the same axial position as the scattering volume. The other data recorded by the oscilloscopes include: (1) the unintegrated magnetic field probe traces taken at the 150 and 175 cm positions of the shock tube--used to determine the shock speed by dividing the probe separation (25 cm) by the time of transit measured from these traces; (2) the integrated magnetic field probe

Channel number;
Center wavelength ($\text{\AA}$);
and Scope sensitivity (mV/div).

1;6845 $\text{\AA}$;10mV/div.

2;6898 $\text{\AA}$;20mV/div.

3;6943 $\text{\AA}$;50mV/div.
(center wavelength)

4;6995 $\text{\AA}$;20mV/div.

5;7041 $\text{\AA}$;20mV/div.

6;7088 $\text{\AA}$;10mV/div.

Magnetic field
plus laser monitor

Grid scale

Oscilloscope traces of the output from the six polychromator channels for one data shot. The arrow indicates time of laser pulse. The large burst of plasma light in channel 5 beginning .7 μsec after the laser pulse has been identified as CuI spark line at 7039 $\text{\AA}$. Sweep speed 1 $\mu\text{sec/div}$.

Figure 4.1

signal at the 175 cm position, triggered by the firing of the shock tube-- to give the total time of transit of the shock from launching wall to the scan point; and, (3) the (integrated) output of a Rogowsky coil mounted behind the launching wall of the shock tube-- to monitor the total current into the shock tube.

4.1.a) Electron temperature determination. The spike on each photomultiplier trace at the time of laser pulsing is due to free electron Thomson scattered laser light, shifted from the laser wavelength by Doppler broadening. A Maxwellian distribution of electron velocities will yield a Gaussian profile for the scattered light spectrum. The origin of this scattered spectrum is discussed in section 2.2.a. After introducing appropriate corrections for the different channel sensitivities, we ask whether the resulting profile is Gaussian and if so what is the full half width of the profile? In the experimental conditions observed in this work the Salpeter parameter, $\alpha = \frac{1}{k}\lambda_D$, ranges from about .05 to 0.6. Thus uncorrelated scattering by free electrons should occur and produce a Gaussian profile, although as $\alpha \rightarrow 0.6$ some correlation effects begin to appear. Section 5.4 deals briefly with some of the observed deviations from a Gaussian spectral profile in this experiment. As shown in section 2.2.a, for a Gaussian profile with half width $\Delta\lambda_{\frac{1}{2}}$, the temperature is given

by the formula $T_e = [\Delta\lambda_t/32.35]^2$ eV in the experimental arrangements used here.

Figure 4.2 shows several profiles of polychromator signal (corrected for channel response and normalized to one at the laser wavelength channel) as a function of channel shift in wavelength from 6943 Å. The graphs (a) and (b) of this figure are typical for this experiment, while graphs (c) and (d) are unusual cases which are discussed in section 5.4. For the graphs of figure 4.2 the channel half widths and center wavelengths were as follows:

Channel number:	1	2	3	4	5	6
Half width, $\Delta\lambda_c$ (Å):	20	21	23	24	25	22
Center wavelength (Å):	6845	6898	6943	6995	7041	7088
Center shift from 6943 Å, $\Delta\lambda$ (Å)	-98	-45	0	+52	+98	+145.

The other polychromator characteristics are discussed in section 3.4.a. The system was originally designed to resolve temperatures from about 2 to 1000 eV by changing the spacing of channels. However nearly all data were taken with the above channel spacing, suitable for $T_e \sim 2$ to 75 eV.

The data in the form shown in figure 4.2 were obtained after multiplying the measured scattered light pulse height by calibration factors for photomultiplier sensitivity, channel bandwidth, and oscilloscope sen-

SCATTERED LIGHT SPECTRAL DISTRIBUTIONS
FOR FOUR DATA SHOTS

Figure 4.2

sitivity and frequency response. On the center channel it was first necessary to subtract the stray light signal. The system as a whole was calibrated at least three times per week of data taking. Little change in these calibrations was noticed (less than 10 per cent overall per channel during the entire experiment).

Once the signals from each channel were evaluated and normalized, a plot of the logarithm of the signal, $\log_{10} S$, as a function of the square of the channel wavelength deviation from 6943 Å, $(\Delta\lambda)^2 \text{ Å}^2$, was made. A Gaussian profile has the form $S(\Delta\lambda) = A \exp\left[-\frac{(\Delta\lambda)^2}{B}\right]$, where $A \propto n_e / \sqrt{T_e}$, and $1/B = 2.66 \times 10^{-3} / T_e$ (n_e in cm^{-3} , T_e in eV). For a given shot, A is constant so a semilogarithmic plot of $S \propto \exp\left[-\frac{(\Delta\lambda)^2}{B}\right]$ versus $(\Delta\lambda)^2$ gives two straight lines rising symmetrically from left and right to a peak at $\Delta\lambda = 0$.

From such a plot it is simple to evaluate $\Delta\lambda_{\frac{1}{2}}$, the profile half width. Typically, however, the data did not produce such neat straight lines. Plot points deviated somewhat randomly above or below a least squares fit on the semilog plot. A "best fit" by means of the principle of least squares for a straight line on the graph $\log_{10} S$ versus $(\Delta\lambda)^2$ was made for all of the first series of data. The corresponding calculated limits of error in slopes of the lines were used to obtain extremal values of $\Delta\lambda_{\frac{1}{2}}$, and thus of T_e . It was

found on these data that equally good limits for T_e could be found by drawing lines of greatest and least slope (eliminating wild points) through the center channel plot point and the points on the wings of the semilog plot. It is these values for T_e error limits which are used in nearly all plot points shown in this chapter. Errors in T_e were less than 25 per cent (usually $\pm 15\%$) using either this procedure or the analytic evaluation of errors by the method of least squares on the semilog plot. It is possible to do a least squares fit to a Gaussian curve rather than to the straight line plot of $\log_{10} S$ versus $(\Delta\lambda)^2$ but the added arithmetic complications are not justified for this experiment.

Of the plot points on the $S(\Delta\lambda)$ curve, the greatest error was in the center channel. This is due to the frequent small level of the scattered signal compared to stray light. Errors in the four closest channels to the center were relatively small when T_e was high enough to give a signal in them. In these experiments a signal was not often noticed in the sixth channel (with $\Delta\lambda = 145 \text{ \AA}$), since the highest electron temperature measured was $43 \pm 10 \text{ eV}$. The traces of figure 4.1 show this channel (No. 6) to have electrical noise almost as large as the signal (noise later cleared up) because it is being recorded at the highest oscilloscope sen-

sitivity. The other five channels have signals well above electrical noise and plasma light fluctuations.

A typical Thomson scattering signal for a channel centered 50 to 100 Å off center was approximately 10 mV. Across the terminating resistor of $50\ \Omega$ at the scope this is a current maximum of 0.2 mA. The photomultiplier anode sensitivity for the RCA type C31025C is approximately 2.7 kA/W⁵⁵, so the incident power on the photocathode is 8×10^{-8} W. Over the 20 nsec of the laser pulse this is equivalent to approximately 1.6×10^{-15} J, or 6×10^3 photons. At a photocathode quantum efficiency of 10 per cent this gives 600 photoelectrons per channel--a good statistical level. These photomultipliers have low noise CuBe dynodes and photomultiplier noise is negligible at these signal levels and frequencies. Hence, errors for the high signal channels (numbers 1, 2, 4, and 5 of the previous chart), for $T_e \sim 10$ eV, are chiefly due to errors in reading the pulse height.

4.1.b) Background light: limitations of the technique for this experiment. An estimate of the ratio of Thomson scattered signal to plasma bremsstrahlung in a 20 Å wide detector channel for the densities and temperatures of this experiment (actual measurements-- $n_e \sim 2 \times 10^{16}$, $T_e \sim 15$ eV) is made in appendix C. These estimates, for full laser power (500 MW) give a ratio of 80 to 1. If only one third of the laser energy reaches the viewed scattering volume (Measurements of power transmitted through the shock tube beam ports give 300 MW or roughly two thirds of the laser light transmitted) we should expect signal to bremsstrahlung ratios of 25:1. Actual values never exceeded 4 or 5 to 1. Thus, we are viewing significantly more light through the collecting optics than a classical bremsstrahlung calculation for a pure deuterium plasma would imply.

Possible explanations for this significantly higher plasma light include: (1) radiation from high Z (highly ionized) impurities in the plasma, (2) recombination radiation, (3) line emission from plasma components or spark spectra of wall components, and (4) molecular band emission from hydrogen. We shall consider each of these in relation to the experiments carried out on the shock tube after first discussing other signal to noise problems.

Oscilloscope traces of photomultiplier scans of plasma light from the shock tube are not smooth. They have a jagged appearance, even when they are of high

absolute magnitude (i.e., are above electrical noise). The apparent period of these fluctuations is short, usually less than 100 nsec. When a laser Thomson scattering signal is not of a frequency near this fluctuation frequency, it is easily picked out from the surrounding noise. If the noise is of short period, the laser pulse, if of small amplitude, may be completely hidden. Figure 4.1 shows these fluctuations as well as a large sharp spike in the plasma light. For faster shocks, the initial spike may be sharper and of much larger amplitude (or may not even exist; very capricious device), thus swamping any laser signal. Examples of the plasma light output can be seen as well in figures 4.4 and 4.6. The high frequency fluctuations of plasma light generally seen may be due to turbulence in the observed plasma volume or to non-equilibrium effects in the shock front.

In order to obtain shock speeds greater than 25 cm/ μ sec, it is necessary to use a larger firing voltage on the shock tube. The usual firing voltage with an air fill in the switch gaps is 35 kV. To go higher than this, SF_6 is used in the switches, thereby allowing firing voltages of 75 to 85 kV for the same gap spacing. A shock speed of up to 45 cm/ μ sec was observed in deuterium, which, calculations indicate, should give an order of magnitude higher plasma temperature, $T_e + T_i$. For a higher T_e , and roughly the same n_e , the bremsstrahlung level per channel should be less because of a temperature

dependence of bremsstrahlung as $T^{-\frac{1}{2}}$. However, background light levels were observed to increase by roughly 20 times. Thus, any Thomson scattered signal is completely swamped by the high background level. Repeated attempts, without success, were made to observe scattered signals under these conditions. It is possible that more impurities will enter the plasma from the wall because electric currents and fields will be higher in the high voltage case. More wall erosion has indeed been observed to occur in the high voltage firings. That the currents and fields should be higher follows from the fact that the main bank current is near its maximum (quarter period of the bank is $5\mu\text{sec}$; $I_{\text{max}} \sim 2\text{ MA}$) and the voltage across the tube is still high ($\sim 40\text{ kV}$) when the shock reaches the test point (total propagation time from launching wall to reflecting wall is 5 or $6\mu\text{sec}$ for high firing voltages). On the other hand, for the 35 kV shots the bank voltage and current are well past peak output (total shock propagation time $\gtrsim 9\mu\text{sec}$, with the same $5\mu\text{sec}$ bank quarter period). The higher light levels appear to have different structures in time corresponding to different wavelengths and have been observed at other axial positions of the shock tube. Hence the higher light levels are not unique to the observation point used in this experiment and are probably not only due to bremsstrahlung emission since the plasma light intensity variation with wavelength differs from that for bremsstrahlung.

