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PHYSICS OF SHORT WAVELENGTH LASER DESIGN

by

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Submitted to the Department of Electrical Engineering and Computer Science on January 30, 1981, in partial fulfillment of the requirements for the Degree of Doctor of Philosophy

ABSTRACT

We describe the physics and design of VUV and soft X-ray lasers pumped by ICF class high intensity infrared laser drivers (for example, the SHIVA laser facility at LLNL). Laser design and physics issues are discussed in the case of a photoionization pumping scheme involving Ne II and line pumping schemes involving H-like and He-like neon.

The Ne II laser scheme is novel and is pumped by preferential photoionization of 2s electrons from Ne I, resulting in a 2s-2p inversion at 27 eV in Ne II. The target design implementing the scheme is based on thin film flashlamp and filter systems which emit strong X-radiation near 500 eV when irradiated with short pulse (90 psec) 1.06μ radiation at 10^{14} W/cm². The design simulations include thin film irradiation calculations (carried out on the LLNL laser fusion code) and two dimensional time-dependent population kinetics studies. Gains in excess of 50 cm⁻¹ are calculated.

The line pumping schemes involve higher energy transitions (55 eV, 151 eV and 189 eV), and are pumped by resonant L-M lines of medium Z (Z near 28) elements. A detailed kinetics model is developed using atomic physics data available in the literature, and from Hartree Fock and distorted wave models. A novel approach to laser target design is presented, based on the production of very large velocity gradients which result when the strong suprathermal electron momentum density is coupled into the ion fluid momentum density via ionization gradients set up when shadows are present. The velocity gradients relieve radiation trapping problems and allow greater penetration of resonant pump radiation. Conditions under which gain of 50 cm⁻¹ can be achieved are discussed.

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TABLE OF CONTENTS

I. Introduction	23
A. Background	26
B. Scope of Present Research.	40
C. Thesis Overview.	45
D. Bibliography	60
1. Short Wavelength Laser Schemes	60
a. Reviews.	60
b. Collisional Pumping.	62
c. Recombination Pumping.	65
i. Hydrogen	65
ii. Hydrogen-like Ions	66
iii. Hydrogen-like Carbon	68
iv. Helium-like Ions	70
v. Li-like.	71
vi. Other Systems.	72
d. Photoionization Pumping.	73
e. Photoexcitation Pumping.	76
f. Charge Exchange.	80
g. Grasers.	83
2. Mirrors for Short Wavelength Lasers.	84
3. Other Background References.	86
a. Laser Pump Sources	86
b. Z-pinch Sources.	86
c. Future Applications.	87

II. Neon Atomic Processes	88
A. Energy Levels	92
B. Wavelengths and Oscillator Strengths for 2s-2p Transitions.	96
C. Collisional Excitation for 2s-2p Transitions.	102
D. Photoionization	115
E. Electron Collisional Excitation and Ionization Processes	126
F. L-shell Double Processes.	138
G. K-shell Auger Processes	143
H. Bibliography.	144
1. Energy Levels, Wavelengths, Oscillator Strengths.	144
2. 2s-2p Collisions.	149
3. Photoionization	150
4. Electron Collisional Excitation and Ionization Processes	154
5. L-shell Double Processes.	157
6. K-shell Auger Processes	160
III. Neon Laser Design	162
A. Material Selection.	169
B. Flashlamp-filter Operation.	175
C. Suprathermal Electrons Within the Gain Medium	187
D. Laser Kinetics and Gain	190
E. Bibliography.	222
1. Short Wavelength Laser Papers	223

2. Molybdenum Spectra Data	223
3. LASNEX Physics.	225
4. Laser-Disk Analysis	226
5. Light Absorption.	228
6. Suprathermal Electron Generation.	231
7. Transport Inhibition.	233
IV. Resonantly Pumped Photoexcitation Schemes	235
A. The Spectroscopy Problem.	244
B. Atomic Physics Parameters	264
C. Line Transfer and Hydrodynamics	277
D. Laser Design in the Case of the Ne IX 4-3 Laser	289
E. Laser Design in the Case of the Ne IX 3-2 Laser	343
F. Laser Design in the Case of the Ne X 3-2 Laser.	357
G. Bibliography.	371
1. X-ray Laser References.	371
2. Spectroscopy.	372
a. Reviews and Compilations.	372
b. He-like Lines	374
c. Li-like Lines	376
d. Be-, B-, C-like Lines	377
e. O-like Lines.	379
f. Fluorine-like Lines	382
g. Neon-like Lines	383
3. Laser Plasma L-shell Spectra and Intensity, etc	385
4. Non-LTE Models.	386
a. H-like Ions	386

