SPECTROSCOPIC DIAGNOSTICS OF LASER PRODUCED PLASMAS

BY

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Abstract of Dissertation Presented to the Graduate School of the University of Florida in Partial Fulfillment of the Requirements for the Degree of Doctor of Philosophy

SPECTROSCOPIC DIAGNOSTICS OF LASER PRODUCED PLASMAS

Norman David Delamater

December 1984

Chairman: Charles F. Hooper, Jr.
Major Department: Physics

A detailed phenomenological study of the spectroscopic analysis of laser produced plasmas is presented in this dissertation. High resolution X-ray spectroscopic data have been obtained in pellet implosion experiments using the 24 beam Omega laser system at the University of Rochester Laboratory for Laser Energetics. Line broadening theory was applied to diagnose plasma core temperatures and densities from observed spectral line shapes; detailed atomic physics models were constructed to diagnose core parameters from sensitive line intensity ratios. Previous experiments have been designed to minimize the effects of plasma opacity, which tends to distort both the lineshape and line ratio diagnostics. This study attempts to systematically observe the effects of opacity on the diagnostics by a self consistent analysis of X-ray spectra obtained from implosions of argon and neon filled glass microballoons differing in diameter, wall thickness, and Ar/Ne fill ratio. The subsequent analysis then includes opacity effects through the use of a radiation transport model for the spectral lines;
atomic physics models are studied and the final results systematically compared with the experimental results.

Typically, the X-ray spectra from dense, high temperature plasmas show prominent satellite lines arising from doubly-excited autoionizing states which lie near the strong resonance lines of hydrogenic and helium-like ions. These satellites are well resolved spectroscopically and are seen to vary in relative intensity over the range of densities achieved in these experiments. A model is presented which has been used to calculate the density dependence for several important satellite line ratios and hence to diagnose electron densities in these experiments. For the first time, a detailed comparison is presented between these satellite line diagnostics and a line broadening analysis including the effects of opacity.

The results of this analysis have relevance for implosive and non-implosive plasma diagnostics. For inertial confinement fusion plasmas, the spectroscopic information represents the critical feedback validating hydrodynamic and atomic physics models, as well as atomic data. Current models of pellet implosions which couple hydrodynamics and radiation transport are highly uncertain; the constraints placed on such models by the present results are discussed.
CHAPTER I
EXPERIMENT

Glass microballoons of various sizes were filled with mixtures of argon and neon and imploded using the 24-beam Omega Nd-glass laser system at the University of Rochester Laboratory for Laser Energetics. Experimentally observed X-ray spectra and pinhole images were analyzed to determine plasma core parameters and to study the effects of opacity on spectroscopic diagnostics. The proportion of argon in the fill was varied from 5% to 100% with the intention of systematically observing the effects of opacity on radiation transport of the principal series lines of hydrogenic and helium-like argon. Electron temperatures and densities, and ion densities were determined by both detailed fits of theoretical lineshapes to experimental line profiles, and by the measurement of various line intensity ratios which are shown to be temperature or density sensitive. The prominent helium-like and lithium-like satellite lines situated respectively on the red wing of the hydrogenic and helium-like resonance lines were observed as additional plasma diagnostics; the results were compared with the previous methods.

The shots analyzed provided high resolution time-integrated spectral data in the region 3 - 4 keV region covering the range of hydrogenic and helium-like argon emission. Laser pulses of 1.06 microns, about 100 picoseconds in duration, and with an energy of approximately 600 joules per pulse were used to symmetrically irradiate
the glass targets. Pellets of sizes 100μm (diameter) × 1μm (wall thickness), 100μm × 3μm, 100μm × 5μm, and 150μm × 2μm were filled to 10 atmospheres total pressure with argon/neon fill ratios of 5/95, 25/75, 50/50, and 100/0 in order to systematically control the opacity in the argon lines. Table 1 provides a list of shot parameters for those shots providing useful data which were analyzed. A fairly complete variation of fill ratio has been obtained for the 100μm × 1μm, and the 100μm × 3μm targets, while less complete data were obtained for the other targets due to various experimental problems.

A flat X-ray crystal spectrometer was used to record the time-integrated plasma emission. The instrument resolution is determined using Bragg's law,

\[ \frac{\Delta E}{E} = \Delta \theta \cot \theta_B; \]  

1-1

if the angular spread \( \Delta \theta \) is determined primarily from the finite source size, then

\[ \frac{\Delta E}{E} = \frac{\Delta X}{L} \cot \theta_B, \]  

1-2

where \( \theta_B \) is the Bragg angle at energy \( E \), \( \Delta X \) is the linear source size, and \( L \) is the source to film distance. Using the source sizes determined with the X-ray microscope images, and the known crystal spacing and experimental geometry, the instrument resolution is determined to be approximately 3.5 eV. In the analysis, the theoretical lineshapes were convolved with a gaussian shape of full width at half maximum equal to 3.5 eV to approximately account for instrumental broadening. This was
not a major effect on the lineshapes studied since in this experiment most of them were already considerably larger than this instrumental width.

The spectra were recorded on Kodak 2497 film, and the film was microdensitometered at the Lawrence Livermore National Laboratory. The digitized neutral film density versus film distance was converted to relative intensity versus energy in order to accurately compare the data with theoretical lineshapes. The energy scale was established using both argon and nearby impurity lines from the glass shell constituents whose energies have been previously calculated and reported in the literature.\textsuperscript{1,2} A linear least squares fit was performed to determine the dispersion (eV/mm) in each spectral region. It was assumed that at the densities attained in these experiments line shifts due to plasma effects were negligible so that the line energies calculated for the isolated ions would be appropriate. The Kodak film used has been absolutely calibrated and the results reported in the literature.\textsuperscript{3,4} It has been shown that for sufficiently small film density, the film density is linearly proportional to exposure (intensity x time). As the film densities in this experiment do fall within the linear range and only relative intensities are used in the analysis, the conversion from film density to radiation intensity is linear and absolute intensities are not calculated.

Experimentally determined lineshapes and line intensity ratios must be compared with theoretical calculations in order to produce reliable plasma diagnostics. In Chapter II, several detailed plasma models are presented which were used to characterize the laser produced plasmas in this experiment. The models provide steady state ionic population
densities; in turn these then determine a number of line intensity ratios which depend on electron density, temperature, or both. Included in these plasma models are an analysis of autoionizing level populations which give rise to satellite line emission. Chapter III considers in detail the shapes of spectral lines emitted from dense, laser produced plasmas. Line broadening theory is used to calculate the intrinsic shapes for the principal series lines from hydrogenic and helium-like argon. These shapes must be modified in order to account for the radiation's escape from the plasma. An approximate solution to the problem of line transfer is presented. A method is presented which systematically compares theoretical lineshapes with all observable lines in a spectral series in order to determine electron and ion densities which are consistent with the opacity model.

Chapter IV presents the principal results of this analysis. Line broadening theory and line intensity ratios are used to predict plasma core parameters for a number of plasmas covering a range of densities, temperatures, and opacities. Chapter V discusses the significance of these results with respect to current models of pellet implosions.
### TABLE 1. Shot parameters

<table>
<thead>
<tr>
<th>shot no.</th>
<th>target size (μm)</th>
<th>Ar/Ne fill (atm)</th>
<th>laser pulse (ps)</th>
<th>absorbed energy (J)</th>
<th>Absorbed specific energy (J/μg)</th>
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<tr>
<td>7499</td>
<td>102.00x0.87</td>
<td>10/0</td>
<td>96.0</td>
<td>197.8</td>
<td>2.507</td>
</tr>
<tr>
<td>7560</td>
<td>105.00x0.93</td>
<td>5/5</td>
<td>116.0</td>
<td>135.7</td>
<td>1.577</td>
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<tr>
<td>7525</td>
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<td>2.5/7.5</td>
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<td>106.0</td>
<td>184.8</td>
<td>0.824</td>
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<td>125.7</td>
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</table>
CHAPTER II
PLASMA MODELS

During the last decade, high power laser systems have been used to drive pellet implosions in an attempt to produce plasma conditions sufficient to generate significant nuclear fusion reactions in the component hydrogen (D-T) fuel. Seed gases such as neon or argon have been used to provide diagnostic information of plasma core parameters and can modify the implosion dynamics by providing a mechanism for radiative cooling. An analysis of the line spectra emitted from these plasmas generally requires some hydrodynamic model for the pellet implosion and a detailed atomic physics model possibly including radiation transport to predict spectral emission.

Hydrodynamic simulation of spherically symmetric pellet implosions can be used to produce detailed predictions of the plasma development during the implosion and predict total ion densities, ion temperatures and electron temperatures as a function of space and time. Currently the detailed atomic physics models are computed in a post-process procedure using the hydro code results as input data for the overall ionization balance calculation. This procedure has been used to predict the overall spectral emission with some success for certain well characterized implosions. Major uncertainties exist in the models of laser energy deposition in the target, electron transport through the plasma, generation of non-thermal high energy electrons, and plasma
instabilities. When these effects become important, code predictions can become inconsistent with the observations.

The hydro code simulations have shown that laser driven implosions of microballoon targets result in "exploding pusher" or "ablative driven" implosion (or an intermediate type). The exploding pusher target generally has a thin (< 1 μm) glass shell which is subjected to a short (100 ps) intense laser pulse. The absorption of laser energy results in the production of hot electrons which instantaneously explode the shell, sending half the material outward and half the material inward thus compressing the gas fill. Due to the effects of the hot electrons and shock waves the gas fill can become preheated thus reducing the maximum compression. Prior to the disassembly phase, a stagnation phase is achieved in which the gas fill temperature is relatively high (several keV) and density low (< liquid DT) compared to the ablatively driven implosion. The ablative implosion is more desirable since shock and hot electron preheating are minimized and a slower compression is driven by the thermal conduction front, with the achievement of higher compression and lower temperature at stagnation. An ablative implosion is partially achieved by increasing shell thickness such that the hot electrons cannot explode the shell. Also reducing the laser intensity or operating wavelength has the effect of decreasing the hot electron production and creating a more ablative implosion. In the experiments to be described in this dissertation it is expected that implosions will be obtained which are of the "exploding pusher" type, however the thick shell targets should have an ablative character and be of an intermediate type.
An important prediction emerging from hydro code simulations is that spherically symmetric implosions reach a stagnation phase and the expected range of temperature and density is such that line emission from hydrogenic and helium-like argon ions should be prominent during this phase. This implies that a steady-state uniform plasma model can be used to predict the spectroscopic emission. Such models have been previously applied successfully to predict the behavior of neon-filled microballoon,\(^8,9\) and comparisons to a more detailed model show only small differences in the diagnosed parameters.\(^6,10,11\)

While a steady-state uniform plasma model may be adequate when time-integrated spectra is to be described, the more detailed simulations will be required when time-resolved spectra is analyzed. Equipment now becoming available is able to provide time resolution of \(~10\) picoseconds compared to implosion times of \(~\text{nanoseconds}\). These new, time-resolved observations should provide critical tests of the hydrodynamic codes and their inherent uncertainties.

In the experiments described here, since time-resolved spectroscopy is not being used, a steady state atomic physics model should be adequate to characterize the conditions in the plasma during peak emission. In particular, the spectral line emission will be analyzed to determine an electron temperature, \(T_e\), electron density, \(N_e\), and ion density, \(N_i\), characteristic of the plasma at peak emission. This analysis considers both lineshape and line ratio diagnostics. While lineshape diagnostics are primarily dependent on electron density, line ratio diagnostics depend on both \(N_e\) and \(T_e\) through the ionic level populations, the ratio of two spectral lines being given by
\[ \frac{I_{ij}}{I_{mn}} = \frac{N_i \nu_i A_{ij}}{N_m \nu_m A_{mn}} \]

where \( N \) indicates upper state level populations, \( \nu \) is energy of the transitions, and \( A \)'s are Einstein coefficients for the transitions. The level populations must be calculated from an atomic physics model.

**Local Thermodynamic Equilibrium**

The simplest model to assume is that of local thermodynamic equilibrium, or LTE. Particle collisions are dominant, rates of any collisional processes are exactly balanced by their inverse process, and ionic populations are functions of the local values of electron temperature and density. In fact all local ionic population densities are given by the Saha-Boltzmann equation

\[ N e^{-g Z+1} = (6.04 \times 10^{23}) N \frac{g Z+1}{g_z} \frac{e^{-I_z/kT}}{TeV} \]

which is valid for any two successive ionization stages where the \( g \)'s are the statistical weights and \( I_z \) is the ionization potential of stage \( Z \).

It is known that LTE is not a good approximation for all the ionic levels, but that it may be valid for some of the upper ones. An approximate criterion for a given level of principal quantum number \( n \) to be in LTE with the ground state of the next highest ionization stage is
\[ N_e > 7 \times 10^{18} \frac{z^7}{n^{17/2}} \left( \frac{kT}{Z^2 E_H} \right)^{1/2} \text{ cm}^{-3}, \]

which is obtained by comparing the total radiative decay rates from level \( n \) to the collisional deexcitation rate into level \( n \). The condition for the collision rate to be dominant gives the above relation. The assumption of LTE for some of the levels greatly simplifies the calculation of total level populations with a more detailed model.