Therefore, in these experiments data on T_e and n_e for shock speeds in deuterium of 5 to 25 cm/ μ sec have been obtained, while at the faster speeds available in this device, Thomson scattering data could not be obtained because of the high background levels.

We now consider the reasons for the high background light observed in the Thomson scattering experiments. The possibility of reduced laser output was considered and the power which reaches the beam dump (i.e., passes through the scattering volume) was measured to be approximately 300 MW when the laser is run at its full power of 500 MW. The collecting optics view roughly 75 percent of the total focussed width. Thus, the attenuation of laser power due to beam handling and scattered light pinhole stop could reduce expected signal to background from the calculated (appendix C) 80:1 to 36:1. This is still much greater than the usually observed 3 or 4 to 1.

The plasma in the shock tube is in contact with the walls. Thus one could expect ablation of and possible excitation or ionization of some wall material. A small amount of such impurities near the wall could cause a significant increase in bremsstrahlung from the plasma layer near the wall. This is clear from the bremsstrahlung emissivity which varies as $\epsilon_\lambda \propto Z^2 n_e^2 / T_e^{1/2}$ and from recombination a similar increase may be expected as the

recombination rate is proportional to $Z^4 n_e n_i / T^{3/2}$.

The mechanism of recombination radiation is such that it is not likely to occur for hydrogen at the inner plasma temperatures of 10-30 eV. Recombination might occur in very low energy deuterium ions in the partly ionized unshocked gas in the first few millimeters of the viewing port collimator. Hence it is most probable that the enhanced continuum emission observed in this experiment is from added impurity electrons at low temperatures near the walls of the shock tube.

When the shock tube had not been used for a few days or had been filled with air, it was noticed that line spectra of neutral and singly and doubly ionized N, O, Al, and Si can be observed in photographically (integrated over time) recorded plasma spectra. The number and intensity of these lines decrease after the shock tube has been cleaned by firing about ten times. Relatively few such lines are noticeable at the red end of the spectrum, but below 6200 Å many such lines occur. In the oscilloscope trace for channel 5 ($\lambda_c \sim 7041$ Å) in figure 4.1, there is a large, wide (1 μsec) burst of light beginning 0.7 μsec after the laser pulse. This is the 7039 Å Cu I spark line which may appear at any time after the reflected shock has arrived at the scan position, whether the tube is clean or not. The magnitude of this burst ranges from unnoticeable to many times the light level at that time for the other channels. This appears to be

from the copper in the collecting port collimator in the tube wall and is not observable if an aluminum port is used. Similarly, less intense lines from Al, B, N, and Ca (the last three in the boron nitride reflecting wall) may occur, but none show up in high resolution photographic spectra at wavelengths in the range 6200 to 6900 Å.

Stark and Doppler broadening of H_{α} is not sufficiently large (only a few tens of Angstroms) to reach any detector channel used in this experiment.⁵⁶

Charge exchange between the plasma ions and the neutral deuterium molecules in the viewing collimator might produce excitation in neutral or singly charged deuterium ions. A study of Geissler tube discharges in hydrogen⁵⁹ found about 150 lines and bandheads of H_2 emission between 6615 and 7212 Å, some of them fairly intense. Such lines and bands would appear as a continuum in this system and would require considerably better resolution to study their contribution to the plasma light. A brief discussion of molecular continuum measurements on a hydrogen plasma beyond the Balmer series limit is in a note by Engelhardt⁵⁷. The effective cross section for charge exchange of protons in both atomic and molecular hydrogen is approximately 10^{-15} cm² for ion velocities $\sim 10^8$ cm/sec⁵⁸, a velocity which is the right order of magnitude for the value of ion energies measured on this shock tube²¹. Hence this

effect might occur in the cold un-ionized gas inside the collimator.

We conclude that the enhanced background light is probably mostly due to impurities in the plasma. The wall damage studies on the shock tube by B. Feinberg³⁵ indicate that pits of several microns depth in the outer shock tube wall occur over a period of a few hundred shots. This implies that a significant number of impurity particles may appear in the plasma.

4.1.c) Electron density determination. The necessary conditions and system calibrations for evaluation of n_e by Thomson scattering were discussed in section 3.4.b. These conditions were well met by this apparatus. Thus, for each T_e value a corresponding n_e value was obtained. The center channel stray light signal was from 1 to 20 times the Thomson scattering signal, and the latter signal could be found by subtracting the stray light level (found at the beginning of the data run and normalized for laser output power) from the total laser signal.

The stray light level in the center channel is equivalent to the Rayleigh scattering level from a certain density of argon test gas. This (scattering equivalent) argon density is determined from a plot of Rayleigh scattering output versus argon gas fill pressure, as in figure 3.8. This determination was made daily. Argon and, occasionally, dry nitrogen were used as the Rayleigh test gas. The cross sections for Rayleigh scattering from argon and for

Thomson scattering from free electrons are σ_R and σ_T , respectively. The ratio σ_T/σ_R is known⁹ to be 427.

Let the measured stray light level be S_{st} and the center line Thomson scattering signal be S_{sig} . Take $\Delta\lambda_c$ as the (measured, Gaussian) center channel half width, and $\Delta\lambda_{\frac{1}{2}}$ as the (measured, also Gaussian shape profile) half width of the Thomson scattered light spectrum (measured to obtain T_e). Then the plasma electron density is given

by
$$n_e = \frac{S_{sig}}{S_{st}} \frac{\sigma_R}{\sigma_T} n \frac{\Delta\lambda_{1/2}}{\Delta\lambda_c} \text{ cm}^{-3},$$
 where n

is the measured Rayleigh scattering equivalent (to stray light level) density of argon (cm^{-3}). This assumes that $\alpha \sim 10^{-8} \sqrt{\frac{n_e(\text{cm}^{-3})}{T_e(\text{eV})}}$ is less than 1 which is the case for nearly all the measurements carried out in this thesis. For the exceptional cases where $\alpha \rightarrow 1$

the temperature could be evaluated, but not the density, by eliminating the center line signal. See section 5.4 for a discussion of these exceptional cases. Typical errors in evaluated electron densities are within $\pm 30\%$.

This arises from the error in evaluating the center channel signal and $\Delta\lambda_{\frac{1}{2}}$, the scattered spectrum half width.

4.2 PRESENTATION OF DATA

Each time Thomson scattering was observed from the shock heated plasma, a value of T_e and n_e was obtained. In order to make a profile of T_e and n_e in time as the plasma passes, it is necessary to have a consistent time base relative to which the plot points may be placed. This problem arises because, from shot to shot, the timing of the shock front varies as do the fine features of the magnetic field and (more strongly varying) plasma light at the scan point. Nevertheless, two features are consistently observable in the azimuthal magnetic field traces (B_θ)-- the sharp rises in field upon arrival of both the incident and reflected shocks. Hence we use the beginning of rise in B_θ for the incident shock as the origin of time in the T_e and n_e profiles of the passing plasma. In the B_θ profiles there are seven distinct regions which show up as well in the T_e and n_e profiles. These regions are:

- (1) shock front-- sharply rising (100-200 nsec) B_θ ;
- (2) post shock plasma-- steady B_θ ;
- (3) magnetic piston-- sharp but small rise in B_θ , maybe just a small bump;
- (4) a post-piston but pre-reflection rarefaction region ("vacuum") seen for some cases;
- (5) reflected shock front-- large rise in B_θ , usually slower than the incident rise;
- (6) reflected shock heated plasma; and
- (7) rarefaction region after the reflected shock plasma has passed. For

each laser data shot, a monitor signal was added algebraically to the B_θ profile for the laser scan axial position to permit a sharp definition of the time of laser pulsing relative to magnetic field features.

On nearly all shots it is thus possible to place the laser pulse in one of the seven magnetic field regions and to subdivide the region to find the relative pulse times within these regions. A few magnetic field profiles are diffuse and indicate a poorly formed shock structure. T_e and n_e values found in these cases are close to those found for well formed shocks, but because of timing difficulties are not included in the profiles of T_e and n_e shown here.

Average values of measured plasma properties are shown in tables 4.1 and 4.2. The initial parameters (fill pressure of D_2 , bias magnetic field, and shock speed) are listed to define the various cases studied. The average magnetic compression ratios due to incident and reflected shocks are given. Next in these tables are shown density and temperatures as measured for the incident, vacuum, and reflected plasmas. The snowplow theory values for non zero initial magnetic field are those calculated by R. Taussig as described in section 2.1 of this thesis. The maximum density values given in brackets are unusually high occasionally observed values of density which are not included in the n_e profiles.

Following these tables are profiles of T_e and n_e made up from the T_e and n_e values obtained by Thomson scattering on the shock heated and compressed plasma, scanned from a point 5 cm in front of the reflecting wall. A typical B_θ trace for the same axial position is shown , and (for two cases) a trace of plasma light seen by a detector of 20 Å bandwidth centered at 6900 Å . The time of initial compression of magnetic field serves as the time origin for all five sets of profiles. More profiles were not plotted because at least 20 points are necessary to achieve a reasonably significant profile. The smaller number of data points for these cases are, however, included in tables 4.1 and 4.2.

These profiles are the first detailed density and temperature plots obtained for this type of shock heated plasma. The following chapter discusses the significance of the information shown here.

INITIAL FILL PRESSURE, $P_1 = 50$ MILLITORR DEUTERIUM					
INITIAL MAGNETIC FIELD, B_1 (kG)		3.6	1.0	0	-3.6 (BIAS AND DRIVE CROSSED)
SHOCK SPEED, U_s (cm/ μ sec)		23 ± 2	$12.5 \pm .5$	7 ± 1	< 3
MAGNETIC FIELD B_2/B_1 COMPRESSION B_5/B_1		$1.75 \pm .25$ $3.3 \pm .5$	$3 \pm .2$ $5.3 \pm .8$	$(\frac{B_5}{B_2} = 2 \pm 4)$	—
REMARKS		SEE FIG. 4.3	POORLY DEFINED	SEE FIG. 4.6	DIFFUSE SHOCK, NO DATA
INCIDENT SHOCK	FRONT T_e (eV) POST SHOCK PISTON	10 ± 3 7 ± 2 9 ± 3	— 7	13 ± 5 6	—
	SNOW PLOW THEORY ($T_e + T_i$)	140 ± 40	60	30^*	—
	$n_e (\times 10^{16} \text{ cm}^{-3})$ AVERAGE MAX.	1.5 $2.3 \pm .3$	1.0 2.9	$\sim .75$ $1.8 \pm .2$	—
VACUUM		T_e (eV) $n_e (\times 10^{16} \text{ cm}^{-3})$	— $< .2$	— $< .5$	< 3 $\sim .5$
REFLECTED SHOCK	FRONT T_e (eV) AVERAGE BACK	32 ± 4 15 20	NO DATA	— ~ 6 —	—
	SNOW PLOW THEORY ($T_e + T_i$)	260 ± 90	120	—	—
	$n_e (\times 10^{16} \text{ cm}^{-3})$ AVERAGE MAX.	2.0 3.5 [7 ± 3]	NO DATA	$\sim .5$	—

* ESTIMATE IF ALL ENERGY OF SHOCK MASS MOTION WENT INTO THERMAL MOTION

ELECTRON TEMPERATURE AND DENSITY FOR SHOCK WAVES INTO 50 mTORR DEUTERIUM

Table 4.1 (For the crossed fields case T_e was measured to be a few electron volts with no structure resolved.)

INITIAL FILL PRESSURE, P_i = 100 MILLITORR DEUTERIUM					
INITIAL MAGNETIC FIELD, B_i (kG)		3.6	1.0	0	-1.0 (BIAS AND DRIVE CROSSED)
SHOCK SPEED, U_s (cm/ μ sec)		16 ± 2	9 ± 4	$4.5 \pm .5$	$3.8 \pm .2$
MAGNETIC FIELD B_2/B_i COMPRESSION B_5/B_i		$2 \pm .4$ 4.5 ± 1.0	$2.5 \pm .5$ $5.2 \pm .5$	$(\frac{B_5}{B_2} = 2.8 \pm .2)$	—
REMARKS		SEE FIG. 4.4	SEE FIG. 4.5	POORLY DEFINED SHOCK	SEE FIG. 4.7
INCIDENT SHOCK	FRONT T_e (eV) POST SHOCK PISTON	15 ± 5 10 ± 5 25 ± 10	8 ± 2 5 ± 1 15 ± 5	PRESHOCK ~ 3 POSTSHOCK ≤ 6	AT $B=0$: 11 ± 7 BEFORE: 6-10
	SNOW PLOW THEORY ($T_e + T_i$)	80 ± 20	~ 70	NOT FULLY IONIZED	— *
	$n_e (\times 10^{16} \text{ cm}^{-3})$ AVERAGE MAX.	$1.5 \pm .2$ [3.2]	$.4 \pm .1$ $.6 \pm .2$	1.0 3.7	AT $B=0$: $3.2 \pm .7$ BEFORE: 3.0
VACUUM	T_e (eV) $n_e (\times 10^{16} \text{ cm}^{-3})$	10, 20 .8	2 $\sim .2$	< 3 ~ 1	3 ± 1.5 $.6 \pm .3$
REFLECTED SHOCK	FRONT T_e (eV) AVERAGE BACK	27 ± 5 29 ± 3 40 ± 10	24 ± 6 ~ 15 ~ 10	≤ 9	—
	SNOW PLOW THEORY ($T_e + T_i$)	140 ± 20	~ 125	—	—
	$n_e (\times 10^{16} \text{ cm}^{-3})$ AVERAGE MAX.	2.5 $\{ 3.1 \pm .6$ [6.0]	$.7 \pm .2$ $\{ .9 \pm .2$ [1.8]	~ 1.9	—

* estimated $T_e \sim 6.5$ eV, if all the energy of shock front mass motion went into thermal motion and subsequently $T_e \rightarrow T_i$.

ELECTRON TEMPERATURE AND DENSITY FOR SHOCK WAVES INTO 100mTORR DEUTERIUM

Table 4.2

Figure 4.3 Transverse Ionizing Shock Wave

$P_1 = 100 \text{ mTorr, D}_2$
 $B_1 \sim 3.6 \text{ k Gauss}$
 $U_S \sim 16 \text{ cm}/\mu\text{sec}$

Figure 4.4 Transverse Ionizing Shock Wave.
 For this case all the original plot points are
 shown for T_e (without averaging for clusters);
 only four representative points are shown for n_e .

Figure 4.5 Transverse Ionizing Shock Wave

Figure 4.6 Gas Shock (Zero initial magnetic field). The dotted line in the T_e plot corresponds to the observations made when there is a small initial dip in magnetic field.

Figure 4.7 Crossed Bias and Drive Fields

CHAPTER V RESULTS AND CONCLUSIONS

In this chapter we first summarize the results of this experiment as illustrated in figures 4.3 through 4.7 and in tables 4.1 and 4.2. The most prominent feature of the data is that the observed electron temperatures for the transverse ionizing shocks (i.e., those with an initial magnetic field perpendicular to the direction of shock propagation) are much lower than the computed equilibrium shock temperatures and measured ion temperatures. The ratios of computed equilibrium ion temperatures to measured electron temperatures are within $4 < T_i/T_e < 10$. Energy relaxation between ions and electrons in this plasma is examined and found not likely to occur within the times and energy ranges characteristic of these experiments. In fact, the Coulomb relaxation theory is not valid for the electron and ion temperature differences observed in the initial plasma following a strong transverse ionizing shock. Strikingly good agreement with collisional relaxation behavior is found for the (relatively slower) ionizing gas dynamic shocks propagating into a magnetic field free region. For the case of antiparallel piston and initial magnetic fields, theoretical expectations are found within the limited detail of the measurements made.

Resistive heating of the electrons explains satis-

factorily the observed peaks in electron temperature for the transverse shocks. The relationship between T_e and n_e is considered and appears to show an interesting interdependence for some of the shocks studied. The ratio of magnetic field compression to plasma density was found to be smaller for fast shocks than for slow shocks studied in this experiment. Possible deviations of the observed scattered light profiles from a Gaussian shape are considered as evidence of correlated scattering due to ion thermal motion and as an indication of a non-Maxwellian velocity distribution of the plasma electrons.

Finally, some related experiments that could be performed on this kind of shock heated plasma are considered.

5.1 THE DATA

5.1.a) Discussion of figures 4.3 to 4.7. Each of the profiles of T_e and n_e is made from 25 to 45 separate data points. Typical B_e and plasma light (for some cases only) traces are shown as recorded by oscilloscope traces of magnetic field probes and photomultiplier responses. The data are presented in order of decreasing shock speed.

Typical error ranges in T_e and n_e are shown for some of the points indicated on the graphs. The indicated plot points near maxima in T_e and n_e are composites of several data shots. These maximal points in the plasma were studied more intensely than those which seemed to give relatively low and stable values. Consider the form of these profiles, which show the different modes of operation of the shock tube. The first three (4.3 - 4.5) are for transverse ionizing shocks. The fourth (4.6) is for a gas shock (no bias field), and the last (4.7) is for antiparallel (opposite polarity) initial and piston (driver) fields.

The first three sets of profiles (4.3, 4.4, and 4.5) are for transverse ionizing shocks. Consider figure 4.3 as typical of the others. The regions of plasma are labelled and correspond to the incident shock, the post shock plasma, the piston, the rarefaction region, the arrival of the reflected shock, the reflected

shock heated plasma, and the expansion fan. The front and back of the incident plasma have significantly higher values of T_e than the intermediate post shock plasma. As will be shown, this higher value of T_e results from Joule heating of plasma electrons by currents in the plasma. Early in this work some profiles of T_e and n_e were made without a reflecting wall in the shock tube. These profiles have the same structure as as the incident plasma shown in figures 4.3 to 4.5. The dotted line on the T_e profile for figure 4.4 corresponds to a less distinct separation of incident and reflected plasma which was occasionally observed. The plot points for this particular time lay originally in two groups of three points each. Rather than take an average for the curve, two curves were drawn. As the lower (solid line) of the two curves slopes downward, it is exhibiting the same reduction in T_e (and n_e) which occurs if there is no wall in the tube and all of the plasma has passed the observation point. Separation of incident and reflected plasmas is more distinct in the other cases. The part beyond the dip in T_e corresponds to the wall reflected plasma propagating in the opposite direction to the incident shock which has passed this point at an earlier time. As the reflected shock arrives there is a rise in T_e and n_e to values as much as several times those in the incident plasma.

For the transverse ionizing shocks the high initial T_e values rise in a time of .2 to .3 μ sec (corresponding to a space dimension of 3 or 4 cm), which is comparable to the magnetic field rise time. We shall see that this temperature rise is due to Joule heating of the electrons by the currents which cause the magnetic fields. It is doubtful that the data imply any kind of damped oscillations following this initial heating (as observed by Martone and Segre on a theta pinch ⁶⁰), because the error levels in data for the post shock plasma are too high to notice small oscillations and because shot to shot variation of fine features in B_θ (our timing standard) make it difficult to locate plot points accurately.

The fourth set of profiles (figure 4.6) corresponds to a slower shock propagating into a region with no initial magnetic field. There is a quite clear signal from plasma light and electron density before the magnetic field starts to rise. Occasionally, as has been noted by others ²⁸, a small initial negative dip in magnetic field may occur prior to the expected positive rise. This is possibly due to a small bias field created by inductive (transformer) effects between the drive circuit and the unused bias circuit. The dip in B_θ is indicated in the figure by a dotted line on the B_θ profile, and the early rise in T_e when this occurs is similarly noted. The reflected plasma has a much lower value of T_e and n_e .

than the incident plasma, presumably because the plasma has decayed (With much lower magnetic fields the confinement time is less) and the shock wave itself has decayed. The presence of a long post shock region of plasma observed when a bias field is employed is not seen here. The magnetic piston may not be so clearly defined (confined to a narrow axial dimension), or may be less effective as an "impermeable" contact surface between the plasma and the vacuum behind it (thus allowing plasma to leak through). This gas dynamic shock has arrived at the observation point long ($\sim 20 \mu\text{sec}$) after the bank current has been at its peak (bank quarter period $\sim 5 \mu\text{sec}$). The slow rise in n_e at late times is apparently due to reionization of deuterium and impurity atoms as the main capacitor bank circuit rings.

The last set of profiles (figure 4.7) is for the case of bias magnetic field antiparallel to the direction of the drive magnetic field^{32,61}. No reflecting wall was in the shock tube when this study was made. The piston magnetic field is larger in magnitude than the bias field, so it raises the initial field through zero to a positive value at the back of the shock. The density and temperature rise to a maximum when $B \sim 0$. The piston field is not large enough to cause a sharp, distinct rise in B , or, apparently, to cause the sharper, larger jump in n_e and T_e anticipated due to field annihilation⁶¹.