**Non-LTE Steady-State Models**

A more general approach to determine ionic level populations in a plasma is to assume that overall plasma conditions are changing slowly compared to collisional and radiative timescales so that a steady-state condition is reached. Then the rate equations governing the population in each level of the model can be solved for steady-state populations, assuming all relevant collisional and radiative rate coefficients are known. This model, also called a collisional-radiative model, contains both the LTE model and the corona (low-density) model as limiting cases.

Consider the total populations \( N_j \) and \( N_{j+1} \) in two successive stages of ionization, \( j \) and \( j+1 \); then the solution of the rate equations gives for the population ratio,

\[
\frac{N_{j+1}}{N_j} = \frac{S_j}{\alpha_{j+1} + D_{j+1} + Ne_{j+1}}
\]

where \( S_j \) is the total electron collisional ionization rate from
j + j + 1, \( \alpha_{j+1} \) is the total radiative recombination rate into j, \( D_{j+1} \) is the total dielectronic recombination rate into j, \( \beta_{j+1} \) is the total three-body recombination rate into j.

The three-body recombination process is just the inverse of collisional ionization and is related by detailed balance as

\[
N_j^* S_j = N_{j+1}^* N_e \beta_{j+1}
\]

where the * indicates LTE population. It can easily be seen from the last two equations that in the limit where

\[
N_e \beta_{j+1} \gg \alpha_{j+1} + D_{j+1}
\]

that the population ratio becomes equivalent to that given by the Saha-Boltzmann equation, the LTE result. Also, in the opposite limit where

\[
N_e \beta_{j+1} \ll \alpha_{j+1} + D_{j+1}
\]

the conditions of coronal equilibrium are attained, where the ratio is independent of electron density, but still dependent on temperature.

A useful plasma model for spectroscopic analysis must contain a number of levels within each ionization stage. In practice, levels with principal quantum \( n < 6 \) usually need only be considered since the rest may be grouped with the next ionization stage and assumed to be in LTE with the ground state using the criterion of equation (2-3).

An effective cutoff of principal quantum number is obtained by the lowering of the ionization potential of an ion in a plasma, which leaves
as unbound those states with \( n > n_{\text{max}} \). This effect is actually a reduction of binding energy of an ionic electron due to the presence of nearby ions and electrons in the plasma. It has been calculated in several limiting cases using Debye-Huckel and Thomas-Fermi theory\(^{13}\), though the results are uncertain in application to laser produced plasmas in current and future experiments.

The rate equation for the \( i \)th level will have the form,

\[
\frac{dN_i}{dt} = -N_i \text{ (depopulation rates)} + \sum_j N_j \text{ (population rate from } j) \tag{2-6}
\]

The system of equations for all the levels is solved in the steady-state, i.e. the time derivative is set equal to zero. Given the electron temperature and density, the relevant rates in equation (2-6) can be determined and the system of equations solved for the populations.

**Atomic Processes**

The conditions generated in many laser-driven implosions are such that many important hydrogenic and helium-like spectral features of the gas fill are easily observed. Thus appropriate atomic physics models generally include explicitly the bare, hydrogenic, helium-like and lithium-like ionization stages, and group the remainder in one composite level, whose population is usually negligible. This fact, and the criterion of equation (2-3), is useful in keeping the total number of
levels in the model small and the solution to the equation (2-6) tractable.

The following processes must be included in the collisional-radiative model (for ionization species, $Z$, and ionic level $i$):

a) electron impact excitation (de-excitation)

$$e^- + N_{z,i} + e^- + N_{z,j} \quad i<j$$

b) electron impact ionization (3-body recombination)

$$e^- + N_{z,i} + N_{z+1} + e^-$$

c) photoionization (radiative recombination)

$$N_{z} + h\nu \rightarrow N_{z+1} + e^-$$

d) spontaneous emission

$$N_{z,j} + N_{z,i} + h\nu \quad \text{i<j}$$

e) Dielectronic recombination (autoionization)

$$N_{z+1} + e^- + N^{**}_z \quad \text{doubly excited}$$

$$N^{**}_z + N_{z,i} + h\nu$$
Appendix A gives additional details on the specific rates which have been used in the calculations reported in this dissertation. Hydrogenic approximations are used for photoionization rates, new calculations of collisional excitation and ionization rates are available in the literature;\textsuperscript{14,15,16} total dielectronic recombination rates were calculated using the method discussed in the next section with autoionization rates available in the literature.\textsuperscript{1}

An additional process which must be included in an optically thick plasma is photoexcitation; this is generally most important for resonance lines. The photoexcitation rate between levels with \( n=1 \) and \( 2 \) is given by

\[
R = \frac{\hbar \nu}{4\pi} B_{12} N_1 \int I_{(v,\Omega)}(v) L(v) \, dv \, d\Omega
\]

where \( N_1 \) is the ground state density and \( B_{12} \) is the Einstein coefficient and \( L(v) \) is the intrinsic lineshape.

The fact that the photoexcitation rate depends on the radiation field greatly complicates the analysis since the radiation field depends on level populations through the transfer equation, which in the two-level atom approximation is,

\[
\overrightarrow{\Omega} \cdot \overrightarrow{\mathbf{v}} = \frac{\hbar \nu}{4\pi} \left[ N_2 A_{21} - (N_1 B_{12} - N_2 B_{21}) \right] I L(v)
\]

In principle, it is necessary to couple the transfer equation to the level rate equations to obtain a consistent solution for populations and the final emergent spectral radiation. In practice it is possible instead to modify the rate equations using the escape factor
method. In this approach, terms in the rate equations proportional to spontaneous radiative decay in the optically thick transition are multiplied by an escape factor, \( g \), where \( g \) ranges from 0 to 1. The escape factor is defined to approximately include the effects of photoexcitation in the rate equations. A rate equation for level \( i \) can be written

\[
\frac{dN_i}{dt} = -N_i A_{ij} + N_j B_{ji} \frac{hv}{4\pi} \int I(\nu, \Omega) L(\nu) \, d\nu \, d\Omega + \text{other terms}
\]  

since the Einstein coefficients are related we can group two of the terms in equation (2-9) and obtain

\[
\frac{dN_i}{dt} = -N_i A_{ij} g(\tau_0) + \text{other terms}
\]  

where \( g(\tau_0) \), the escape factor, is given by

\[
g(\tau_0) = 1 - \frac{N_i}{N_j} \frac{\omega_i}{\omega_j} \frac{c^2}{2h\nu_0^3} \int I(\nu) L(\nu) \, d\nu
\]  

With the assumption of a uniform plasma source and a lineshape determined by the Doppler effect, Holstein obtained the result

\[
g(\tau_0) = \frac{1}{\tau_0 / \pi \ln \tau_0}
\]  

where \( \tau_0 \) is the optical depth at line center. In this approximation, the escape factor must also be included in the calculation of the emitted intensity (equation (2-11)).
Model Calculations

The methods of the previous sections have been applied to calculate ionic populations for steady-state argon plasmas under conditions expected in the experiments considered in this dissertation. Figure 1 gives some results for total hydrogenic and helium-like populations as a function of temperature under conditions of LTE, coronal, and collisional radiative equilibrium with the electron density fixed at $1 \times 10^{23} \text{ cm}^{-3}$. Coronal equilibrium rates were obtained from Shull and von Steenberg\textsuperscript{19} and collisional-radiative calculations were done at Lawrence Livermore National Laboratory using the code RATION.\textsuperscript{20} The code RATION was modified to include opacity effects of the resonance lines approximately in the rate equations with an escape factor. These results are included in Figure 1. It can be noted that the inclusion of opacity effects in the collisional-radiative model tends to drive the populations away from the coronal result and towards LTE.

More complete, time dependent models are currently being developed\textsuperscript{5,6,11} which couple hydrodynamic results with the rate equations and transfer equation. Specifically, the hydro code produces total ionic density and temperature as a function of time and space for a particular implosion. This information can then be used as input to the rate equation and transfer equation which can then be iteratively solved to yield the final emitted, time integrated spectral intensities. Currently these approaches have not directly coupled the radiative transport with the hydrodynamics but use LTE opacities during the hydrodynamics calculations and assume that implosion dynamics are
FIGURE 1
Ionization Fraction Calculation

(a) Ionization fraction is shown as a function of electron temperature for a highly ionized argon plasma calculated in LTE for an electron density of $10^{23}$ cm$^{-3}$. The curves show the fraction of argon ions which are in the bare through Li-Like ionization stages.
(b) Ionization fraction is shown as a function of electron temperature for a highly ionized argon plasma calculated in the corona approximation. The curves are labeled as in the previous figure.
(c) Ionization fraction is shown as a function of electron temperature for a highly ionized argon plasma calculated using a collisional-radiative steady-state model with fixed electron density of $10^{23}$ cm$^{-3}$. The curves are labeled as before, and opacity effects are included with an escape factor approximation.
not sensitive to non-LTE effects. While this assumption may be adequate for D-T plasmas with small argon impurities, it is apparently not true in the present case where the targets are argon/neon filled microballoons. The results using the hydrodynamic code LILAC$^5$ and a post-processing atomic rate equation package were inadequate to model these experiments. In general, the LILAC code did not treat the radiation effects in these optically thick plasmas correctly and produced results grossly inconsistent with the observations.

**Line Ratios**

Line intensity ratios can in general be dependent on electron density and temperature. Suitably chosen line ratios which are primarily dependent on either density or temperature can also provide useful diagnostics.

To form a temperature sensitive line ratio, we consider two transitions to the same ground state; for example Ly-$\beta$ and Ly-$\gamma$ lines:

$$\frac{I_{31}}{I_{41}} = \frac{N_3 A_{31} \nu v_{31}}{N_4 A_{41} \nu v_{41}} \quad 2-13$$

If the coronal approximation is valid, collisional excitation from the ground state is balanced by radiative decay for each transition and we can write

$$\frac{I_{31}}{I_{41}} = \frac{C_{13} \nu v_{31}}{C_{14} \nu v_{41}} \quad 2-14$$
From the form of the collision rate, $C$, given in Appendix A it is seen that the above ratio is dependent on temperature and independent of electron density. This is also the case in LTE [equation (2-1)].

To obtain a density dependence, we can consider two lines, one of which has a long lifetime (low $A$ value) such that the process of collisional de-excitation can become important. For two optically thin lines, the rate equations for the level populations are

\[
\frac{N_i A_{ij}}{N_i A_{kg} + N_k N_C e^{-k/T e}} = \frac{N_i N C e^{-k/T e}}{N_i N C e^{-k/T e}}
\]

where the $C$'s are collisional excitation and de-excitation-rates. It is evident that the intensity ratio $I_{4g}/I_{4g}$ is a function of electron density, and possibly temperature. An example of this situation occurs in the helium-like argon resonance line of this experiment. The resonance line transition is

\[
\text{1s2p } ^1P + \text{1s}^2 \text{1S}
\]

and can be compared with the intercombination line

\[
\text{1s2p } ^3P + \text{1s}^2 \text{1S}
\]

where the intercombination (ic) line has a low transition probability (dipole forbidden) and fulfills the conditions of equation (2-15). As the resonance line is affected by photoexcitation, it is also convenient
to consider the more nearly optically thin line ratio $I(1s3p-1s^2)/I(ic)$ which is also density dependent, and slightly dependent on temperature. In addition, it is useful to compare lines from adjacent ionic species and use a collisional-radiative model in the analysis. For example, the ratio $I(\text{Ly}\beta)/I(\text{He}\beta)$ compares a prominent hydrogenic line with a permanent helium-like line, neither of which is greatly affected by opacity. Figures 2 and 3 give some results of line ratio calculations using a collisional-radiative model. These will be referred to in chapter IV when comparing to the experimental data.

**Satellite Line Models**

Satellite lines arising from doubly-excited autoionizing levels are prominent in high temperature plasmas. These lines are usually optically thin and are useful as additional diagnostics of temperature and density. Their presence also indicates the importance of an additional atomic process, dielectronic recombination, which must be included in plasma models. This process was first recognized in an astrophysical context\(^{21,22}\) and the importance for laser produced plasmas was soon recognized.\(^{23,24}\)

The process of satellite line formation can be explained by reference to Figure 4, which gives schematically some ionic levels for three different ionization stages. Dielectronic recombination is a two step process: first, a doubly excited atomic state is formed by direct electron capture from a ground state or excited state of the next ionization state, i.e. for Argon,
FIGURE 2

Line intensity ratio of (1s3p-1s^2) to Lyβ shown as a function of temperature, and calculated with a collisional-radiative model and a pure argon plasma.
FIGURE 3

The resonance to intercombination line ratio shown as a function of electron density and calculated with a collisional-radiative model and a pure argon plasma.
FIGURE 4

Energy level diagram over three ionization stages showing levels giving rise to satellite line emission (indicated by arrow). The dashed lines indicate series limits, with levels above these limits being autoionizing.
The doubly excited $2\ell 2\ell'$ state may autoionize and return to the initial state, or emit a photon and form a bound state,

$$N(\text{Ar}^{+16}_{2\ell 2\ell'}) + N(\text{Ar}^{+16}_{1s2\ell}) + \text{hv}$$

The net effect of these two steps is thus a recombination process leading to a change in ionic species from $\text{Ar}^{+17}$ to $\text{Ar}^{+16}$.

In a steady-state low density approximation the rate equation for the doubly excited level population $N_1$ is given by balancing electron capture with radiative decay and autoionization,

$$N_1 \left( \Gamma_1 + \sum_k A_{1k} \right) = D_1 N_e N_H$$

where $\Gamma_1$ is the autoionization rate, $D_1$ is the electron capture rate, $N_H$ is the initial ground state level population, and $A_{1k}$ the Einstein transition probability for any possible radiative decay. Autoionization rates have been calculated from atomic structure codes for a number of helium-like and lithium-like autoionizing levels (see Appendix). Electron capture rates can then be calculated from the principle of detailed balance. Using equation (2-1) and assuming capture from a hydrogenic ground state into a helium-like doubly excited level,

$$D_1 = \frac{\sigma_1/\sigma_H}{6.04 \times 10^{-21} \text{ ev}^{-3/2}} \Gamma_1 \exp(-E_1/kT)$$
where $E_i$ represents the energy of the doubly excited level above the ionization limit (see Figure 4). Using the previous two equations it can then be seen that the intensity of a satellite line from level $i$ is proportional to

$$I_{ij} \propto \frac{\Gamma_i N_e}{\Gamma_i + \sum_k A_{ik}} \exp\left(-\frac{E_i}{kT}\right) \quad 2-18$$

Due to the exponential in the above equation, $kT$ must be large compared to $E_i$ for satellite emission to be prominent. The above equation also predicts that the ratio of two satellite lines, in the coronal approximation, is

$$\frac{I_{ij}}{I_{mn}} = \frac{q_{ij}}{q_{mn}} \exp\left(-\frac{E_{ij}}{kT}\right) \quad 2-19$$

where

$$q_{ij} = \frac{g_i \Gamma_i A_{ij} \hbar \nu_{ij}}{\Gamma_i + \sum_k A_{ik}}$$

and $E_{ij} = E_i - E_m$. In this approximation the relative satellite intensities are independent of electron density.

The satellite line ratio in the high density, collision dominant limit can easily be written down, since the populations are in the ratio of the statistical weights of the levels

$$\frac{I_{ij}}{I_{mn}} = \frac{r_{ij}}{r_{mn}} \exp\left(-\frac{E_{ij}}{kT}\right) \quad 2-20$$

where
\[ r_{ij} = g_i A_{ij} h \nu_{ij} \]

It was first noticed by Seely\textsuperscript{24} that a particular group of helium-like satellite lines (from the $2p^2 \ 3p$ level) was more intense than the coronal model prediction. The coronal model is also not valid for lithium-like satellite lines of the helium-like resonance line.

A more general collisional-radiative model for the autoionizing level populations will yield satellite line intensity ratios as a function of density and yield the previous results in the low and high density limits. The model will include the following important collisional processes (see Figure 4):

1) electron induced collisional excitation and de-excitation among the $2\ell 2\ell'$ (or $1s2\ell 2\ell'$) doubly excited levels.
2) electron induced collisional ionization from the $2\ell 2\ell'$ levels to $2\ell$ (or $1s2\ell$) state of the next ionization stage
3) electron induced inner shell excitation from bound helium-like (lithium-like) levels.

A method similar to that of Jacobs and Davis\textsuperscript{25} was used to set up and solve a system of rate equations for the helium-like and lithium-like level populations giving rise to the prominent argon satellite lines seen in this experiment. All the previously mentioned collisional and radiative processes were included in the model in order to predict overall satellite line emission and to compare with the experiment.

The doubly excited levels included in the analysis for both helium-like and lithium-like satellite lines are shown in Figure 5. In either case, six levels were included in the system of rate equations, which
(a) The six helium-like autoionizing levels included in the satellite line calculations are shown on a relative energy level diagram.
FIGURE 5 (continued)

(b) The six lithium-like autoionizing levels included in the satellite line calculations are shown on a relative energy level diagram.
were then solved for doubly excited level populations in terms of the hydrogenic (helium-like) ground state population. It was assumed that the population within each group could be distributed among the sublevels of angular momentum j according to the statistical weights, \((2j+1)\). Atomic data for the particular argon doubly excited levels and satellite lines were obtained from the literature, and the specific forms for the collision rates are given in Appendix A.

The steady state rate equation for the \(i^{th}\) doubly-excited level can be written in a more detailed form than that given in equation (2-6):

\[
\sum_{i', \neq i} Q_{i,i'} N_{i'} = D_{i,g} N_e N_e + T_{B_{i,b}} N_e N_e^2 + \sum_{j} C_{ex(i,j)} N_j N_e
\]

where \(D(i,g)\) is the electron capture rate from the ground state of the higher ionization stage, \(T_{B_{i,b}}\) is the three body recombination rate from the \(n=2\) bound state of the higher ionization stage, \(C_{ex(i,j)}\) is the inner shell collisional excitation rate from lower bound levels. The rate matrix \(Q_{i,i'}\) has diagonal elements which represent the total depopulation rate of level \(i\) due to autoionization, radiative decay, and collisionally excited transitions, i.e.

\[
Q_{i,i} = \sum_{i', \neq i} N_e C_{(i+i')} + \Gamma_{i} + \sum_{f} A_{if} + N_e C_{i}(i+b)
\]

where \(C_{i}\) is the inverse process of \(T_{B_{i,b}}\), and \(\Gamma_{i}\) is the inverse of \(D\). It is the collisionally induced mixing transitions among the doubly excited levels which is particularly responsible for driving the populations away from coronal equilibrium and toward LTE values.
The off-diagonal terms of $Q(i,i')$ represent collisional excitation or de-excitation transitions from other doubly excited levels, $i'$, into the level $i$,

$$Q(i,i') = -N_e C(i',i)$$  \hspace{1cm} (2-23)

The solution for level population $N_i$ of the system of rate equations defined in equation (2-21) is given by

$$N_i = \sum_{i,j} Q^{-1}(i,i') [D(i',g) N_g N_e + TB(i',b) N_b N_e^2]$$

$$+ \sum_{i,j} Q^{-1}(i,i') C_{ex(i',j)} N_j N_e$$  \hspace{1cm} (2-24)

which depends on the ground state, $N_g$, and first excited state, $N_b$, of one ionization stage and the excited bound states, $N_j$, of the next lower ionization stage. The relative populations, $N_b/N_g$ and $N_j/N_g$, were calculated using less detailed models of the type previously described and also using several experimentally obtained line ratios. Thus the level populations can be written in terms of $N_g$, ground state populations. The relative intensities of several satellite lines will then be independent of $N_g$.

It should be noted that this solution is guaranteed to reproduce both coronal and LTE populations in the appropriate density limits. Also the total dielectronic recombination rate may be defined as

$$\alpha_D = \sum_{ij} A_{ij} N_i$$

$$\frac{N_i}{N_g}$$  \hspace{1cm} (2-25)
where $N_1$ is given in equation (2-24). We can then write

$$q_D = \sum_{i,i',j} A_{ij} Q^{-1}_{(i,i')} D(i',g)$$  \hspace{1cm} (2-26)

This process is an important recombination mechanism for the total ionization balance in a high temperature plasma.

**Satellite Line Calculations**

The methods of the previous section were used to calculate the total relative satellite line intensities for argon plasmas over a density range covering the transition from corona to LTE conditions. Both the helium-like $2\pi2\pi'$ satellites of the Lyman alpha resonance line and the lithium-like $1s2\pi2\pi'$ satellites of the helium-like resonance line were included in the calculations. Equation (2-1) is used to calculate the relative line intensities using the level populations from equation (2-24).

The overall satellite emission is composed of a blend of many transitions from the group of six doubly excited levels. Using the calculated populations and the atomic data for each transition the total relative satellite emission was obtained by summing the intensities over all frequencies of satellite line emission for all lines and normalizing to the intensity of the strongest satellite line. Each individual satellite component was assumed to have, for simplicity, a Lorentzian shape whose width was determined primarily from source broadening, which was shown in Chapter I to be approximately 3.5eV. The Lorentzian shape has the form
\[ L_\nu \propto \frac{\Gamma}{(\nu - \nu_0)^2 + (\Gamma/4\pi)^2} \]

where \( \Gamma \) is the full width at half the maximum intensity. Voigt profiles were also assumed for satellite line shapes but these in general did not prove to fit the data as adequately as did the Lorentzian shapes.

Figures 6 and 7 present the overall blend of satellite emission for the helium-like and lithium-like satellites in argon plasmas and show both low density (corona) to high density (LTE) emission. It can readily be seen that the emission illustrated in these two figures changes dramatically; this fact will be used as a density diagnostic for experiments analyzed in this dissertation. Each of these figures is labeled giving the component transitions of the blend in spectroscopic notation. The usual spectroscopic notation in LS coupling is used throughout, i.e. \((2S+1)L_J\).

The increased emission with increasing density, shown in the helium-like satellites of Figure 6, comes from the \(2p^2 \, 3F\) levels which have low autoionization rates (see Appendix for selection rules). These levels become populated by collisional transitions from other doubly excited levels, a process which becomes dominant as density increases. The calculations show that this effect will be noticeable precisely when the condition expected in conditions of this experiment are realized.

The satellites lines of Figure 6 can just be resolved into four major component blends. The resolvable ratio of these turns out to be a good measure of the density effect just described. Thus the methods described in the previous section were applied to calculate the ratio, \( R \).
Figure 8 shows the smooth transition of this ratio from the coronal to LTE limits. Also shown are the effects of inner shell excitation which are predominant at low density and the effects of collisional ionization, important at very high density.

The density effect apparent in Figure 7 for the lithium-like satellites can also be measured by a resolvable line ratio. Using the shorthand notation introduced by Gabriel, the density-dependent line ratio is

$$R = \frac{I(abcd) + I(qr)}{I(jkl)}$$

where a,b,c,d refer to transitions of the type \(1s2p^2 \, ^2P - 1s2s \, ^2P\); q,r refer to transitions of the type \(1s2s(1P)2p \, ^2P - 1s2s \, ^2S\); and j,k,l refer to transitions of the type \(1s2p^2 \, ^2D - 1s2p \, ^2P\). This line intensity ratio was calculated with the satellite line model described in the previous section and the results are shown in Figure 9.

The results given in Figures 8 and 9 will be used in Chapter IV where the respective line ratios are measured from the data and compared to the calculations for a density diagnostic.

Higher Order Satellites

The analysis of the previous section considered only n=2 doubly-excited levels. Higher order satellites have previously been observed
FIGURE 6

The blend of satellites of the Ar+17 Lyα line is shown as a function of electron density. The numbered components represent the transitions:

1. $2p^2 \, 1D - 1s2p \, 1P$
2. $2p^2 \, 3P - 1s2p \, 3P$
3. $2s2p \, 3P - 1s2s \, 3S$
4. $2s2p \, 1P - 1s2s \, 1S$
5. $2p^2 \, 1S - 1s2p \, 1P$
FIGURE 7

The blend of satellites of the Ar$^{+16}$ resonance line (1s2p–1s$^2$) is shown as a function of electron density. The numbered components represent (predominantly) the transitions:

(1) j,l  \hspace{1cm} 1s2p$^2$ 2$^2$D $\rightarrow$ 1s$^2$2p 2$^2$P
(2) k  \hspace{1cm} 1s2p$^2$ 2$^2$D $\rightarrow$ 1s$^2$2p 2$^2$P
(3) a,b,c,d  \hspace{1cm} 1s2p$^2$ 2$^2$P $\rightarrow$ 1s$^2$2p 2$^2$P
(4) q,r  \hspace{1cm} 1s2s (1$^1$P) 2p 2$^2$P $\rightarrow$ 1s$^2$ 2s 2$^2$S
(5) m,n,s,t  \hspace{1cm} 1s2s (3$^3$P) 2p 2$^2$P $\rightarrow$ 1s$^2$ 2s 2$^2$S \\
\hspace{3cm} 1s2p$^2$ 2$^2$S $\rightarrow$ 1s$^2$ 2p 2$^2$P
(6) higher order 1s2\ell3\ell' satellites

The letters are the shorthand notation introduced in reference
FIGURE 8

The helium-like satellite line intensity ratio of equation 2-28 is shown as a function of electron density. Effects of inner shell excitation and resonance line opacity are indicated.
FIGURE 9

The lithium-like satellite line intensity ratio of equation (2-29) is shown as a function of electron density. Effects of resonance line opacity are indicated.
which arise from \( n=3 \) and higher levels. In Figure 4 some of the \( n=3 \) levels are shown, namely \( 2\ell 3\ell' \) and \( 1s2\ell 3\ell' \) in the helium-like and lithium-like stages, respectively. These levels have several channels for radiative decay. For example, the \( 2\ell 3\ell' \) level may follow

\[
2\ell 3\ell' + 1s3\ell' + h\nu
\]

or

\[
2\ell 3\ell' + 1s2\ell + h\nu
\]

to form a satellite of Lyman \( \alpha \) or Lyman \( \beta \), respectively. Generally, the higher order satellites of the resonance line will not be resolved and hence will tend to broaden the resonance line.