5.1.b) Late time behavior. For times up to $20 \mu\text{sec}$ beyond the time intervals shown in figures 4.3 to 4.7, Thomson scattering in the shock tube can still be observed. The temperature and density of this residual plasma seem to lie somewhat randomly in the range 1 to 4 eV and $.2$ to $.8 \times 10^{16} \text{ cm}^{-3}$. A few (3 or 4 out of about 25) apparently very high values (up to $8 \times 10^{16} \text{ cm}^{-3}$) values of n_e were observed at times from 5 to $20 \mu\text{sec}$ after the incident shock arrival. Such high scattering signals may not really be from Thomson scattering by equilibrium plasmas-- it may be enhanced scattering from fluctuations. More likely they are due to scattering from large concentrations of ionized wall impurities which have drifted, after long periods of time, into the center of the plasma volume. Plasma light levels seem low at these late times compared with the levels when the shock heated plasma is present. It is probable that this late presence of electrons is due to reionization, as the main condenser bank circuit rings, of deuterium and impurity atoms which are in the shock tube long after the shock heated plasma has died away.

An estimate of plasma cooling time may be obtained from the half cooling time $t_{\frac{1}{2}} \sim l^2/\chi$, where l is a linear plasma vessel dimension and χ is the thermal diffusivity, defined in terms of K the thermal conductivity by $\chi = K/\rho C_p = K/\frac{3}{2}nk$. The half cooling time is the time for a body at initially uniform temperature T , with cold bound-

aries, to cool to a temperature $\frac{1}{2}T$. There is a magnetic field parallel to the plasma boundaries (azimuthal bias and drive fields). Consequently the transverse conductivity (which is smaller than the usual conductivity for a plasma by the factor $(1 + \omega_c^2 \tau^2)$ where ω_c is the cyclotron frequency and τ is the particle collision period) is the appropriate thermal conductivity for this plasma. Using the classical value of the thermal conductivity ⁶², we get, for deuterium,

$$t_{\frac{1}{2}} \approx 1.1 \times 10^3 \ell^2 T^{1/2} B^2 / n_i \ln \Lambda \text{ sec},$$

where T is in eV, B in gauss, n_i (ion density) in cm^{-3} and $\ln \Lambda$ is the Coulomb logarithm (equal to 8-12 for the plasmas of interest here). For $T_i \sim 300$ eV, $B \sim 3$ kG, $\ell = 2.5$ cm, and $n_i \sim 2 \times 10^{16} \text{ cm}^{-3}$, $t_{\frac{1}{2}} \sim 5 \mu\text{sec}$. From the temperature profiles, it appears that the plasma cools faster than this. Actually, the magnetic piston pushes incident and reflected plasmas beyond the scan point and we cannot determine how long the dense, shock compressed plasma remains in the shock tube. Some part of the late time density behavior may be due to rarefied plasma which has expanded from the once dense localized post shock plasma. However, given the long time span over which electrons are present and the lower magnetic fields at late times (so the half cooling time decreases, too) it is most probable that atoms in the shock tube have been reionized by the ringing of the bank.

5.1.c) Shock front width. Defining the shock front rise time as the time of rise from initial magnetic field to full magnetic compression, the data indicate that T_e and n_e rise in the same time scale (.2 to .3 μ sec, with n_e rising slightly more slowly). For the zero bias magnetic field case, T_e and n_e rise more slowly because of precursor effects. The antiparallel fields case also indicates a diffusive, slowly rising shock structure and an extensive precursor.

A collisional shock front, by definition, is one with a width greater than a few times the mean free path evaluated in the preshocked gas. For hydrogen molecules the collisional cross section may be estimated from the molecular diameter $d \sim 2.5 \times 10^{-8}$ cm, which gives $\sigma = 2^{\frac{1}{2}} \pi d^2$, and a mean free path, $\lambda = \frac{1}{n\sigma}$, of 0.22 and 0.11 cm for 50 and 100 mTorr D_2 . The transverse ionizing shock front thickness (given by the shock transit time times the shock speed) is on the order of 3 to 4 cm with respect to the jumps in B , T_e , and n_e . The shock is therefore collisional in the sense that the shock jump occupies many mean free paths.

5.2 THE PLASMA TEMPERATURES

5.2.a) Predictions and measurements of shocked plasma temperatures. Snowplow model calculations of the single fluid temperature for incident and reflected plasmas, with a transverse bias field, are shown in figure 2.2. The data are representative of the initial pressure and magnetic field cases studied in this experiment. As can be seen by examining either the appropriate curve of this figure or the snowplow model values shown in tables 4.1 and 4.2, all the measured values of T_e in this experiment lie well below the values indicated in these calculations. The value of T_e and T_i which should result from equilibration of these two temperatures following shock heating is given by $\frac{1}{2}T$ where T is the single fluid value of the snowplow model. The errors indicated in the table of snowplow temperature values are intended to cover the upper and lower values of T in the range of shock speeds that the initial conditions give. For speeds less than 30 cm/ μ sec the calculated temperatures are very sensitive to shock speed and may increase an order of magnitude or more for an increase in speed of 10 cm/ μ sec. Peak measured values of T_e range from 10 to 45 per cent of the expected $\frac{1}{2}T$ of these calculations. The post shock plasma regions in general have noticeably lower electron temperatures than the peaks, so that T_e for the bulk of the post-shock plasma will be 5 to 25% of the expected $\frac{1}{2}T$.

S.H. Robertson ²¹ has measured the plasma ion temperature to be of the order of 300 eV for a shock speed of approximately 20 cm/ μ sec in deuterium. He has also found that the ion temperature is of the order of 500 eV for the reflected shocks. These measurements have been made using a charge exchange neutral particle energy analyzer and are integrative in time (so no time resolution over the shock structure is provided). We assume that the ions reach the measured value rapidly, i.e., within the shock front. The agreement of this ion temperature with the theory is reasonably good. Hence, these experiments have shown that this method of shock heating a plasma creates hot ions and cold electrons. This is indeed what collisional shock theory predicts for the immediate post shock region where the ions should be hotter by collisional heating. It does not, however, explain the lack of equilibration of T_e and T_i at later times ($t > 0.5 \mu$ sec).

The temperatures expected in the zero bias field and the antiparallel field cases are low compared to those expected for the transverse ionizing shocks. A crude estimate based on the assumption that all shock front mass motion becomes ion thermal motion gives a temperature for the zero bias, 50 mTorr filling pressure case ($U_s \sim 7$ cm/ μ sec) after energy equipartition of $T_e = T_i = 20$ eV. This compares well with the measured value of 13 ± 5 eV. A more interesting comparison of this measured

profile can be made with one that was calculated in detail by S.H. Schneider ²⁶ on the basis of the model of Jaffrin and Probstein ⁶³. His result is shown in figure 5.1. This figure shows a two temperature gas shock profile for a $9 \text{ cm}/\mu\text{sec}$ shock in hydrogen (equivalent to $6.4 \text{ cm}/\mu\text{sec}$ in deuterium). In the profile of figure 5.1, T_e and T_i become equal to 10.5 eV after $\sim 0.2 \mu\text{sec}$ (Preheating of electrons is shown, along with the much more rapid ion heating.). The Thomson scattering profile of figure 4.6 shows $T_e = 12 \text{ eV}$ for $U_s \sim 7 \text{ cm}/\mu\text{sec}$ after $0.75 \mu\text{sec}$ (about $0.5 \mu\text{sec}$ when the initial self-bias magnetic field appears). Schneider's simulation does not show later time behavior. Apparently, for the slow gas shock case observed in this experiment, equipartition has occurred.

No such computational estimate of T_e and T_i is available for the case of antiparallel fields studied here. However, Taussig ⁶¹ has shown that there should be significant and rapid increases in T_e and n_e within the magnetic piston when $B=0$ due to annihilation of opposed fields and conversion of field energy into electron energy. For a $4 \text{ cm}/\mu\text{sec}$ shock, all of whose mass motion was converted into ion thermal energy and subsequently underwent equipartition with the electrons, T_e would be 6.5 eV. This is less than the measured $11 \pm 7 \text{ eV}$ at $B=0$. Thus, the electrons seem to be at least as hot as the crudest model would imply.

TWO TEMPERATURE GAS SHOCK ($B=0$) PROFILES IN LAB-FRAME COORDINATES WITH REAL VISCOSITY FOR $6.4 \text{ cm}/\mu\text{sec}$ SHOCK IN DEUTERIUM

Figure 5.1
(From ref. 26)

Only for the slowest propagation speeds ($U_s \sim 3$ to $7 \text{ cm}/\mu\text{sec}$) without preshock magnetic fields do the electrons reach the predicted single fluid temperature. A discussion of ideas as to why the electrons do not relax to a higher temperature comparable to the ion temperature in the faster shocks is given in the next section.

5.2.b) Electron-ion temperature equilibration. At the rear of a fast collisional shock front structure, the ion temperature has, in theory, reached its maximum value while the electrons have a very low temperature. The separate ion temperature, T_i , and electron temperature, T_e , are rapidly achieved by their individual thermalization to a Maxwellian velocity distribution by ion-ion and electron-electron Coulomb collisions, respectively. Assuming that the collision mean free path is less than the characteristic length of the plasma, energy exchange collisions will cause the equipartition of energy among the plasma particles according to the following formulas (see, for example, references 64, 62, and 58). The electrons will reach a Maxwellian velocity distribution with final temperature T_e (eV) from any initial distribution of velocities in a deuterium plasma of density n (cm^{-3}) in a time

$$\tau_{ee} \approx 3.3 \times 10^5 T_e^{3/2} / n \ln \Lambda \text{ sec}, \quad (5.1)$$

where $\ln \Lambda$ is the Coulomb logarithm (value 8 to 12 for the plasmas studied here). The deuterium ions will thermalize to a final temperature T_i in a time (sec)

$$\tau_{ii} \approx 2.02 \times 10^7 T_i^{3/2} / n \ln \Lambda = \sqrt{\frac{m_i}{m_e}} \tau_{ee}(T_i). \quad (5.2)$$

And for initial Maxwellian distributions (at T_e and T_i) of electrons and ions to relax to a common temperature, it takes a time

$$\tau_{eq} \approx \frac{3 m_i m_e k^{3/2}}{8 \sqrt{2\pi} n Z^2 e^4 \ln \Lambda} \left(\frac{T_i}{m_i} + \frac{T_e}{m_e} \right)^{3/2}, \quad (5.3)$$

which for $T_1 \ll T_e (m_i/m_e)$ can be written (Sec)

$$\tau_{ei} \approx 6.1 \times 10^8 \frac{T_e^{3/2}}{n \ln \Lambda} = \frac{m_i}{2 m_e} \tau_{ee}(T_e), \quad (5.4)$$

where T_e is the initial electron temperature. The following condition must hold if a (quasi-) equilibrium* state is to result from an arbitrary state⁶⁴:

$$\tau_{ee} \ll \tau_{ei} \text{ and } \tau_{ii} \ll \tau_{ie}. \quad (5.5)$$

The present measurements of T_e provide a value of T_e immediately following the shock jump, at which time T_1 has presumably reached its maximum value. As an example, take the value $T_e \sim 10$ eV at the peak of the initial jump of figure 4.3 ($U_s \sim 23$ cm/ μ sec, $P_1 = 50$ mTorr, $B_1 \sim 3.6$ kG, and $n_e \sim 1.5 \times 10^{16}$). Then, using formulas 5.1 through 5.3, $\tau_{ee} = 7 \times 10^{-11}$ sec. For the measured ion temperature, $T_1 \sim 300$ eV, $\tau_{ii} = 7 \times 10^{-7}$ sec. From the temperature 10 eV for the electrons, equilibrium should be reached with the ions in the time $\tau_{ie} = 1.3 \times 10^{-7}$ sec. From these estimates it is clear that the electrons should classically thermalize to a Maxwellian velocity distribution. On the other hand, $\tau_{ie} \leq \tau_{ii}$, which does not satisfy the requirement for attainment of a quasi-equilibrium that $\tau_{ii} \ll \tau_{ie}$.