The populations of these higher order doubly excited levels can usually be taken to be close to the LTE limit, as seen from the criterion of equation (2-3). Then the ratio of two satellite line intensities arising from these levels is given by equation (2-20), the high density limit. As will be shown in Chapter IV, these high order satellites are seen in the experimental spectra.

**Temperature Diagnostics**

It has already been shown in equation (2-18) that the presence of satellite line emission already implies a high temperature. The intensity ratio of a satellite to resonance line provides a useful additional temperature diagnostic. Consider an optically thin resonance
line whose upper state population is determined by collisional
excitation from the ground state. For a satellite line in the coronal
approximation, it can be shown that

\[
\frac{I_{\text{sat}}}{I_{\text{Res}}} = \frac{A_s}{C_{s2p}} \frac{D_s}{\Gamma_s + \Sigma A_{is}}
\]

where \( C_{s2p} \) is the collisional excitation rate for the resonance line
upper state; \( D \) and \( \Gamma \) are the electron capture rate and autoionization
rate of the satellite level. It is evident that this line ratio is a
function of temperature and atomic constants, and is not dependent on
electron density. This ratio was calculated for the strong satellite
transition,

\[
2p^2 \, 1D + 1s2p \, 1P
\]

and compared to the Ly-\( \alpha \) resonance line and is shown in Figure 10.
Although the resonance line is not optically thin, this ratio is still a
useful temperature diagnostic (see Chapter IV).
Intensity ratio versus temperature for the Ar$^{+16}$ $2p^2 1D - 1s2p 1P$ satellite to the Ly-$\alpha$ resonance line of Ar$^{+17}$ the model described by equation (2-30) was used.
$^{1}D$ Satellite to Ly $\alpha$ Resonance

$R$

$T$ (eV)

500 700 900 1100 1300 1500
CHAPTER III
SPECTRAL LINESHAPES IN DENSE PLASMAS

Introduction

Theoretical shapes of spectral lines emitted from dense plasmas are determined by the convolution of several processes. These include natural broadening of the line due to the finite lifetime of the atomic states, Stark broadening due to the collisional processes between the radiating ion and electrons and other ions, and Doppler broadening due to thermal motion. Stark broadening theories yield lineshapes for transitions in ions whose energy levels are perturbed by the electric fields due to dense plasma constituents. Under conditions generated in current laser-produced implosion, Stark widths tend to dominate the other causes of line broadening. The fact that Stark lineshapes are especially sensitive to electron density and relatively insensitive to temperature means that observed spectral lineshapes are good measures of achieved densities in these experiments.

While observed lineshapes are primarily due to the Stark effect, several other causes of line broadening are relevant in these plasmas. Doppler broadening at these high temperatures must be convolved with the Stark shape, though it is usually a small effect. More significant is the additional opacity broadening of lines which occur when radiation passes through the plasma. If more absorbing ions are present, or if the plasma covers a larger region of space, there is a greater chance of
photon absorption and subsequent re-emission within the plasma with a slight change in frequency. The net effect on the escaping line profile is described through the equation of transfer, the result is generally termed "opacity" broadening. This effect is most important for the resonance line transitions which have the largest photon reabsorption probabilities. Finally, all observations are affected by instrument and source broadening, as described in Chapter I. This is generally a small effect, basically smoothing the profile near line center.

In order to accurately compare theoretical lineshapes with observations it is necessary to include all the above processes in the theoretical calculations. The model theoretical lineshapes were produced by including the broadening processes in the following order: Stark + Doppler + Opacity + Instrument. This chapter discusses how theoretical lineshapes are generated and applied in the analysis of the data presented in the next chapter.

In addition, a systematic method is presented to analyze several spectral lines in the same series while self-consistently including opacity effects. This method attempts to formalize a procedure which have been presented, in part, elsewhere.\textsuperscript{26} Previous theoretical developments have generally assumed that profiles of lines originating from high \( n \)-states are optically thin and yield a better estimate of final core density than do those originating from lower \( n \)-states. However, optically thick resonance line is analyzed to obtain the ground state ion density. Here, we demonstrate how all the observed lines in a given series can be used to obtain the "best" consistent electron and ion densities.
Stark Broadening

A number of Stark and Doppler broadened lineshapes have been calculated for the principal series transitions of Ar$^{+17}$ and Ar$^{+16}$ in plasmas consisting of various argon/neon mixtures. The principal series transitions considered are

\[
\begin{align*}
\text{Ar}^{+17} & \quad 1s - n\ell \quad n < 6 \\
\text{Ar}^{+16} & \quad 1s^2 - 1s2\ell \quad n < 5
\end{align*}
\]

The calculations covered the electron density range

\[
5 \times 10^{22} < N_e < 1 \times 10^{24} \text{ cm}^{-3}
\]

and the temperature range

\[
800\text{eV} < T_e < 1200\text{eV}
\]

The lineshape formalism of Tighe and Hooper\textsuperscript{27} was used with extension to helium-like lineshapes using the method of Joyce, Woltz and Hooper.\textsuperscript{28} Fine structure splitting of the upper state level, an important source of asymmetry in the lineshape of the first few series members, was included in these calculations; it was found to be especially important for Lyman $\alpha$ and $\beta$ lines.\textsuperscript{29}

Lineshapes were calculated using the static ion approximation. That is, it was assumed that except for times long compared to an electron collision time the ions could be treated as creating a static electric field at the radiating ion. In this approximation the
intensity profile is given by

$$I(\omega) = \int_{-\infty}^{\infty} P(\varepsilon) J(\omega, \varepsilon) \, d\varepsilon,$$

where $P(\varepsilon)$ is the ion microfield probability distribution\(^{27}\) and $J(\omega, \varepsilon)$ is the electron broadened line profile for a radiating ion in an ion microfield $\varepsilon$. The total profile is then given by an average over all possible ion microfields. The dipole approximation is used for the radiator-perturber interaction and the electron broadened line profile is calculated with a second order perturbation theory in the no-quenching approximation.\(^{30}\) Although most of these approximations can be removed, they should be valid over the range of densities considered here.\(^{29,30}\)

The Doppler effect is included through a convolution with the Stark profile

$$I(\omega) = \int_{-\infty}^{\infty} I(\omega') \phi(\omega-\omega') \, d\omega' ;$$

and the Doppler profile is given by the gaussian form,

$$\phi = \frac{1}{\beta \sqrt{\pi}} \exp \left[ -\frac{(\omega-\omega_0)^2}{\beta^2} \right]$$

The parameter $\beta$ is the Doppler half-width given by

$$\beta = \frac{(2kT_{\text{eV}})^{1/2}}{mc} \omega_0$$

$$\approx 10^{-3} T_{\text{eV}}^{-1/2} Z^{3/2} \text{ Ryd}$$
In this treatment, it is assumed that electron collisions have no effect on the radiating ion's motion during a radiative lifetime. This is generally the case since the electron mass is much less than the radiator's mass.31

Under the conditions produced in the current experiments, the Doppler effect generally produces a small change about the center of the overall Stark dominated profile. The conditions under which Doppler broadening would predominate can be estimated using equation (3-4) and the Stark broadening results. The condition for the Doppler effect to dominate assuming temperatures typical for the implosions analyzed here and assuming that we have a hydrogenic ion of charge Z is

\[ N_e < 5 \times 10^{-3} \left( \frac{\alpha}{a_0} \right)^3 Z^{21/4} n^{-3} \], \hspace{1cm} 3-5

where \( \alpha \) is the fine structure constant and \( a_0 \) the Bohr radius (cm) for transitions from principal quantum number, \( n \). For the Ly-\( \alpha \) line of \( \text{Ar}^{+17} \) and \( T \sim 1 \text{ keV} \) this condition gives

\[ N_e < 7 \times 10^{21} \text{ cm}^{-3} \]

which is not satisfied in these experiments.

Radiation Transport

The passage of radiation through an absorbing medium may affect the lineshapes finally observed, depending upon whether the material is optically thick or thin to the radiation field \( I(\nu) \). In the optically
thin case, where the mean free path for photons to be absorbed is larger than the size of the absorbing medium, the radiation field is relatively unaffected by the medium. In the optically thick case, the observed radiation has been subject to a number of absorption and re-emission processes in the medium. The equation of transfer is used to describe these processes; it may be written as,

\[
\frac{1}{c} \frac{\partial I}{\partial t} + n \cdot \vec{V}I = \eta - \chi I,
\]

where \(\eta\) is the total emissivity for the medium (energy emitted per unit volume per unit time per unit frequency per unit solid angle), \(\chi\) is the opacity or absorption coefficient (units of cm\(^{-1}\)), and the photon mean free path is \(1/\chi\). For the cases of planar and spherical geometry the directional derivative may be written

\[
-\eta \cdot \vec{V}I = \mu \frac{\partial I}{\partial x} \quad \text{planar} \quad (3-7)
\]

\[
= \mu \frac{\partial I}{\partial x} + \frac{1}{r}(1-\mu^2) \frac{\partial I}{\partial \mu} \quad \text{spherical}
\]

here, \(\mu=\cos \theta\) and \(\theta\) is the angle between the ray and local normal. The time dependent term in equation (3-6) can be neglected in this analysis. This is true because the overall plasma conditions do not change significantly during a photon travel time through the medium (< 1 picosecond) in these experiments.

In the planar case, the transfer equation may be written in the form

\[
\mu \frac{\partial I}{\partial t} = I - S
\]

\(3-8\)
where the following definitions are used

\[ S = \frac{n}{\chi} \]
\[ dt = -\chi \, dz \]

and \( S \) is called the source function, \( \tau \) is called the optical depth.

The transfer equation in the form of equation (3-8) can be formally solved by using an integrating factor, \( \exp(-\tau/\mu) \). The normally emergent intensity from a planar medium is given by

\[ I = \int_{0}^{T} S(t) \, e^{-t} \, dt \]

where the variable \( t \) represents the variable of optical depth and \( T \) is the total optical depth of the medium. This result says that the normally emergent intensity is given by the Laplace transform of the source function.

While the source function is not in general known as a function of optical depth, the above equation may be solved readily for a constant source function, \( S \), yielding

\[ I_{\nu} = S[(1-\exp(-T_{\nu}))] \]

In general the source function is not constant but depends implicitly on the unknown radiation field. However in LTE, the source function is simply given by the Planck function \( B_{\nu} \),

\[ B_{\nu} = \frac{2\nu^3}{c^2} \left(e^{\frac{\nu}{kT}} - 1\right)^{-1} \]
It is useful to rewrite the transfer equation in microscopic form for a resonance line. If we consider only a two-level atom, the atomic processes to be considered are spontaneous emission, stimulated emission, and photoexcitation. Then we obtain

\[ \mu \frac{3I}{\partial z} + \left[ n_2 A_{21} - (n_1 B_{12} - n_2 B_{21})I \right] L(\nu) \frac{hv}{4\pi} \]

which has the form of equation (3-6) where the right hand side represents the difference of emission and absorption processes for the resonance line transition of intrinsic lineshape \( L(\nu) \). The absorption coefficient can be rewritten using the definition of the Einstein \( B_{ij} \) coefficient. We obtain,

\[ \chi = (n_1 B_{12} - n_2 B_{21}) L(\nu) \frac{hv}{4\pi} \]

\[ = \frac{ne^2}{mc} f_{12} n_1 \left( 1 - \frac{g_1 n_2}{g_2 n_1} \right) L(\nu) \]

For a uniform slab of thickness, \( z \), the total opacity is obtained using equation (3-9)

\[ \tau_\nu = \chi_\nu z \]

\[ = \frac{ne^2}{mc} f_{12} n_1 z L(\nu) = 1.64 \times 10^6 f_{12} n_1 z L(\nu) \]

if the slab thickness is given in \((\text{cm})\), the intrinsic line profile has units \((\text{eV})^{-1}\) and the density of absorbing ions has units \((\text{g/cm}^3)\). The constant is calculated with the argon ion mass. This form for the
optical depth assumes that the stimulated emission term of equation (3-14) could be neglected. This is the case since $g_1 n_2 < g_2 n_1$ unless other pumping mechanisms become important.

The expression for the intensity emitted from a uniform slab [equation (3-11)] can now be written in terms of optical depth at line center, $\tau_0$, as

$$I_\nu = S_\nu [1 - \exp \left( -\frac{\tau_0 L(\nu)}{L(\nu)} \right)]$$

which is true since the ratio $\tau_\nu / L(\nu) = \tau_0 / L(\nu)$ for frequencies within a given lineshape, $L(\nu)$. In this equation $L(\nu)$ refers to profile value at line center, and $L(\nu)$ is the profile value at frequency $\nu$. Note that in the optically thin limit, since the exponential argument is small, we obtain,

$$I(\nu) \propto L(\nu)$$

and the emitted intensity is simply a reflection of the intrinsic lineshape, $L(\nu)$. It will be seen that the frequency variation of the source function is not significant over the lineshape. In the optically thick case, when $\tau >> 1$, we obtain

$$I(\nu) \propto S(\nu)$$

which indicates a constant emission for frequencies in the line which are optically thick. This gives rise to an overall profile which is flat topped for optically thick frequencies near line center. Further
out in the wing of the line the decrease in the profile intensity leads to an optically thin condition. The total lineshape obtained from equation (3-16) will have a broader profile than the optically thin intrinsic case and possibly with a flat top at line center. The effect of the opacity model of equation (3-16) on a theoretical Ly-α profile is shown in Figure 11 using several values of \( \tau_0 \). The theoretical profile includes the fine structure splitting as is apparent from the double peaked profile in the result for lowest optical depth.

This model of opacity broadening was used to obtain the ratio \( \tau_0/L(0) \) from a series of fits to the experimental data for the hydrogenic and helium-like argon lines. The validity of this model for opacity broadening will in fact be tested in this analysis since the extent to which it can be successfully applied depends in part upon the achievement of nearly uniform conditions during the time of maximum emission of the spectral lines. Severe temperature and density gradients occurring during the implosion will therefore be expected to affect the results obtained in the next chapter.

The solution to the transfer equation for the case of a uniformly emitting and absorbing sphere differs from the planar solution given in equation (3-11) by the inclusion of a \( \cos \theta \) factor in the exponential, where \( \theta \) is the angle between the outgoing ray and the sphere radius. The length in the definition of optical depth [equation (3-15)] becomes the sphere diameter and the total emitted flux is obtained by integrating the intensity over the sphere,

\[
I = \int_0^{\pi/2} S_\nu \left[ 1 - e^{-\tau \nu \cos \theta} \right] 2\pi \cos \theta \sin \theta \, d\theta
\]  

3-17
Opacity broadening of Ar$^{+17}$ Ly-α resonance line is shown for several values of optical depth, $\tau$, using the opacity model of equation (3-16). The theoretical Ly-α profile is for an electron density of $2 \times 10^{23}$ cm$^{-3}$, temperature of 100 eV and includes the effects of fine structure.
It can be shown by the evaluation of the flux\(^{26}\) that the intensity emitted from a slab of length \(z\) yields equivalent overall profiles to the flux emitted from a sphere of diameter \(1.5z\). Thus the slab size used in equation (3-15) was chosen to equal the measured core diameters multiplied by the \(1.5 \times \) correction factor to approximate for total emission from the spherical core.

**Non-LTE Opacity Model For Two-Level Atoms**

The model given in the previous section is inherently non-LTE, since the source function is not required to be the LTE Planck function and level populations are not specified. However, since the source function depends on the unknown level populations, it is possible to explicitly include the equations of statistical equilibrium for the level populations into the definition of the source function to obtain a solution for the emitted intensity which is also consistent with statistical equilibrium equations. In the two level atom approximations this solution can be obtained by following a straightforward numerical procedure. While the method is inherently non-LTE it includes the additional constraint of the statistical equilibrium equation explicitly in the source function; simultaneously solving the transfer and rate equations.

The source function defined in equation (3-9) can be written in the form for the two-level atom,
\[ S = \frac{n_2 A_{21}}{n_1 A_{12} + n_2 B_{21}} \]
\[ = \frac{2\hbar \nu^3}{c^2} \left[ \frac{n_1 g_2}{n_2 g_1} - 1 \right]^{-1} \]

where the final form uses the Einstein relations. In this approximation, population mechanisms from higher lying levels and the continuum is neglected. The ratio of level populations \( n_1/n_2 \) is obtained from the statistical equilibrium equations for the two levels,

\[ n_2 A_{21} + n_2 B_{21} \int L_\nu J_\nu \, d\nu + n_2 C_{21} \]
\[ = n_1 B_{12} \int L_\nu J_\nu \, d\nu + n_1 C_{12} \]

where \( L_\nu \) is the Stark broadened lineshape and \( J_\nu \) is the mean intensity of the radiation field,

\[ J_\nu = \int I_\nu \, d\mu \]

Combining equations (3-18) and (3-19) we can obtain an expression for the source function,

\[ S = (1-\varepsilon) \int L_\nu J_\nu \, d\nu + \varepsilon B \]

where the coefficient \( \varepsilon \) is given by

\[ \varepsilon = C_{21} \left[ C_{21} + A_{21} \left( \frac{e^{h\nu_0/kT}}{e^{h\nu_0/kT} - 1} \right) \right]^{-1} \]
and ε represents the probability per scattering that a photon will be lost from the line by collisional de-excitation and thus thermalized. Note that in the LTE limit where \( C_{21} \gg A_{21} \), the source function becomes the Planck function as ε becomes equal to one. The scattering term in equation (3-20) represents photon frequency diffusion, i.e. photons which are absorbed at frequency \( v \) can be reemitted at frequency \( v' \) and eventually escape after multiple scatterings.

The above form for the source function can be combined with the transfer equation (3-8); and the resulting equation can be solved by the method of discrete ordinates for both the intensity and source function as functions of optical depth. Then equation (3-10) can be used to find the emitted intensity in the spectral line. The procedure described by Avrett and Hummer\(^{35} \) has been used to describe the non-LTE radiative transfer of spectral lines using Stark broadened lineshapes. The results in this case are similar to that of the reference, which used Doppler and Voigt lineshapes.

The results were obtained for plane-parallel geometry with the approximation of a homogeneous, isothermal plasma and constant Planck function. The opacity broadened Lyman alpha lineshape is shown in Figure 12 using the two opacity broadening methods described. It is seen that this non-LTE method shows a self reversal near line center due to the non-coherent scattering term in the source function. This effect is similar though not equivalent to that of a cooler absorbing region. However, the overall shape yields a similar fit in the line wings to that calculated from equation (3-16) for the same optical depth \( \tau \). Since the wings of the lineshape are most important in obtaining
Opacity broadening of $\text{Ly}\alpha$ with electron density of $5 \times 10^3$ cm$^{-3}$ and optical depth at line center, $\tau_0$, equal to 50. The solid line shows the uniform slab results and the dashed line the non-LTE two level atom result.
lineshape fits to experiment; and since the theoretical is most uncertain near line center, it is concluded that the method of equation (3-16) can be used. It is further noted that self-reversals of the optically thick Ly-α lines are not seen in the experiment (except as can be explained by fine-structure corrections).

**Line of Best Fit Method**

A systematic approach was used to determine electron and absorbing ion densities self-consistently when analyzing all lines in a given spectral series. In particular, the principal series lines of Ar$^{+17}$ and Ar$^{+16}$ were analyzed by fitting theoretical lineshapes to the experimental data. Theoretical profiles used in the fitting included the effect of Stark, Doppler, opacity, and instrument broadening as previously described. The fitting procedure employed is similar to that introduced by Kilkenny and Lee$^{26}$ and can be called the "line of best fit" method. This method recognizes that for a given spectral line there are many values of ion and electron densities which can produce fits of equivalent "quality". Recall from equations (3-15) and (3-16) that a larger ion ground-state density yields a broader lineshape and that Stark broadening theory predicts broader profiles for plasmas of increasing electron densities. Then as a first step in our procedure a line is assumed to be optically thin and fit with a Stark profile which has no opacity correction. This optically thin profile provides an upper limit to the inferred electron density. Next, a series of additional fits for the same line is made, where the electron density is progressively lower and the opacity correction (ion density) is
progressively larger. This series of fits can then be plotted in the electron density-ion density plane to map out a line of best fit (see Figure 13). Similar lines are constructed for each observable spectral series member; for example Figure 13 shows Lyman α, β, γ, and δ lines of best fit for an idealized case. If it is conjectured that all of these lines are emitted under the same plasma conditions (so that each is characteristic of the same time history) then their point of intersection should unambiguously identify the values of electron density and the $\text{Ar}^{+17}$ ground state density which are consistent for all these lines. In practice, each line of best fit has an inherent uncertainty depending on the quality of the fit; then the lines define a region of consistency from which ion and electron densities can be estimated. Previous studies have observed a marked disagreement of the Lyman-α line of best fit with the other Lyman lines, and have relied upon higher series members for approximately optically thin electron density determinations. It is clear from the figure that such an approach will tend to produce an overestimate of density, except for the highest series members which are truly optically thin but probably not intense enough to be observable.

An additional constraint may be placed on the electron ion density plot with an ionization limit line. For example, all ionization equilibrium models indicate that the maximum density in the hydrogenic ionization state is approximately 50%, then

\[
\frac{n_H}{n_{\text{total}}} < \frac{1}{2}
\]

\[
< \frac{1}{2} \frac{n_e}{Z}
\]
FIGURE 13

Idealized line of best fit, plotted on the electron density-ion density plane. The intersection of the line of best fit for each Lyman series line gives the unique values of electron and ion densities which characterize the plasma conditions.
The above inequality defines an excluded region in the electron density-ion density plot. For any given value of electron density, there is a certain maximum amount of opacity broadening which can be obtained. This excluded region is indicated on Figure 13. A more complete collisional-radiative model also defines an excluded region similar to the above condition.
CHAPTER IV
RESULTS

The methods described in the previous two chapters were used to analyze the experimental spectra obtained from the pellet implosions of the targets listed in Table 1. The spectra covered the 3-4 keV energy range which included the principal series lines ($ls-n\ell$ or $ls^2-lsn\ell$ for $n<6$) of $Ar^{+17}$ and $Ar^{+16}$. A region of the spectrum which includes the hydrogenic and helium-like Lyman alpha transitions is shown in Figure 14, a densitometer trace from shot 7497, which imploded a 150$\mu$m $\times$ 2$\mu$m Ar/Ne filled target. Note that the 3.5eV instrumental resolution allows separation of the fine structure components in $Ar^{+17}$ Ly-$\alpha$, and detailed structure is also visible among the satellite line components.

Lineshape fits have been made for all the hydrogenic and helium-like lines observed following the "line of best fit" method. Intensity ratios of lines were also measured, including satellite lines, which are sensitive to temperature or density. Results from the several different diagnostics were then compared for consistency.

Results of Line of Best Fit Method

A number of lineshape fits were carried out for each observable series line so that a line of best fit could be constructed on a $Ne-N_{ion}$ plot. Each theoretical profile was fitted to the data using a least squares procedure to determine relative intensity factors necessary to
obtain a high quality fit over the entire line profile. The underlying continuum was fit both by the least squares method and by manual estimation. It was found that the most satisfactory results were obtained by manually inserting the background coefficients. Quality of a fit was defined as a normalized mean square deviation, i.e.

\[ Q = \frac{\sum (y_{\text{exp}} - y_{\text{th}})^2}{\sum y_{\text{th}}^2}, \]

where \( Q < 0.05 \) represents a good quality fit. Although no specific provision was made to ensure that area normalization was maintained, the results of our fitting procedure were in good agreement with that constraint. The fits were complicated in some cases by the overlapping of neighboring lines (e.g., Ca or K impurities from the glass shell or lines of helium-like Ar\(^{16}\)). The effects of overlapping lines were treated by adjusting relative intensities, again according to a least squares procedure. In Figures 15 (a-d) some Ar\(^{17}\) Lyman series line fits for shot 7499 are shown. These line fits are representative of those used in the construction of a line of best fit for each spectral series line. Recall that these lines of best fit are generated by a sequence of line fits in which the opacity and electron density are systematically varied.

Image analysis of X-ray microscopic data determined the approximate diameters of the imploded cores to be 40-45 microns. These core sizes were then included in the expression for opacity [equation (3-15)] to calculate the density of Ar\(^{17}\) ions in the ground state where the ratio \( \tau_0/L(0) \) was determined from the best fit for a given electron density. Electron temperatures of approximately 1 keV were deduced from
FIGURE 14

Densitometer trace showing the spectral region, including the resonance lines of \text{Ar}^{+7} and \text{Ar}^{+6}.
FIGURE 15

Fits to the Lyman series lines from shot 7499 are shown. Values for optical depth at line center, quality of fit, electron temperature and density are indicated.
the ratio of the Lyman alpha satellite to resonance line (see Chapter II).