What has happened in this case is basically that the condition for applicability of the rate equation for temperature equilibrium, $\frac{dT}{dt} = \frac{T_e - T}{\tau_{ie}}$ (5.6)

has not been satisfied. This condition requires that

the ion and electron temperatures should be within an
* i.e., electrons and ions have Maxwellian distributions as $T_e \rightarrow T_1$.

order of magnitude of one another for the rate equation and the relaxation formulas above to be correct. Here T (i.e., T_1) is greater than T_e by ~~more~~ than an order of magnitude. If $T_1 \ll (m_1/m_e)^{1/3} T_e$, then $\tau_{11} \ll \tau_{1e}$. In this case $(m_1/m_e)^{1/3} = 15.3$, while $T_1/T_e = 30$; thus, $T_1 > (m_1/m_e) T_e$. Sivukhin ⁶⁴ shows that the condition

$$T_1 \ll (m_1/m_e)^{1/3} T_e \quad (5.7)$$

is sufficient to ensure that the electron and ion velocities are Maxwellian throughout the time of energy exchange, with the possible exception of a short transient stage at the beginning.

The condition (5.7) is not both a necessary and sufficient condition for the appearance of a quasi-equilibrium state (T_e and T_1 both Maxwellian on the way to their equal temperatures). By assuming that the ion temperature is initially Maxwellian (despite the fact that the theory indicated that $\tau_{11} \gg \tau_{1e}$) and since $\tau_{ee} \ll \tau_{1e}$, it is reasonable to assume that the interaction of ions and electrons with one another will not interrupt their respective Maxwellian distributions. Thus, the ion and electron distributions may be Maxwellian throughout the ion-electron energy exchange ⁶⁴.

This assumption is not customarily made, and in general is probably not justified. In particular, the case with $T_e = 0$ and T_1 any finite value raises the question of how to substitute in the formulas (5.1)

and (5.4) and obtain a physically meaningful relaxation time. A search of literature on this subject indicates a lack of theory to cover the large initial divergence in T_i and T_e (E.g., references 62, and 58 do not discuss such a case; only 64 mentions it.) observed in this experiment for the fastest shocks. It is clear that in this plasma T_e does not approach T_i in a time of the order of 0.2-0.5 μ sec as would be predicted by theory for the initial temperatures $T_e=10$ eV and $T_i=300$ eV. From the T_e profiles it appears that T_e does gradually increase over time, but it is not following the relaxation times defined above. It is possible that in the post shock plasma the electrons and ions do not undergo collisions (That the collision lengths and times of classical formulas require collisions to occur is seen in the following paragraph.), and the temperature structure is dominated by the bulk plasma dynamics, i.e., adiabatic compression and expansion in post shock and expansion fan regions, or radial flows (Calculations in reference 24 show a radial profile of the wall reflected plasma in which the plasma is much denser near the center electrode of the shock tube, and this dense region moves radially outward in time.). Investigation of such behavior is beyond the scope of the present experiments.

The above collisional energy transfer times may be used to calculate the mean free paths λ_{ee} , λ_{ii} , and

λ_{ie} for electron-electron, ion-ion, and ion-electron collisions. Electrons of temperature 10 eV have an average thermal speed $v_e \sim 2.12 \times 10^8$ cm/sec and deuterium ions with temperature 300 eV have thermal speed $v_i \sim 1.91 \times 10^7$ cm/sec. Since $v_i \ll v_e$ we may consider the ions to be essentially at rest relative to the electrons, so the relative velocity for ion-electron collisions is v_e . The mean free path for collisions is given by the velocity times the collision time. For electron-electron collisions, the mean free path is thus $\lambda_{ee} \sim 10^{-2}$ cm; for ion-ion collisions, it is $\lambda_{ii} \sim 15$ cm; and for ion-electron collisions, $\lambda_{ie} \sim 24$ cm. When the strong azimuthal magnetic field is taken into account, the electrons and ions are constrained to move in orbits with centers on the field lines, which are concentric with the shock tube electrodes. These orbits have radii given by the electron and ion Larmor radii. For the temperatures $T_e = 10$ eV and $T_i = 300$ eV and a magnetic field of approximately 7 kGauss, the Larmor radius for electrons is 1.7×10^{-3} cm, and for ions it is 0.57 cm. Both these dimensions are less than the distance to the shock tube walls from the Thomson scattering scan point (roughly midway between the inner and outer conductors). Hence, many ion-electron collisions should occur in the observed plasma lifetime (of at least $1 \mu\text{sec}$).

In section 5.4 the possibility of a two electron temperature distribution in the plasma is considered, on the basis of an occasional appearance of rather broad wings in the spectrum of scattered light. This deviation from a Maxwellian, if significant beyond the errors in data taking, might suggest a partial collisional equilibration of the electrons with the hotter ions. If this occurs, it involves relatively few electrons.

It is interesting to note that the equilibration of electron and ion temperatures is to be expected for the case of the slow, zero bias magnetic field, shocks in deuterium, for which the profiles of T_e and n_e are shown in figure 4.6. For this case, the electron temperature was measured to be 10 eV, and the ion temperature is predicted by elementary theory to be 40 eV. The electron ion relaxation time is found by the previous formulas to be $\tau_{ie} \sim 0.25 \mu\text{sec}$, and the mean free path for collisions between ions and electrons to be slightly less than 1 cm. The shock thickness is several centimeters and the time of rise to maximum electron temperature is approximately 0.5 to $0.75 \mu\text{sec}$, in good agreement with this estimate of relaxation lengths and times.

The conclusion to be derived from these electron measurements, and the Coulomb collision theory results above, is that electron-ion collisions do not heat the plasma electrons in the time available. We proceed

on the assumption that the electrons are heated by other means than collisions with hot ions.

5.3 ELECTRON HEATING AND OTHER EFFECTS

5.3.a) Joule heating. A mechanism for electron heating is resistive dissipation. Figure 5.2 illustrates the configuration of shock tube currents and magnetic fields. At the piston and Shock ^{front} Λ axial positions of the shock magnetic profile there are jumps in the magnetic field and a current flows across the annular gap of the shock tube. It is interesting to note that the electron temperatures for all the transverse ionizing shocks show a sharp increase at the points where such magnetic jumps occur, followed by a drop in T_e just after the rise. The power dissipated per unit volume by resistive heating is the force on the electrons due to collisions with ions (nearly stationary ions, as $v_1 \ll v_e$) times the drift velocity of electrons relative to ions.

Consider the profile data of figure 4.3. B_θ rises uniformly from 3.6 kG to 8.5 kG in approximately .13 μ sec, as does T_e from 0 to 10 eV in the same time. The shock front width is approximately 3 cm. Both B_θ and T_e are measured at a position $R = 9$ cm from the axis of the shock tube. Figure 5.2 illustrates the relevant dimensions. The current in the center conductor (which branches off into the shock front) that produces the magnetic field in the post shock region is (MKS

Figure 5.2

DIAGRAM OF CURRENTS AND ASSOCIATED MAGNETIC FIELDS IN SHOCK TUBE

units for this section) $I = 2\pi \Delta B R / \mu_0 = 2.2 \times 10^5 \text{ A}$, where ΔB is the jump in B ($4.9 \text{ kG} = 0.49 \text{ T}$) measured at $R = 0.09 \text{ m}$ from the axis. At the jump, this current propagates radially and uniformly (in its axial distribution) giving a current density of approximately $j = I / 2\pi R U_s \Delta t = 1.3 \times 10^7 \text{ A/m}^2$, where $U_s t$ is the shock thickness (3 cm). Assume that the plasma particles stay in the current sheet for a time $\Delta t = 0.13 \mu\text{sec}$ (equal to the time for the shock front to pass a stationary point in the lab frame) after first being swept up by the advancing shock front.

The power expended in resistive heating of a volume ΔV is given by $P = \frac{j^2 \Delta V}{\sigma_{\perp}}$. The classical value of conductivity, given by $\sigma = T^{3/2} / 6.53 \times 10^8 \text{ mho/m}$ (T in $^{\circ}\text{K}$, and $\ln \Lambda \sim 10$), has to be replaced by the transverse conductivity when the current is perpendicular to a strong magnetic field. This is because the conducting electrons tend to stay in small (Larmor radius) orbits about the magnetic field lines. The transverse conductivity is given by $\sigma_{\perp} = \frac{\sigma}{(1 + \omega_{ce}^2 \tau_{ei}^2)}$, (5.8) where ω_{ce} is the electron cyclotron frequency and τ_{ei} is the collision frequency for electrons with ions.

$$(\omega_{ce} \tau_{ei})^2 = 2.36 \times 10^{33} \frac{B^2 T^3}{n^2 \ln^2 \Lambda}, \quad (5.9)$$

where B is the magnetic field in Teslas, T is in $^{\circ}\text{K}$, n is density in m^{-3} , and $\ln \Lambda \sim 10$. This gives

$$\sigma_{\perp} = 6.5 \times 10^{-37} \frac{n^2 \ln^2 \Lambda}{B^2 T^{3/2}} \text{ mho/m}. \quad (5.10)$$

For the case at hand we take B to be the average field in the jump, ~ 0.61 T, $n \sim 1.5 \times 10^{22} \text{ m}^{-3}$. Hence,

$$P = 1.54 \times 10^{36} \frac{j^2 B^2 T^{3/2}}{n^2 \ln^2 \Lambda} \Delta V \text{ Watts} \quad (5.11)$$

This is set equal to the energy necessary to heat $n \Delta V$ electrons to a temperature T in unit time. Thus,

$$\frac{n \Delta V k T}{\Delta t} = \frac{j^2}{\sigma_1} \Delta V, \text{ or}$$

$$T^{1/2} = \frac{6.5 \times 10^{-37} k n^3 \ln^2 \Lambda}{j^2 B^2 \Delta t} \quad (T \text{ } ^\circ\text{K}) \quad (5.12)$$

In the case of figure 4.3, using the values above, we obtain for the temperature jump in the shock front,

$$T = 1.4 \times 10^5 \text{ } ^\circ\text{K} = 12 \text{ eV.}$$

This compares well with the observed 10 eV.