The lines of best fit for shot 7497, which employed a \(150\mu m \times 2\mu m\), 50/50 Ar/Ne target, are shown in Figure 16(a). It is observed that a density inference of \(n_e = 1.5 \times 10^{23}\) cm\(^{-3}\) can be made on the basis of a consistent best-fit analysis for the Lyman-\(\alpha\), \(\beta\), \(\gamma\), and \(\delta\) lines; similarly a consistent density of Ar\(^{+17}\) ions in the ground state was found to be \(~3 \times 10^{21}\) cm\(^{-3}\). It is estimated that a 20% uncertainty in the theoretical line profiles results in a corresponding uncertainty in the lines of best fit (LBF). Errors in emission region size will directly affect the ion density determination but the electron density and the general trends indicated are not sensitive to this dimension. It was assumed that all Lyman series lines were emitted from the same emission region, so all the LBF'S would be affected equally by errors in size.

The ionization-limit curves shown in Figure 16(a) correspond to 30% and 50% of the total argon ions in the ground state of Ar\(^{+17}\). An upper bound on the Ar\(^{+17}\) ionic ground state density appears to be 50%, regardless of the ionization model. Hence the 50% line defined in the \(N_e - N_i\) plane represents a limit: all physically realizable results are excluded from the region to the right of that curve. The LBF analysis of shot 7497 implies an Ar\(^{+17}\) ground state density which includes roughly 30% - 50% of the total argon ions.

The lines of best fit for shot 7496, which employed a \(100\mu m \times 3\mu m\) 100% Ar filled target, are shown in Figure 16(b). It is observed that the Ly-\(\alpha\), -\(\beta\), and -\(\gamma\) LBF are consistent with \(N_e = 2.5 \times 10^{23}\) cm\(^{-3}\) and \(N_i\)
\[ 2.1 \times 10^{21} \text{ cm}^{-3} \]. As in the previous case (shot 7497), less than 50\% of the total argon ions are in the \( \text{Ar}^{+17} \) ground state.

In the implosions of these thick shelled target, the mean free path of hot electrons is reduced, resulting in less heating of the core and a slower more ablative implosion, generally reaching higher densities.\(^8\) The hydrogenic x-ray spectra from the implosions of 150\( \mu \text{m} \times 2\mu \text{m} \) and the 100 \( \times \) 3\( \mu \text{m} \) targets allowed a self consistent electron and ion density determination from the analysis of all observed Lyman series members. In fact, this is the first time that analysis of the optically thick Lyman alpha line \((\tau_0 = 100)\) has been consistent with that for the other series members. We attribute this improvement, at least in part, to the inclusion of fine structure corrections and to an extended treatment of the line wings. In these more ablative implosions the conditions of line formation appear to be equivalent, or nearly so, for all the Lyman series members; time integration effects of the spectral data produces no serious inconsistencies. Hence, in the "right" system an optically thick line can be used as a density diagnostic.

Next we consider two examples of experiments which employed 100\( \mu \text{m} \times 1\mu \text{m} \) targets. Figures 16 (c) and (d) illustrate lines of best fit for shots corresponding to Ar/Ne mixtures of 100/0 and 25/75, respectively. In these shots the highest series lines, \( \delta \) and \( \epsilon \), are observed to correspond to average densities that are best determined to be \(< 3 \times 10^{22} \text{ cm}^{-3}\). This density is significantly lower than that inferred from the Lyman-\( \alpha \), -\( \beta \), and -\( \gamma \) lines which imply \( N_e > 5 \times 10^{22} \text{ cm}^{-3} \), assuming consistency with the 50\% ionization limit. For lower electron densities fits for these lines would require ion ground state densities in the forbidden region, e.g. to the right of the 50\%
FIGURE 16

Results of the line of best fit procedure applied to the analysis of Ar$^{+17}$ Lyman series line, produced in four different shots. Consistency is achieved only in the (a) and (b) results. The dashed limit lines refer to constraints imposed on $N_I$ through the use of ionization equilibrium models. $\tau_0$ signifies the calculated opacity at the Lyman-\(\alpha\) line center which is characteristic of the indicated point on the line.
line. In any event the δ and ε lines are much too narrow to agree with the density obtained from the β and γ lines (also see Figure 15).

The results of the line fitting procedure when applied to each of the 100 μm × 1μm targets did not lead to single valued estimates of the electron and ion densities. In these implosions the thin glass shell is exploded by hot electrons which deposit their energy throughout the shell. The results suggest that different Lyman series lines are most intense at different time intervals, and hence their profiles are representative of different temperature-density regimes. Thinner shells tend to increase preheat of the core, so we may be observing the effects of altered level populations due to fast electron preheat in the 100μm × 1μm cases. That is, the most intense emission from the higher lying Lyman series members may occur earlier in the thin shell targets and at higher temperatures and lower densities. High resolution time resolved spectroscopy is required to determine the details of these observed differences. The electron density inferred from the intersection of the Lyman alpha LBF and the 50% ionization limit line was found to be about $1 \times 10^{23} \text{ cm}^{-3}$, regardless of fill ratio. Similarly, density inferences from the other Lyman lines showed no sensitivity to fill ratio but were also inconsistent with the Lyman-α result. This suggests that opacity may not be as important in the diagnosis of these shots as are the effects of implosion dynamics.

Our analysis has shown that spectroscopic plasma diagnostics using Stark and opacity broadening of the Lyman series lines from the compressed gas fill produces consistent results with the thick-shell (> 2 μm thickness) targets, and that inconsistencies occur with the thin-shell targets. It is suggested that these inconsistencies are due
to the effects of implosion dynamics which result in the lines from high-lying \( n \)-states implying significantly lower densities than do those from the lower \( n \)-states. The time integrated spectra is consistent with a conjecture that during these implosions the line emission occurs in more than a single burst, each corresponding to different temperature density regimes. Thus, although it had been conjectured that lines from the highest observable series members, being the most optically thin, would provide the most accurate determination of core density, this study has been shown that lines from the lower lying \( n \) states (e.g. Ly-\( \alpha \), -\( \beta \), -\( \gamma \)) are probably more correct in determining final core densities.

Our analysis of shots with similar target size but varying fill ratio showed similar spectral characteristics. While it was shown that opacity effects must be included for all lines, the variation of fill ratio over the range of fill pressures considered did not produce significant effects on the validity of the overall diagnostics. Opacity may not be as important in the diagnosis of these shots, however, as are the effects of implosion dynamics. Future high resolution time resolved spectroscopy should provide further insights into the implosion dynamics.

The results of the opacity broadening analysis of hydrogenic series lines are summarized in Table 2. The results for opacity broadening analysis of the principal series lines of \( \text{Ar}^{16} \) are less complete due to the fact that the potassium and calcium lines from the glass shell obscured both the helium-like \( \beta \) and \( \gamma \) lines in several shots. The helium-like \( \delta \) line was also blended with the hydrogenic \( \text{Ar}^{17} \) Lyman \( \beta \) line. An attempt was made, however, to apply the opacity broadening
method to all usable data and the results are summarized in Table 3. In most cases, only two helium-like lines (α and β) could be analyzed except for shot 7497 which also included the γ line. Though the data is incomplete, the indications are that the helium-like emission occurs at electron densities almost equivalent to that of the hydrogenic emission, while the helium-like ground state densities seem to be somewhat lower than the hydrogenic ground state density. For temperatures above 1 keV, this observation tends to agree with results based on the optically-thick, collisional-radiative steady-state model discussed in Chapter II.

Satellite Line Results

Also in Chapter II, there was a discussion of the dependence of satellite line emission on electron temperature and density in the plasma. In particular, for the argon lithium-like and helium-like satellites, it was shown that relative satellite line ratios are density dependent (Figures 8, 9) and the satellite-to-resonance line ratios are temperature dependent. These features were used to analyze the satellite emission for the shots given in Table 1 in order to diagnose electron temperature and density. The satellite line diagnostics were then compared to those from the line broadening analysis for each shot.

Two complementary methods were used in the satellite line analysis: (a) fitting the data to theoretical calculations of total satellite emission, and (b) de-convolution of the data to extract measurements of several satellite line ratios which are density dependent. Both methods utilize the same theoretical calculations of the satellite emission and produce similar results. The first method
<table>
<thead>
<tr>
<th>Shot no.</th>
<th>Target size (µm)</th>
<th>N_e (line broadening) (cm⁻³)</th>
<th>N_{Ar+17}^{ground state} (cm⁻³)</th>
<th>N_e (satellite) (cm⁻³)</th>
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<tbody>
<tr>
<td>7499</td>
<td>102.00x0.87</td>
<td>&lt;1x10^{23}</td>
<td>&lt;1x10^{21}</td>
<td>2x10^{23}±30%</td>
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<tr>
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<td>105.00x0.93</td>
<td>&lt;1x10^{23}</td>
<td>&lt;1x10^{21}</td>
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<tr>
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<td>105.00x0.93</td>
<td>&lt;1x10^{23}</td>
<td>&lt;1x10^{21}</td>
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<thead>
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<th>shot no.</th>
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<th>$N_{\text{Ar}^{+16}}$ ground state</th>
<th>$N_e$ (satellite)</th>
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<td>$&lt; 5 \times 10^{23} \alpha, \beta, \gamma$</td>
<td>$&lt; 6 \times 10^{20}$</td>
<td>$3.5 \times 10^{23} \pm 15%$</td>
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<td>7500</td>
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<td>$3 \times 10^{23} \pm 30% \alpha, \beta$ only</td>
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<td>7496</td>
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<td>$2 \times 10^{23} \pm 20%$</td>
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**TABLE 3:** Electron and ion densities as determined from line broadening and satellites of the helium-like argon series lines.
utilizes a least-squares fitting procedure over the entire range of satellite emission, while the deconvolution method considers line ratios over a part of the satellite emission. The deconvolution method, however, is based in a more formal numerical procedure and produces more accurate error estimates.

Several assumptions have been made concerning satellite line shapes and widths. Individual satellite lines were assumed to be optically thin with Voigt or Lorentz profiles. The effects of opacity can generally be ignored due to the low populations of the autoionizing levels. However, under conditions of strong satellite emission, opacity effects may be evident in the strongest lines and this can affect line ratio diagnostics. \textsuperscript{36} Calculations of Stark broadened lineshapes for lines from autoionizing levels have not yet been made, but some previous work\textsuperscript{37} has shown that the isolated satellite line width can be estimated by adding the partial widths arising from autoionization, spontaneous emission, and collisions. For this experiment it appears that individual satellite lines have shapes determined primarily from instrumental and source broadening, which when convolved with the natural and collisional shape produces a Voigt profile (or Lorentzian for simplicity) of full width at half-maximum of approximately 5 eV. Due to the uncertainties in shape and width of satellite components, the sensitivity of the results to shape and width was tested. In general, the line ratio measurements were not found to be sensitive to shape and width when the width was allowed to vary within a narrow range of the expected width (as determined from isolated, unblended satellite lines).
The method of satellite lineshape fitting used the same least squares fitting procedure described previously, where the overall satellite line shape was multiplied by a scaling factor determined by a least squares fit to the data. A lineshape fit using this method is shown in Figure 17 for shot 7499. Although all lineshape details are not matched, a trend that is evident from this and other fits is that the peak satellite emission shifts toward the resonance line as density increased. This is evident in the results of theoretical calculations shown in Figures 8 and 9. Due to the poorer qualities of fit obtained in attempting to match the entire satellite lineshape to the data, this method has somewhat greater estimated errors than the line ratio method.

The relative intensities of the satellite components necessary to determine the ratios indicated in Figures 9 and 10 have been measured by an effective deconvolution process using an iterative peak finding code, RLCFIT. This code fits the background in terms of a Fourier series and determines individual peaks by minimizing a weighted sum of the difference between the data and background with respect to peak heights, positions, and widths. Standard peak shapes employed were either Lorentzian or Voigt. This method serves to enhance the data, but is different from a similar procedure discussed recently by Seely.

This method must be used with caution due to the non-uniqueness of the curve fitting problem. It is therefore necessary to introduce as constraint as much physical information about the spectra as possible. Thus, standard lineshapes were chosen along with widths which were allowed to vary by several electron volts. Though the code could be run with a variable width option, a fixed width was chosen where the standard peaks of fixed width were made to fit the data by varying the
FIGURE 17

Lineshape fit of Ly-α and satellites observed in shot 7499. In this fit, the satellite line profile was calculated for an electron density of $3 \times 10^{23} \text{ cm}^{-3}$ and the Ly-α line profile was calculated for an electron density of $2.5 \times 10^{23} \text{ cm}^{-3}$ and optical depth $\tau_0 = 90$. The quality of fit is 0.002.
background, peak heights and peak locations. It was initially found that the satellite blend was composed of four predominant peaks (Figures 6 and 7). This additional information was incorporated into the code by fixing the spacing of the satellite components, but not the overall location. The line ratios found in this case agreed with those found without this constraint, but the calculated error in the ratio was lower.