Temperatures of similar magnitude follow from the same examination of the other transverse ionizing shock profiles. The closeness of this temperature estimate to the measured value indicates that relatively little heating of electrons in the shock plasma has occurred due to collisions with the shock heated ions.

5.3.b) Adiabatic compression.

Adiabatic

compression as a contribution to initial electron heating is negligible for this plasma. Adiabatic compression implies $Tn^{1-\gamma}$ constant, where γ is the ratio of specific heats. Thus across the shock front, we have $T_2 = T_1 \left(\frac{n_2}{n_1} \right)^{\gamma-1}$, or for the maximal compression observed in this experiment (~ 7) and $\gamma = 5/3$ (for a monatomic gas) an increase of about four times over the initial temperature of the deuterium to 0.1 eV.

Here we consider the possible role of adiabatic compression and expansion in relation to the plasma behavior following resistive heating of the electrons in the shock front. In section 5.2.b, it was noted that relatively little heating of the electrons is due to collisions which transfer energy from the ions to the electrons. Initially the electrons are heated in resistive dissipation at the shock front, and then there is a drop in T_e . In following the graphs of (particularly) figures 4.4 and 4.5, there seems to be an increase in T_e when there is one in n_e as time increases. This suggests a functional relation between T_e and n_e . Figure 5.3 shows a plot of T_e on a semilogarithmic scale versus n_e for $B_1 \sim 1$ kG and $P_1 \sim 100$ mTorr (Figure 4.5 is the source of the plot point values), for incident plasma, reflected plasma, and intermediate rarefaction regions. If one seeks a relation of the form $T_e \propto n_e^{\gamma-1}$, this graph gives $\gamma = 2.94$, which corresponds to 1 degree

Figure 5.3

of freedom (γ is related to the number of degrees of freedom of the constituent particles of a gas by

$$\gamma = \frac{m+2}{m}, \text{ where } m \text{ is the number of degrees of freedom.})$$

for the plasma electrons. (A plasma in a very strong magnetic field would have $m=1$.) A similar plot for the plasma of figure 4.4 gives $\gamma \sim 1.43$ to 1.75 , or $m \sim 2.5$ to 4.5 , in a fairly good curve fit bounded by these values. For figure 4.6, the T_e versus n_e plot gives $m \sim 3 \pm 1$ with a poor fit of the plot points. The significance of this consideration of temperature and density is not the actual exponent values, but the fact that there appears to be a relationship between the temperature and density of electrons throughout the plasma observation time. The electrons do not appear to be heated, other than by compression, or cooled, other than by rarefaction, except at the incident and reflected shock fronts, where Joule heating occurs.

5.3.c) Magnetic field and density compression. Ideally an MHD shock would conserve the quantity (B/ρ) because of frozen flux. This is not so for the ionizing shock case due to changes in chemical composition across the shock front. For the limiting case of high shock speed both the density compression and the magnetic field compression in an ionizing shock approach 4.

Figure 5.4 shows a plot of plasma electron density compression versus magnetic field compression. The electron density compression is defined as the electron density in the plasma divided by the density of bound electrons in the un-ionized initial gas. Magnetic field compression is given in terms of the ratio to the bias field (An exception in the plot is point number 10, for which the ratio B_5/B_2 is indicated -- the compression between incident and reflected shocks).

The numbered plot points correspond to the following conditions:

	P_1	B_1	U_s
(1) incident shock	50 mTorr,	3.6 kG,	23 cm/ μ sec
(2) reflected shock	"	"	"
(3) incident shock	"	1.0 kG	12.5 cm/ μ sec
(5) " "	100 mTorr	3.6 kG	16 cm/ μ sec
(6) reflected "	"	"	"
(7) incident "	"	1.0 kG	9 cm/ μ sec
(8) reflected "	"	"	"
(10) " "	50 mTorr	0	7 cm/ μ sec

The large error bars are on the plot because there is wide variation of n and B over the entire plasma region included in the plot point.

The graph shows that for different initial conditions the ratio of compressions is not the same. Higher

Figure 5.4 See section 5.3.c for the conditions applicable to each numbered plot point.

MAGNETIC AND DENSITY COMPRESSION
(n_{e1} ARE BOUND ELECTRONS)

compression ratios are evident for the faster transverse ionizing shocks. For these shocks the density compression is larger than the magnetic field compression. For the slower shocks, on the other hand, the magnetic field compression is greater than the density compression.

5.4 THE SCATTERED LIGHT SPECTRAL DISTRIBUTION AND ITS RELATION TO THE ELECTRON VELOCITY SPECTRUM

It was not intended to study the spectral distribution of scattered light, other than to obtain the half width of the profile in order to evaluate T_e . However, in the course of data evaluation occasional irregularities were noticed. Most of these irregularities could be corrected upon testing system components, but a few seemed to persist even though they could possibly be explained away by asserting that the apparent deviation lies within the limits of error for the profile points.

Figure 4.2 shows four examples of scattered light spectral profiles from Thomson scattering from the plasma. Consider the location of the plot points, and ignore the shape of the curves drawn in the figure which tend maybe to reflect the author's whims in the slight upcurving on the wings of the curve. Almost all shots had signals on five channels, but relatively few shots showed Thomson scattering signal on all six channels. Very few showed significant deviation from a Gaussian plot-- drawings (a) and (b) of figure 4.2 are typical. Drawings (c) and (d) are two good examples of atypical traces which led to two questions about the electron temperature data obtained.

(1) Is there really an asymmetry in some plots as shown in (d) in which there appears to be a slight shift to the short wavelengths? Can it be a drift motion or is it

an instrumental or theoretical difficulty? (2) Do the wings of the plot curve tend to spread slightly more than they might for a Gaussian plot-- is there a two electron temperature distribution with a Gaussian profile for each temperature, or does correlated scattering cause an enhancement of the centerline signal.

The first question may be answered by considering the apparent deviation to fall within the error limits of plot point evaluation. However, as in figure 4.2.(d) this was observed^{or} many of the higher temperature shots and exchange of photomultipliers, shift of spectrograph channels, and comparing data taken with the shock tube bias and main condenser bank polarities reversed did not change the sense of the apparent deviation. This eliminates radial or azimuthal electric field drifts which would change direction upon current reversal and shift the light toward the red. The average apparent shift for temperatures in the range 25 to 50 eV is between 0 and 10 Å toward shorter wavelengths. Higher resolution (closer channel spacing) and smaller channel bandwidths would be required to ascertain the actual value (if not zero) of such a shift. An interesting analytical explanation has been offered by J. Sheffield⁴³, who calculates a shift to the blue, due to inclusion of terms of order v/c (usually taken equal to 0) in the evaluation of uncorrelated scattered spectra, by the amount $\Delta\lambda = 1.0 \times 10^{-1} T_e$ Angstroms (T_e in eV) for 90° scattering with

ruby laser light. For $T_e = 30$ eV, the deviation according to this formula would be $3 \text{ \AA}$. Again the few observations of this shift do not numerically confirm any theory.

The other question is more important, and it is fairly certain that center line enhancement has occurred due to collective scattering from a plasma where the Salpeter parameter $\alpha \approx \frac{1}{k} \lambda_D$ (see appendix C) is near 1. Figure 4.2 (c) shows an extreme example of this. The half width of the profile as drawn would imply a temperature of 4 eV; however, if one uses a semilogarithmic plot of the two nearest channels on either side of the center line a good fit for 8 eV results. This is the real electron temperature; however, the shot could not be used to evaluate density accurately (it is about 4×10^{16}), because to do so would require information of finer spectral resolution or knowledge of T_1 .

A few plots for which α was much less than 1 showed points on the wings above the Gaussian plot level. There were indeed close to the noise level $150 \text{ \AA}$ or more off the center line, but could indicate the presence of a higher temperature for some of the electrons. The relative densities could not be evaluated precisely, but less than a tenth of the electrons could have contributed in any such case and the higher temperature of this small group of electrons must be less than an order of magnitude above the main plasma temperature.

5.5 CONCLUSIONS AND SUMMARY OF RESULTS

In this section we summarize the results of this thesis. As topics for possible related investigations arise in this summary, they will be discussed.

The principal results of this research are:

(1) Thomson scattering measurements of electron temperature and density have been made on a coaxial electromagnetically driven shock tube in which incident and reflected shock produced plasmas in deuterium were studied. Profiles of T_e and n_e have been obtained covering the lifetime of plasma in the shock tube.

(2) Throughout the shock plasma, the electron temperature is observed to be less by an order of magnitude than the ion temperature, for strong (9 to 25 cm/ μ sec) transverse ionizing shocks. This supports the idea that collisional ionizing shocks should initially heat ions much more than electrons; however,

the plasma (from strong transverse ionizing shocks) studied here does not undergo Coulomb collisional equipartition of thermal energy-- the electrons remain colder. This is a question on which there is little theoretical information, because in the present case, all the available theory indicated that such equilibration of electron and ion temperatures should occur. The failure of Coulomb collisional relaxation to occur in this plasma requires further investigation in detail.

(3) For the case of no initial magnetic field, the agreement of theory with the results of this experiment is excellent. The gas shocks studied are relatively slow (3 to 7 cm/ μ sec) and their behavior is well described by simple theory and by more sophisticated computer calculations -- the time and distance scales, as well as the absolute magnitude of the shock plasma quantities, fit well.

(4) The magnitude and time behavior of the electron heating in this plasma (transverse ionizing shocks) are well described by Joule heating of electrons in the shock front and piston. The electron behavior appears to be independent (for the most part) of the ion temperature. It is found that the density compression becomes relatively larger compared to the magnetic field compression for faster shocks.

(5) Possible deviations of the scattered light spectrum from a Gaussian shape have been considered. Correlation effects have been observed. When this experiment was originally conceived, it was expected that significantly higher electron temperatures would be observed than have been found to be the case. Hence it was expected that the scattering would be due to electron density fluctuations only. Ion density fluctuation contributions to the scattered spectrum have been observed at the scattering angle used in this experiment. A forward scattering (15° to 20° scattering angle) experi-

could be carried out, although with considerably more difficulty than the present experiment, and the ion temperature measured.

APPENDIX A

Shock Jump Equations

The jump equations are used to obtain the final state values of the post shock variables in terms of the same variables in the unshocked gas. The initial gas quantities are designated by the subscript 1, post-shock equilibrium quantities by the subscript 2, and piston quantities by 3. References 15 and 25 may be consulted for further details on these equations which are simply listed here.