This statistical error is determined by a method where all data points are varied randomly within the standard deviation and the entire fitting calculation is redone. Several such re-fits are performed and the results give error estimates for peak positions, widths, and heights. Thus the error in any line intensity ratio is also obtained.

Figures 18(a,b) illustrate this deconvolution process for several shots. The density sensitive ratio \( R \) is the ratio of the sum of the strengths of components (2+3) to (1) in this case. These shots exhibit a blend of satellite lines of the type \( 2\ell 2\ell' - 1\ell 3\ell \) located near the resonance Lyα transition (component 5 in the Figure). These satellites, which tend to distort the lineshape analysis of the resonance line, have intensities between the LTE and coronal limits. Also seen in the figure is an unidentified component near component (1); which though of low strength is statistically significant according to our fitting routine. The unidentified component can be due to both uncertainties in the standard peak shape, and uncertainties in the background fit. The number of coefficients in the Fourier series for the background can be controlled with the input, and it is found that increasing the number of coefficients raises the background level and reduces or eliminates the unidentified component without affecting the line ratio of interest.
The spectral emission of shot 7496 of Figure 18 is an example of a higher density implosion. Note that component (2) is now stronger than component (1) and the peak of the satellite blend shifts toward the resonance line.

The sensitivity of measured line ratios to standard peak widths was checked by running the fitting code for a number of fixed widths within several electron volts of the estimated width. A plot of peak width versus satellite line ratio is shown in Figure 19. The plot indicates a wide region in peak width where the ratio is stable. Fits outside of this range produced results away from the plateau region shown. It was concluded that the data was known accurately enough to deduce a small range for experimental widths for the resolved satellite line components, and that these widths were consistent with previous estimates.

The density dependent helium-like and lithium-like satellite line ratios were measured using the RLCFIT code and techniques previously described. The deduced densities are listed in Tables 2 and 3, together with those inferred from the line broadening results. While the line broadening results are somewhat incomplete for the helium-like series lines, note that in most cases the satellite line ratio indicates a higher electron density than the line broadening analysis. This is also shown in Figure 20 where the helium-like density dependent triplet/\(^1\)D satellite component line ratio is plotted for several different shots on the theoretical curve yielding the electron densities of table 2. The following figure (Figure 21) shows the measured values of the density dependent lithium-like satellite component ratio (acdb+qr)/(jkl) plotted on the theoretical curve and yielding the electron density result for
FIGURE 18
Deconvolution of Satellite Emission

(a) $^{17}$Ar Ly-α resonance line and helium-like satellites observed in shot 7497. Experimental points are indicated by, dashed and solid lines indicate the individual satellite components and the overall theoretical fit determined from RLCFIT. The line ratio $R$ is the ratio of the strengths of components (2)+(3) to (1). The numbered components represent the transitions:

1. $2p^2 \, ^1D - 1s2p \, ^1P$
2. $2p^2 \, ^3P - 1s2p \, ^3P$
3. $2s2p \, ^3P - 1s2s \, ^3S$
4. $2s2p \, ^1P - 1s2s \, ^1S$
5. $2\lambda3\lambda' - 1s3\lambda'$ higher order satellites
Shot No. 7497

Log Intensity (arb. units)

Energy (KeV)

3.25 3.30 3.35
FIGURE 18 (continued)

(b) \( \text{Ar}^{17} \) Ly-\( \alpha \) resonance line and helium-like satellites observed in shot 7496.
Shot No. 7496

Log Intensity (arb. units)

Energy (KeV)
FIGURE 19

A plot of the satellite line ratio, $^{1}D/\text{triplet}$ of equation 2-28, determined from the data of shot 7497 using the code RLCFIT with fixed standard peak widths. The satellite line ratio is seen to be insensitive to this parameter over the range 4.5 to 6 eV.
Shot 7497 \textsuperscript{1}D/triplet Ratio

![Graph showing \textsuperscript{1}D/triplet Ratio vs. Standard Peak Width (eV).](image)
Helium-like satellite line ratio, $R$, versus electron density, $N_e$, calculated at $T_e = 1000$ eV. Measured values of $R$, and the uncertainty, for four shots are plotted. The electron density and its error can then be determined from the plot.
FIGURE 21

The lithium-like satellite line ratio, \((\text{abcd}+qr)/(\text{jk}l)\), versus electron density calculated for \(T_e = 1000\) eV. Measured values of this ratio and uncertainties are plotted for several shots. The electron density and its error can then be determined from the plot.
these shots in Table 3. The theoretical curve was calculated using the
model described in chapter II and the rate coefficients given in the
Appendix.

The fact that the densities deduced from satellite line ratios and
line broadening analysis disagree is a new result but not entirely
unexpected. It is probably a real effect and due to the fact that
derfrent densities are being sampled in the different spectral
emission. Recent low resolution time resolved results\(^{40,41}\) show that
satellite emission occurs during a much shorter time interval than the
Lyman series emission. Our time integrated results thus sample the
Lyman series emission over a range of conditions, with maximum emission
not necessarily under equivalent conditions as the satellite emission.
Further high resolution time resolved data is required to compare the
line broadening analysis to satellite emission produced at the same time
interval. An additional point to consider is the reliability of the
atomic data and satellite line model. Errors in temperature
determination and collisional excitation and ionization rates can
possibly account for some of the discrepancy in the present results.
Calculations of these rates by several different groups do not seem to
differ appreciably,\(^{1,6,14,46}\) and future time resolved spectra may be
used to verify these rate calculations.

It can be seen from Table 3 that the satellite line results show an
opacity dependence in the 100\(\mu m \times 1\mu m\) targets. In particular the
indicated density progressively decreases as the fill ratio of argon is
decreased. In Figures 22(a) and (b) it can be seen that the satellite
components abcd and qr appear to be weaker in the target of less argon
fill. This effect may be partially explained, however, by the opacity
of the intercombination line which for these shots is estimated to vary from an optical depth at line center which is greater than one to a value that is less than one as the fill is lowered. This extra opacity broadening of the intercombination line may account for some of the apparent intensity in the region of these satellite lines, though it does not appear to be able to account for the entire effect.

The complex of lines in the region of the helium-like resonance provides several important diagnostic possibilities. As shown in Figure 22 this complex includes the helium-like resonance, intercombination, and lithium-like satellites. The intercombination line is also blended with the m, n, s, t satellites and the forbidden component of the resonance transition 1s2s 1S - 1s2 1S. This forbidden component, which is included in the lineshape calculation, is found to increase in relative intensity to the resonance line as the opacity is increased. Thus under conditions of high resonance line opacity, the forbidden component and intercombination lines have similar relative intensities and are also blended with the satellites. This entire complex of lines thus provides estimates of electron density through an analysis of satellite line intensities, and resonance to intercombination ratio; and estimates of ion densities by an opacity analysis of the lineshape and resonance to forbidden component ratio. Such an analysis will become important when high resolution time resolved data is obtained.

It has previously been indicated that opacity in the satellite lines would affect the line ratio diagnostics, since the calculations presented here have made the optically thin approximation. In order to explain the discrepancy between lineshape and satellite diagnostics it can be argued that opacity in the satellites may be a factor. In fact
FIGURE 22
Deconvolution of Satellite Emission

(a) Ar$^{+16}$ 1s2p - 1s$^2$ resonance line and lithium-like satellites observed in shot 7499, a 100µm × 1µm, 100% argon filled target. Experimental points are indicated by dots; dashed and solid lines indicate the satellite components and the overall theoretical fit determined from RLCFIT. The line ratio $R$ is the ratio of the strengths of components (3+4) to (1+2). The numbered components represent the transitions:

1. j,l
2. k
3. a,b,d,c
4. q,r
5. m,n,s,t,+ IC
6. 1s2p - 1s$^2$ Resonance
Shot No. 7499

Log Intensity (arb. units)

Energy (KeV)

3.10    3.15

1 2 3 4 5 6
FIGURE 22 (continued)

(b) $\text{Ar}^{+16}$ $1s^2p - 1s^2$ resonance line and lithium-like satellites observed in shot 7560, a 100μm × 1μm, 50% Argon filled target.
Shot No. 7560

Energy (KeV)

Log Intensity (arb. units)

3.10

3.15

1234 5 6

1 1 1 1 1
there is some evidence that the satellite line widths as determined from
the RLCFIT code are larger (by up to 1 eV) than expected, particularly
on the high density shots. While opacity broadening can account for
this, it is also true that the components determined in the RLCFIT
analysis are not isolated satellites but blends of several satellite
components which also accounts for the additional linewiths.

Other possible evidence for satellite line opacity comes from the
ratio among several satellites arising from the singlet levels. The
weak singlet transition, 2s ^1S - 1s2p ^1P, was resolved from the noise
on only shot 7497. The following line ratios were measured, and the
calculated corona and LTE limits are given in parentheses:

\[
\begin{align*}
I \left( \frac{2s^2 \, 1S - 1s2p \, 1P}{2s2p \, 1P - 1s2s \, 1S} \right) &= 0.25 \pm 0.02 \quad (0.149, 0.121) \\
I \left( \frac{2s^2 \, 1S - 1s2p \, 1P}{2p^2 \, 1D - 1s2p \, 1P} \right) &= 0.086 \pm 0.015 \quad (0.049, 0.039)
\end{align*}
\]

It is seen that the observed ratios are approximately twice as large as
those expected. This might be explained by assuming that the stronger
singlet transitions are affected by opacity.

Temperature Diagnostics

Stark broadened lineshapes are relatively insensitive to
temperature, a fact which is fortunate since the temperature diagnostics
are the most uncertain. Temperature deduced by several methods are
given in Table 4 for a number of shots. Generally, 1 keV temperatures
are obtained though higher temperatures are found for the thinner-
shelled targets. Although different methods did not all agree, the same trends in temperature were found in all cases.

The satellite \(^1D\) to Ly-\(\alpha\) resonance line ratio (Figure 10) is affected both by resonance line opacity and by time integration effects. Satellite line emission generally occurs at the higher temperatures and thus for a shorter time than the Lyman series emission. The ratio of hydrogenic to helium-like \(\beta\) resonance lines can also be used as a temperature diagnostic (Figure 2). It can be seen however that such a prediction depends on the accuracy of the collisional-radiative opacity model and the attainment of a collisional-radiative steady state. In general, the hydrogenic and helium-like emission may arise from differing conditions, at differing times. Hence time integrated data may then yield unreliable line ratio diagnostics for this particular case. However, the hydrogenic and helium-like opacity analysis yielded fairly consistent conditions, so it will be assumed that the collisional-radiative model can be used to describe the experiment considered here. The prediction for this ratio, \((Ly-\beta/ls3p-ls^2)\) depends on whether the plasma is optically thick or thin, and Table 4 gives results for both cases. Optical thickness is defined here by the resonance line opacity, the rate equations are then modified with an escape factor technique. The optically thick case yields temperatures in closer agreement with the satellite line results.

The slope of the continuum emission making up the background can also be used as a temperature indicator; though the measurement is not always accurate due to other continuum radiative processes making up the background, particularly non-thermal emission from high energy electrons which may arise from the implosion. Thermal continuum emission is due
TABLE 4: Electron temperatures

Electron temperature (keV) from:

<table>
<thead>
<tr>
<th>Shot No.</th>
<th>Target Size (µm)</th>
<th>Lyβ/(1s3p-1s²)</th>
<th>Lyβ/(1s3p-1s²)</th>
<th>Lyβ/(1s3p-1s²)</th>
<th>Lyβ/(1s3p-1s²)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>thin</td>
<td>thick</td>
<td>1D satellite/Res</td>
<td>Continuum Slope</td>
</tr>
<tr>
<td>7499</td>
<td>102.00x0.87</td>
<td>1.3±.2</td>
<td>0.9±.1</td>
<td>1.6±.2</td>
<td>1.6±.3</td>
</tr>
<tr>
<td>7496</td>
<td>99.00x3.00</td>
<td>1.4±.2</td>
<td>0.8±.1</td>
<td>1.2±.2</td>
<td>-</td>
</tr>
<tr>
<td>7497</td>
<td>157.00x1.97</td>
<td>1.2±.2</td>
<td>0.7±.1</td>
<td>0.7±.1</td>
<td>1.2±.2</td>
</tr>
<tr>
<td>7501</td>
<td>102.00x4.53</td>
<td>1.4±.2</td>
<td>0.8±.1</td>
<td>0.9±.1</td>
<td>1.1±.1</td>
</tr>
</tbody>
</table>
to radiative recombination, whose emission has been shown to be given by

\[ \varepsilon(v) \propto \exp\left(-\frac{hv}{kT_e}\right). \]

Thus a plot of \( \ln \varepsilon(v) \) vs \( v \) should be proportional to \( h/kT_e \). The recombination continuum beyond the hydrogenic series limit was measured and the slope was determined by a least squares analysis; temperature estimates from this method are given also in Table 4. It is seen that this method produced consistently higher estimates than the other ones, but the trend among shots was consistent with the other methods.
CHAPTER V
CONCLUDING REMARKS

Time integrated, high resolution x-ray spectra have been analyzed by several different techniques, and the resulting diagnostics have been compared. Theoretical models of the atomic collisional and radiative properties of laser produced plasmas have been developed and related to observable properties of the spectral emission. The effects of opacity have been included in the models and a systematic attempt made to observe these effects in argon plasmas by analyzing the spectra produced by similar targets varying mainly in their total argon fill.