- Assumptions:
- (1) steady flow ($\partial/\partial t \equiv 0$),
 - (2) one dimensional flow,
 - (3) p_1, ρ_1, B_1 , and the drive field $B_3 - B_1$ are initial values,
 - (4) $\vec{E}$, $\vec{v}$, and $\vec{B}$ are mutually perpendicular,
 - (5) ignore boundary layers, and neglect radiation and heat flow,
 - (6) assume the post-shock plasma to be fully equilibrated, and
 - (7) the post-shock plasma is force free (i.e., free of electric current).

Given the above assumptions, we can obtain the following

jump equations:

All the following equations are in the laboratory frame. The flux conservation relations are derived in the shock front rest frame, but here converted to lab frame velocities. CGS Gaussian units are used.

- (1) Conservation of mass flux:

$$\rho_2 (V_s - V_2) = \rho_1 V_s \quad A-1$$

- (2) Conservation of momentum flux:

$$(P_2 + B_2^2/8\pi) - (P_1 + B_1^2/8\pi) = \rho_1 V_s V_2 \quad A-2$$

(3) Conservation of energy flux:

$$(P_2 + B_2^2/8\pi) V_2 = \frac{c}{4\pi} E_1 B_1 + \rho_1 V_3 \left[(\varepsilon_2 + B_2^2/8\pi\rho_2) - (\varepsilon_1 + B_1^2/8\pi\rho_1) \right] + \frac{V_2^2}{2} \rho_1 V_3 \quad A-3$$

(4) Equation of state:

$$\text{For initial gas: } \varepsilon_1 = (\gamma_{D_2} - 1) \frac{P_1}{\rho_1}, \quad \gamma_{D_2} = 7/5 \quad A-4a$$

$$\text{For plasma: } \varepsilon_2 = \frac{3}{2} \frac{P_2}{\rho_2} + \frac{D}{2m_D} + \alpha_2 I/m_D \quad A-4b$$

(D is the dissociation energy of D_2 ; I is ionization energy; $\alpha \equiv n_e/(n_e + n_{D_2})$)

(5) Faraday's law (electric field condition):

$$E_1 = \frac{1}{c} V_3 B_1 \left[1 - \frac{B_2/B_1}{\rho_2/\rho_1} \right] \quad A-5$$

(6) Piston pressure balance:

$$B_3^2/8\pi = P_2 + B_2^2/8\pi. \quad A-6$$

Consequences: If $P_1 \ll P_2$, and $\varepsilon_1 \ll \varepsilon_2$, then

$$\frac{\rho_2}{\rho_1} \left(\frac{D}{m_D} + \frac{2I}{m_D} \right) = U_m^2 \left(\frac{\rho_2}{\rho_1} - 4 \right) + \frac{B_1^2}{8\pi\rho_1} \left(\frac{B_2}{B_1} + 5 \right) \left(\frac{B_2}{B_1} - 1 \right) \quad A-7$$

where $U_m^2 \equiv V_3 V_2$. From this equation we may then,

upon substituting for P_2 , obtain a value for T_2

(the last term may be dropped for M large):

$$kT_2 = \frac{m_D}{2(\rho_2/\rho_1)} \left\{ U_m^2 - \left[\frac{B_2^2}{B_1^2} - 1 \right] \frac{B_1^2}{8\pi\rho_1} \right\} + \frac{\rho_1}{2\rho_2} kT_1. \quad A-8$$

APPENDIX B.

Spectral Distribution of Density
Fluctuations in a Maxwellian Plasma

For a Maxwellian distribution of velocities of both ions and electrons (T_e not necessarily equal to T_i), neglecting particle collisions, assuming that the Coulomb interaction energy is small compared to the thermal kinetic energy, and that the observation time is long compared with the period of the fluctuations, the general form of $S(\omega, \bar{K})$ as given by Salpeter³⁶ is:

$$S(\omega, \bar{K}) = \frac{n}{K} \left\{ \left| \frac{1 - G_i}{1 - G_i - G_e} \right|^2 F_e\left(-\frac{\omega}{K}\right) + Z \left| \frac{G_e}{1 - G_i - G_e} \right|^2 F_i\left(-\frac{\omega}{K}\right) \right\}.$$

Here the electron density fluctuations have frequency

$\omega = \omega_i - \omega_s$ and wave vector $\bar{K} = \bar{K}_i - \bar{K}_s$ (where i, s refer to incident and scattered quantities); $F_e(v)$ and $F_i(v)$ are the Maxwell velocity distribution functions for electrons and ions-- $F(v) = \sqrt{\frac{m}{2\pi kT}} \exp\left(-\frac{mv^2}{2kT}\right)$; Z is ion charge; and, $G_e(\omega) \equiv -\alpha^2 [1 - f(x) + i\sqrt{\pi}x e^{-x^2}]$ and $G_i(\omega) \equiv -\frac{ZT_e}{T_i} \alpha^2 [1 - f(y) + i\sqrt{\pi}y e^{-y^2}]$, with $x \equiv \omega/\omega_e$, $y \equiv \omega/\omega_i$, $\omega_e \equiv \sqrt{\frac{2K^2 kT_e}{m_e}}$, $\omega_i \equiv \sqrt{\frac{2K^2 kT_i}{m_i}}$, $f(x) \equiv 2x e^{-x^2} \int_0^x e^{t^2} dt$, and $\alpha \equiv 1/K\lambda_D$ (λ_D is the Debye length of the plasma).

The first term of this equation for $S(\omega, \bar{K})$ is a relatively broad spectral distribution corresponding to a Doppler spread characteristic of electron thermal motion while the second is a narrower distribution cor-

responding to ion thermal motion. The total integrated intensity of the electron and ion components are

$$1/(1+\alpha^2) \quad \text{and} \quad Z\alpha^4/[(1+\alpha^2)(1+\alpha^2+Z\alpha^2\frac{T_e}{T_i})], \quad \text{respectively.}$$

This equation for $S(\omega, \bar{k})$ breaks down into three cases for the value of the parameter $\alpha \equiv \frac{1}{k\lambda_D} = \frac{\lambda_i}{\lambda_D(4\pi\sin\frac{\theta}{2})}$, where θ is the scattering angle.

(i) $\alpha \ll 1$. This corresponds to the case where the incident wavelength λ_i is much less than a Debye length $\lambda_D(\text{cm}) = 744 \sqrt{T_e(\text{eV})/n_e(\text{cm}^{-3})}$, or more precisely, less than $\lambda_D \times 4\pi\sin\frac{\theta}{2}$. In this case the electron term dominates, and the scattered spectrum has the Gaussian spectral distribution characteristic of the random electron thermal motion.

(ii) $\alpha \sim 1$. The electron spectrum becomes flat topped for $\alpha \sim 1$, and as α increases further the spectrum is double humped. A narrow ion peak appears in the center of the distribution. It is possible, using the known shape of spectra as a function of α , to determine both T_e and T_i from the one spectrum ⁵¹.

(iii) $\alpha \gg 1$. This corresponds to collective or correlated scattering from the collective interactions of the plasma. Here, $\lambda_i \ll \lambda_D$, so that the scattering is determined by the motion of the ions about which the electrons move in clouds of radius λ_D . The bulk of the scattered intensity is in the narrow ion center peak whose characteristic width depends on $\beta \equiv Z\sqrt{\frac{T_e}{T_i}}$, just as the electron peak depended on α for small α . Electron satellite lines now occur at $\omega \approx \pm\omega_p = \sqrt{\frac{4\pi ne^2}{m_e}}$ corresponding

to scattering from electron plasma oscillations.

The width of the satellites is determined only by Landau damping, so they may be very narrow.

APPENDIX C

An estimate of the ratio of Thomson scattered signal to plasma bremsstrahlung

The scattered light collection geometry is shown in figure 3.3. The scattering volume viewed through the pinhole in the collecting optics is approximately 0.0023 cm^3 (a cylinder of diameter 1.3 mm and length 2 mm).

An effective volume of bremsstrahlung emitting plasma as viewed by the lens is 1.4 cm^3 (a cylinder $\frac{1}{4}$ inch in diameter by 5 cm long). The solid angle viewed from the scattering volume and subtended by the lens is .005 sr (laser focal spot to collecting lens is 8 cm, lens aperture is $\frac{1}{4}$ inch). The emission coefficient for bremsstrahlung is

$$\epsilon_{\lambda} = 1.51 \times 10^{-37} \frac{n_+ n_e Z^2}{\lambda^2 \sqrt{T_e}} e^{-1.24 \times 10^{-4} / \lambda T_e} \frac{W}{\text{cm}^3 \text{cm}(\lambda) \text{sr}}$$

for wavelengths centered about λ cm, where n_+ and n_e are the ion and electron densities in cm^{-3} , T_e is the electron temperature in eV, and Z is the ionic charge. For typical values of this experiment, $n_+ = n_e \sim 2 \times 10^{16} \text{ cm}^{-3}$, $T_e \sim 15 \text{ eV}$, $Z = 1$, $\lambda \sim 7 \times 10^{-5} \text{ cm}$, channel width $\sim 20 \text{ \AA} = 2 \times 10^{-7} \text{ cm}$, emitting volume $\sim 1.4 \text{ cm}^3$, and solid angle $\sim .0057 \text{ sr}$. Thus the power from bremsstrahlung into one channel at the entrance to the collecting optics is $1.3 \times 10^{-6} \text{ Watts}$.

Total Thomson scattered light into the collecting optics for a 500 MW beam focussed down to a diameter of 1.3 mm

is given by $I_0 \sigma_T n \frac{\delta \Omega}{4\pi} V \approx 6.7 \times 10^{-4} \text{ W}$. $T_e \sim 15 \text{ eV}$

gives $\Delta \lambda_{\frac{1}{2}} \sim 125 \text{ \AA}$ for the scattered light profile half

width, which would be spanned by six adjacent channels of $20 \text{ \AA}$ half width, each getting approximately one sixth of the total scattered light, if the profile were imagined as a step function of width $125 \text{ \AA}$. This gives approximately $1 \times 10^{-4} \text{ W}$ per channel, or a ratio

$$\frac{\text{scattered light into one channel}}{\text{plasma bremsstrahlung into same channel}} \approx 80$$

at the entrance to the collection optics. This ratio remains unchanged at the output of the polychromator, although absolute levels may be down by a factor of approximately 25. This attenuation level is acceptable as the photomultipliers can easily detect the remaining signal.

REFERENCES

1. R.A. Gross, Y.G. Chen, E. Halmoy, and R.R. Moriette, Strong Shock Waves, *Phys. Rev. Lett.* 25, 575 (1970).
2. Y.G. Chen, C.K. Chu, R.A. Gross, E. Halmoy, P. Moriette, and S. Schneider, Plasma Heating by Strong Shock Waves, paper CN-28/J-8 of Proceedings of Fourth Conference on Plasma Physics and Controlled Nuclear Fusion Research, Vienna, IAEA (1971).
3. R.A. Gross, A Shock Heated Wall-confined Fusion Cycle, *Nuclear Fusion* 13, 293 (1973).
4. F.L. Ribe, Fusion Reactor Systems and Associated Technological Problems, Review paper given at APS Plasma Physics Section, November 1973, *Bull. APS*, Ser. II, 18, 1288 (1973).
5. C.K. Chu, and R.A. Gross, Shock Waves in Plasma Physics, article in Advances in Plasma Physics, vol. 2, A. Simon and W.B. Thompson, eds., N.Y., J. Wiley (1969).
6. H.J. Kunze, The Laser as a Tool for Plasma Diagnostics, article in Plasma Diagnostics, W. Lochte-Holtgreven, ed., Amsterdam, North Holland (1968).
7. R.A. Gross, Strong Ionizing Shock Waves, *Rev. Mod. Phys.* 37, 724 (1965).
8. R.A. Gross, The Physics of Strong Shock Waves in Gases, article in Proceedings of the International School of Physics Enrico Fermi, XLVIII Corso, Physics of High Energy Density, Varenna, Italy; New York, Academic Press (1971).
9. A.W. DeSilva and G.C. Goldenbaum, Plasma Diagnostics by Light Scattering, part III of Methods of Experimental Physics (Plasma Physics), vol. IX, part A, N.Y., Academic Press (1970).
10. S.A. Ramsden, Light Scattering Experiments, article in Physics of Hot Plasmas, B.J. Rye and J.C. Taylor, eds. Scottish Universities' Summer School, 1968; Edinburgh, Oliver and Boyd (1970)
11. D.E. Evans and J. Katzenstein, Laser Light Scattering in Laboratory Plasmas, *Rep. Prog. Phys.* 32, 207-271 (1969).

12. Ya.B. Zel'dovich and Yu.P. Raizer, Physics of Shock Waves and High Temperature Hydrodynamic Phenomena, (in Russian), Moscow, 1966. English translation in two volumes, New York, Academic Press (1966, 1967).
13. I.G. Katayev, Electromagnetic Shock Waves, (in Russian), Moscow, 1963. English translation, London, Iliffe (1966).
14. R. Courant and K.O. Friedrichs, Supersonic Flow and Shock Waves, New York, Interscience (1948).
15. R.J. Meehan, Experimental Investigation of Incident and Reflected Transverse Ionizing Shock Waves, Thesis, Columbia University, 1969. Printed as Columbia University Plasma Physics Laboratory report No. 44, February, 1969.
16. R.Z. Sagdeev, Cooperative Phenomena and Shock Waves in Collisionless Plasmas, in Reviews of Plasma Physics, vol. 4, M.A. Leontovich, ed., (in Russian) Moscow, 1964. English translation, New York, Consultants Bureau (1966).
17. H. Völk, Collisionless Shocks, article in Physics of Hot Plasmas, B.J. Rye and J.C. Taylor, eds., Scottish Universities Summer School, 1968. Edinburgh, Oliver and Boyd (1970).
18. J.W.M. Paul, Collisionless Shock Waves, article in Physics of Hot Plasmas, cf. ref. 17.
19. J.B. McBride, E. Ott, and C.E. Wagner, Theory of Turbulent Plasma Heating by Anomalous Absorption of Magnetosonic Waves, Nuclear Fusion 13, 151 (1973).
20. V.D. Shafranov, Structure of Shock Waves in a Plasma (in Russian), ZhETF 32, 1453 (1957).
21. S.H. Robertson, Ion Temperature Measurements of Shock Heated Plasmas, Bull. APS, ser. II, vol 18, 1347 (1973).
22. S. Glasstone and R. Lovberg, Controlled Thermo-nuclear Reactions, Princeton, N.J., D. Van Nostrand (1960).
23. R.T. Taussig, Plasma Heating by Reflected Ionizing Shock Waves, Phys. Fluids 13, 2270 (1970).

24. H.C. Lui, Numerical Studies of Shock Focusing in a Coaxial Electromagnetic Shock Tube, Thesis, Columbia University, 1973. Published as Columbia University Plasma Laboratory Report No. 60 (1973).
25. S.H. Schneider, Numerical Study of Strong Plasma Shock Waves Produced in an Electromagnetic Shock Tube, Phys. Fluids 15, 805 (1972).
26. S.H. Schneider, Numerical Simulation of Strong Plasma Shock Waves Produced in an Electromagnetic Shock Tube, Thesis, Columbia University, 1971. Published as Columbia University Plasma Laboratory Report No. 55 (1971).
27. S.H. Schneider, C.K. Chu, and B.P. Leonard, Numerical Study of the Magnetic Shock Tube, Phys. Fluids 14, 1103 (1971).
28. E. Halmoy, Production and Investigation of Very Strong Ionizing Shock Waves, Thesis, Columbia University, 1971. Published as Columbia University Plasma Laboratory report No. 53 (1971).
29. E. Halmoy, Production and Investigation of Very Strong Ionizing Shock Waves, Phys. Fluids 14, 2134 (1971).
30. P.R. Moriette, A Study of Strong Transverse Ionizing Shock Waves, Thesis, Columbia University, 1971. Published as Columbia University Plasma Laboratory report No. 56 (1971).
31. P.R. Moriette, Experimental Study of Strong Transverse Ionizing Shock Waves, Phys. Fluids 15, 51 (1972).
32. R.T. Taussig, Y.G. Chen, and R.A. Gross, Strong Shocks, Snowplows, and Magnetic Pistons, Phys. Fluids 16, 212 (1973).
33. Y.G. Chen, Density and Magnetic Field Compression by Strong Transverse Shock Waves, Bull. APS, Ser. II, vol. 18, 1347 (1973).
34. D.H. McNeill, Thomson Scattering Measurements of Electron Temperature and Density in a Shock Heated Plasma, Bull. APS, Ser. II, vol. 18, 1347 (1973).
35. B. Feinberg, A Study of Energy Transfer from a Hot Plasma in Contact with a Cold Wall, Bull. APS, Ser. II, vol. 18, 1348 (1973).

36. E.E. Salpeter, Electron Density Fluctuations in a Plasma, Phys. Rev. 120, 1528 (1960).
37. J.A. Fejer, Scattering of Radio Waves by an Ionized Gas in Thermal Equilibrium, Can. J. Phys. 38, 1114 (1960).
38. J.P. Dougherty and D.T. Farley, A Theory of Incoherent Scattering of Radio Waves by a Plasma, Proc. Roy. Soc. A, 259, 79 (1960).
39. M.N. Rosenbluth and N. Rostoker, Scattering of Electromagnetic Waves by a Nonequilibrium Plasma, Phys. Fluids 5, 776 (1962).
40. A.I. Akhiezer, I.G. Prokoda, and A.G. Sitenko, On Scattering of Electromagnetic Waves in a Plasma, (in Russian) ZhETF 33, 750 (1957).
41. I.B. Bernstein, S.K. Trehan, and M.P.P. Weenink, Plasma Oscillations:II. Kinetic Theory of Waves in a Plasma, Nuclear Fusion 4, 61 (1964).
42. G. Bekefi, Radiation Processes in Plasmas, N.Y., J. Wiley (1966).
43. J. Sheffield, The Incoherent Scattering of Radiation from a High Temperature Plasma, Plasma Physics 14, 783 (1972).
44. D.E. Evans, The Effect of Impurities on the Spectrum of Laser Light Scattered by a Plasma, Plasma Physics 12, 573 (1970).
45. J. Sheffield, A Method for Measuring the Direction of the Magnetic Field in a Tokamak Plasma by Laser Light Scattering, Plasma Physics 14, 385 (1972).
46. W.H. Kegel, Light Scattering from Plasmas with a Non-Maxwellian Velocity Distribution, Plasma Physics 12, 295 (1970).
47. G. Fiocco and E. Thompson, Thomson Scattering of Optical Radiation from an Electron Beam, Phys. Rev. Letters 10, 89 (1963).
48. R.A. Niland and C.N. Watson-Munro, Electron Temperatures behind Normal Ionizing Shock Waves in Helium, Physics Letters 34A, 91 (1971).
49. L. Bighel, A.R. Collins, R.C. Cross, I.S. Falconer, and C.N. Watson-Munro, MHD Shock Wave Heating of a Plasma, Bull. APS, Ser. II, vol. 18, 1348 (1973).

50. R.M. Patrick, Thomson Scattering Measurements of Magnetic Annular Shock Tube Plasmas, Phys. Fluids 8, 1985 (1965).
51. W.H. Kegel, Graphs for the Determination of Plasma Parameters by Light Scattering Experiments, Report IPP-6/34 Institut für Plasma Physik, Garching, March 1965 (in German). Translation by J.B. Sykes, Culham Laboratory report CTO/760.
52. R.A. Gross and B. Miller, Plasma Heating by Strong Shock Waves, article in Methods of Experimental Physics (Plasma Physics), vol. IX, part A, H.R. Griem and R.H. Lovberg, eds., N.Y., Academic Press (1970).
53. Yu. V. Makarov and E.K. Chekalin, Physical Processes in Electromagnetic Shock Tubes, Atomizdat, Moscow, 1968 (in Russian).
54. R.S. Post and Y.G. Chen, A 100kV, Fast, High Energy, Non-uniform Field Distortion Switch, Rev. Sci. Inst. 43, 622 (1972).
55. R.C.A., Photomultiplier Manual, Technical Series PT-61. Harrison, N.J., RCA Electronic Components (1970).
56. H.R. Griem, Plasma Spectroscopy, New York, McGraw-Hill (1964).
57. W. Engelhardt, Continuum Measurements on a Hydrogen Plasma, Phys. Fluids 15, 2074 (1972).
58. L.A. Artsimovich, Controlled Thermonuclear Reactions, Fizmatgiz, Moscow, 1963 (in Russian).
59. H.G. Gale, G.S. Monk, and K.O. Lee, Wave-lengths in the Secondary Spectrum of Hydrogen, Astrophysical Journal 67, 89 (1928).
60. M. Martone and S.E. Segre, Electron Heating in Large Amplitude Oscillations in a Theta Pinch, Plasma Physics 13, 173 (1971).
61. R.T. Taussig, R.A. Chervin, and Y.G. Chen, Field Annihilation in Reversed Field Current Layers, Bull. APS, Ser. II, vol. 18, 1292 (1973).
62. L. Spitzer, Jr., Physics of Fully Ionized Gases, 2nd edition, N.Y., J. Wiley (1962).

63. M.Y. Jaffrin and R.F. Probst, Structure of a Plasma Shock Wave, *Phys. Fluids* 7, 1658 (1964).
64. D.V. Sivukhin, Coulomb Collisions in a Fully Ionized Plasma, article in Reviews of Plasma Physics vol. 4, M.A. Leontovich, ed., Moscow, 1964 (in Russian). English translation by Consultants Bureau, N.Y. (1966).