Lineshape fitting using the "line of best fit" method, which includes the effects of Stark and opacity broadening, was used to estimate electron and ion densities in the plasma. Electron densities for the shots analyzed were found to be in the range 1 to $5 \times 10^{23} \text{ cm}^{-3}$. The importance of using all available experimental data was emphasized with a consistent fitting method for all observed spectral lines. Significant differences were observed in the emitted spectra from thin shell and thick shell targets. For the thick shell targets, electron and ion densities were determined consistently from an analysis of all observed series members. For the first time analysis of the alpha resonance line has been found to be in agreement with the other series members. For the thin shell targets the densities inferred from the lineshape analysis of individual series members can differ appreciably. The highest series members which are most optically thin
are not always the best indicators of core or peak density. It was conjectured that this effect may be due to implosion dynamics since a greater number of non-thermal high energy electrons may preheat the core of thin shell target prior to maximum hydrodynamic compression, and lead to more than a single burst of spectral emission. A double peak in the total x-ray emission has previously been noted in neon filled microballoons. Thus, opacity may not be as important in these shots as are the effects of implosion dynamics.

Prominent satellite line emission was analyzed as both a temperature and density diagnostic using appropriate detailed models of satellite line emission. The density dependent satellite line component ratios were shown to yield density estimates which were uniformly larger, a factor of 2 or more, than the corresponding line broadening analysis. Such an effect has been observed before, but at a lower density, and was partially explained by including inner shell collisional excitation in the theoretical calculation. As shown in Figure 8 the effects of inner shell excitation are more important at lower densities and not likely to explain the discrepancies in these experiments. The discrepancy observed in the present experiments is probably a real effect due to the fact that satellite lines are emitted during a shorter time period near peak temperature [recall equation (2-18)] and density than are the resonance series lines. Thus time integration effects are probably important. As was discussed, the possibility of opacity effects in the strong satellite lines could also affect the line ratio diagnostics.

Higher order satellite emission (from 2\ell3\ell' and 1s2\ell3\ell' levels) was found to be prominent and the unresolved components effectively served
to broaden the hydrogenic and helium-like resonance lines. The intensities of these satellites were seen to be between the low and high density limits of the theoretical models.

A collisional-radiative steady-state model was found to adequately represent the plasmas of this experiment during the time of maximum emission. Models which explicitly included photoexcitation effects with an escape factor were described, and it was shown how this affected line ratio diagnostics. The conditions under which hydrogenic and helium-like series line emission occurred were apparently similar, while the satellites appeared to be produced under different plasma conditions. The thin shell targets also appeared to be more affected by time integration effects.

It is now becoming possible to achieve both high spectral resolution and high time resolution (approximately 10 picoseconds) with newly developed spectrometers and streak cameras. When such data becomes available it will be possible to explicitly assess the differences between times of satellite and resonance emission or the importance of the double burst effect on the thin shell target implosions. It will also be more important to use the more detailed models (described in Chapter II) which combine hydrodynamic code results with a collisional-radiative rate equation model. Time-resolved spectroscopic diagnostics will thus provide additional constraints on the hydrodynamic code and rate-equation predictions.
APPENDIX
RATE COEFFICIENTS AND ATOMIC DATA

The solutions to the rate equations of Chapter II, which were used to predict ionization fractions, level populations and line intensity ratios, utilized rate coefficients and other atomic data which were obtained from the literature. In this appendix, a more complete description of the required atomic data is given.

Energies

Energies of ionic levels and spectral lines were obtained from the calculations of Vainshtein and Safranova\(^1\) and Cowan.\(^43\) These calculations use the method of charge expansion perturbation theory which gives the energies for each configuration as an expansion in powers of \((1/Z)\) where matrix elements are evaluated using Coulomb wave functions. This method has the advantage of providing results over an isoelectronic sequence by a simple scaling of \(Z\). The results of Vainshtein and Safranova include energy level terms up to \(n=3\) in hydrogenic and helium-like ions and the autoionizing levels (also including \(n=3\) terms) which give rise to the satellite line emission. Calculations are available in the literature\(^41\) specifically for \(Ar^{+17}\) and \(Ar^{+16}\) which use the results of Cowan's detailed relativistic Hartree-Fock atomic structure code. The results of this model were found to agree (< 1\% difference) with the Vainshtein and Safranova
energy levels. Several other sources\textsuperscript{2,14,45} were also used to obtain energy levels and spectral line energies.

\textbf{Autoionization and Radiative Decay Rates}

Spontaneous radiative decay rates have previously been calculated for the principal series and satellite transitions in argon.\textsuperscript{1,14,44} In these calculations the dipole moment operator was evaluated in terms of the single electron electric dipole radial matrix elements and angular momentum coupling coefficients for the states of interest in each transition; Coulomb wavefunctions were used to evaluate the matrix elements. The spontaneous radiative decay rate is related to the Einstein absorption coefficient and the $f$-value of the transition $i\rightarrow j$ by the well known relations:

\[
A_{ij} = \frac{2\hbar\nu^3}{c^2}B_{ij}
\]

(A-1)

and

\[
\frac{\pi e^2}{mc}f_{ij} = \frac{\hbar\nu}{4\pi}B_{ij}
\]

(A-2)

where $g_iB_{ij} = g_jB_{ji}$. The calculations of Vainshtein and Safranova were used throughout when available.

The process of autoionization, discussed in Chapter II, represents a mixing between doubly-excited bound states and continuum states (see Figure 5). For example a doubly excited bound state of a helium-like ion can be represented $2\ell 2\ell' 1,3L_J$ using LS notation, and the
corresponding continuum state of the same energy can be represented as \( \text{lse} 1,3^3L_j \) where the \( \varepsilon \ell \) electron is unbound. Selection rules for autoionization can be obtained in the LS coupling approximation similarly to electric dipole radiation. It is found\(^{46}\) that autoionization can occur only into specific continuum, such that \( \Delta S=0, \Delta L=0, \Delta J=0, \) and \( \Sigma \ell_{1} \) remains either even or odd. Thus a term like \( 2s2p \, ^3P \) can autoionize into the \( 1seP \, ^3P \) continuum state; but a term like \( 2p^2 \, ^3P \) cannot autoionize into any \( 1se \ell \) continuum.

The transition probability for autoionization can be obtained in a more detailed analysis by the evaluation of the electrostatic interaction matrix elements:

\[
\Gamma_a = \frac{4\pi^2}{\hbar} \epsilon^2
\]

Matrix elements of the electrostatic interaction, \( e^2/r_{12} \), are taken with respect to bound and continuum state wavefunctions. In the results of Vainshtein and Safranova,\(^1\) Coulomb radial wavefunctions were used in the evaluation of eq. A-3 matrix element for the bound states and strict LS coupling was not assumed. In a similar analysis,\(^{47}\) bound state wavefunctions were calculated using a semi-empirical approximation. The continuum electron states were calculated with distorted wavefunctions,\(^{48}\) which account for the non-coulomb electrostatic interaction between the free electron and the screened electrostatic field of the ion. The autoionization rates of Vainshtein and Safranova\(^1\) were used throughout this analysis. These rates were in good agreement with other available calculations.\(^{39}\)
Electron impact excitation and ionization rates are available in the literature for a number of transitions of interest. Some of the rates used in the rate-equation analysis of Chapter II are:

i) collisional excitation and deexcitation transitions among bound states.

ii) inner shell excitation transitions from singly to doubly excited.

iii) collisional excitation and deexcitation transitions among doubly excited levels.

iv) electron impact ionization from singly or doubly excited levels.

Cross sections and collision rates for transitions of this type can be calculated using time dependent perturbation theory with the method developed by Sampson and co-workers.\textsuperscript{14,15,16,49} In this method, the zero order Hamiltonian for a system of a free electron scattering off an ion of nuclear charge $Z$ and $N$ bound electrons is,

$$H_0 = \sum_{i=1}^{N+1} \Delta E - \frac{1}{2} \left( \frac{\mathbf{p}_i^2}{\mathbf{r}_i} + \frac{2}{\mathbf{r}_i} \right),$$

and the perturbation is due to the electrostatic interactions among the electrons plus all relativistic corrections,

$$H^1 = \sum_{i>j} \frac{1}{Z \mathbf{r}_{ij}} + \text{relativistic terms}$$

where $\mathbf{r}_i$ is the position of the $i^{th}$ electron relative to the nucleus and...
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$r_{ij}$ is the magnitude of the distance between the $i^{th}$ and $j^{th}$ electrons. Extended atomic units are used where distances are in units of $a_0/Z$ and energies in units of $2Z^2$ Rydbergs. It is clear that for $Z \gg N$ the perturbation Hamiltonian becomes infinitesimal, and time dependent perturbation theory can be used to obtain an electron impact excitation cross section of the form,

$$Q_{ij} = \frac{1}{16\pi} \frac{k'}{k} |H_{ij}'|^2$$

where $k$ and $k'$ are the initial and final wavenumbers of the free electrons. Matrix elements of $H'$ are evaluated using zero order wavefunctions, which are formed from the $N$ hydrogenic bound state wavefunctions and a free Coulomb wavefunction. Sampson and co-workers have shown how cross sections for transitions in complex ions can be related to hydrogenic ion matrix elements and easily calculated for a wide range of $Z$ ($4 < Z < 74$). In particular, results are available for collisional mixing and inner shell excitation of the $2l2l'$ and $1s2l2l'$ autoionizing levels, collisional excitation among singly excited bound helium-like and lithium-like states, collisional ionization from hydrogenic and helium-like bound states, and others.

Collision rates were obtained from cross sections by integrating over a Maxwellian velocity distribution,

$$C_{ij} = N e \sqrt{\frac{8kT}{\pi m}} \int_0^\infty \frac{E}{kT} e^{-E/kT} Q_{ij} d\left(\frac{E}{kT}\right)$$

The effects of a non-Maxwellian hot-electron velocity distribution would thus affect these collision rates, but it is expected that in
these experiments less than 5% of the electrons are non-thermal. Dense plasma effects have also been neglected in the collision rate calculations. They would affect the autoionization and electron impact excitation rates most directly since the free electron wavefunctions used in these calculations would become distorted by the dense plasma effects.

Collisional deexcitation and three body recombination rates have been calculated by applying the principle of detailed balance to the collisional excitation and ionization rates, respectively. For the collisional excitation and deexcitation rates this takes the form

\[ N_1^* C_{12} = N_2^* C_{21} \]  

where the * indicates LTE populations. For the three body recombination rate the detailed balance equation involves the Saha-Boltzman equation.

**Radiative Recombination**

The rate coefficient of radiative recombination was calculated from the hydrogenic coefficient of continuous emission, with an approximate screening factor to account for non-hydrogenic behavior. The form used here is

\[ R = 5 \times 10^{-14} Z N_e y^{3/2} (1 - \frac{n}{2 y}) e^y \int_y^\infty \frac{e^{-x}}{x} \, dx \, sec^{-1} \]  

where \( R \) is the radiative recombination rate into level \( n \), \( I_\alpha \) is the
ionization potential for level \( n \), \( y = I_n / kT \), and \( P_n \) is the number of electrons in the state with principal quantum number \( n \).

The inverse process, photoionization, can be calculated by detailed balance and the Saha-Boltzmann equation.
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BIOGRAPHICAL SKETCH

Norman David Delamater was born on September 18, 1950, in Harrisburg, Pennsylvania. He graduated from Susquehanna Township High School in June, 1968. In June, 1972, he received a B.S. degree in physics from Rensselaer Polytechnic Institute and in December, 1975, he received a M.S. degree in astronomy from the Ohio State University. In September, 1976, he enrolled in the Graduate School of the University of Florida.

Norman David Delamater is married to the former Yvonne Dorothy Shearer. He has been a member of the American Physical Society since March, 1979.
I certify that I have read this study and that in my opinion it conforms to acceptable standards of scholarly presentation and is fully adequate, in scope and quality, as a dissertation for the degree of Doctor of Philosophy.

[Signature]
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Professor of Physics

I certify that I have read this study and that in my opinion it conforms to acceptable standards of scholarly presentation and is fully adequate, in scope and quality, as a dissertation for the degree of Doctor of Philosophy.

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Professor of Physics

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