b.	He-like Ions	.389
c.	Low Z Elements	.390
5.	Energy, Level and Radiative Data.	.393
6.	Electron Collisional Excitation.	.396
a.	Reviews.	.396
b.	H-like Ions, Coulomb Born Apprx.	.398
c.	H-like Ions, Distorted Wave and Close Coupling Apprx.	.400
d.	He-like Ions	.401
e.	Li-like Ions	.403
f.	Be-like Ions	.407
g.	Other Sequences.	.409
h.	Collections.	.410
i.	Classical Path Model	.412
j.	Numerical Methods.	.414
7.	Electron Collisional Ionization.	.416
a.	H-like Ions.	.416
b.	Bethe Approximation.	.418
c.	He-like Ions	.419
d.	Li-like Ions	.420
e.	L-shell Results.	.421
f.	M-shell Results.	.423
g.	Neutrals	.424
h.	Binary Encounter	.427

8. Ion Collisional Excitation	.429
9. Line Transfer.	.430
10. Hydrodynamics.	.434
11. Rayleigh Taylor Instabilities.	.435
V. Conclusions and Suggestions for Further Work	.436

LIST OF FIGURES

I. Introduction	
I-1	Single Electron Photoionization Cross Sections for Neon 47
I-2	Ne II Photoionization-pumped Lasing Scheme 48
I-3	SPIVA Beam Positioning for Short Wavelength Laser Targets. 49
I-4	Resonantly-pumped Photoexcitation Scheme Involving Helium-like Ne IX 53
II. Neon Atomic Processes	
II-1	Cross Section and Rate Coefficient for the Ne II $2p_{1/2}-2s_{1/2}$ Transition. 109
II-2	Comparison of Theoretical and Experimental 2p Single Electron Photoionization Cross Sections for Ne I 119
II-3	Theoretical Photoionization Cross Sections (Per Electron) for 2p Electrons for Neon Ions . . 120
II-4	Single Electron 2s Photoionization Cross Sections for Neon Ions. 122
II-5	Ne I Single Electron Collisional Ionization Cross Sections and Rate Coefficients. 129
II-6	Collisional Excitation and Ionization Per Electron for Cross Sections Ne I 2p Electrons 130
II-7	Cross Sections Per Electron for Electron Collisional Ionization and for Excitation and Ionization. 132
II-8	Collisional Data (Per Electron) for 2s Ionization and for 2p Ionization Excitation and Ionization of Ne I 133
II-9	Estimates of Total Electron Collisional Excitation and Ionization Data for Neon 2p Electrons (Per Electron) 134

II-10	Estimates of Total Collisional Excitation and Ionization Data for Neon 2s Electrons (Per Electron)	135
II-11	Ratio of Two Electron Ionization Cross Section to Single Electron Ionization Cross Section for Low Energy Ne I Ionization	140
II-12	Empirical Shake-off Cross Sections for L-shell Photoionization of Ne I	142
III. Neon Laser Design		
III-1	Ne II Laser Target Design in Cylindrical Geometry	163
III-2	Angle Averaged Transmission of 1 μ LiF Between 20 eV and 1 keV.	172
III-3	Optical Depth of Constituents for a 1 μ LiF Filter	173
III-4	Benchmark Design for Flashlamp and Filter System Displayed in Planar Geometry.	174
III-5	Density of Flashlamp and Filter System as a Function of Space at Six Times.	179
III-6a	Electron Temperatures of the Molybdenum Zones as a Function of Time.	180
III-6b	Electron Temperatures of the Beryllium and LiF Zones as a Function of Time.	181
III-7	Positions and Velocities of the Molybdenum Zones as a Function of Time.	183
III-8	Calculated Time-integrated Spectra Out the Backside of the Target	184
III-9	Power Radiated Out the Backside of the Target	185
III-10	Photoionization Rates Per Electron at 5×10^{13} W/cm^2 and 10^{14} W/cm^2	191
III-11	Electron Temperatures of the Be Zones	192
III-12	Schematic of Code System Used in X-ray Laser Target Design.	194
III-13	Gain on the 460.7 $\text{\AA}$ Line at $5 \times 10^{13} \text{ W/cm}^2$	196

III-14	Photoionization Rates (sec^{-1}) at the Time of the Peak Gain	198
III-15	Electron Density as a Function of Time at the Target Center.	199
III-16	Fractional Populations as a Function of Time at the Target Center.	201
III-17	Average Energy of Photoejected Electrons as a Function of Time at the Target Center . . .	202
III-18	Small Signal Gain as a Function of Time at the Target Center	203
III-19	Photoionization Rates (sec^{-1}) at 10^{14} W/cm^2	205
III-20	Fractional Level Populations at the Target Center at 10^{14} W/cm^2	206
III-21	Small Signal Gain at the Target Center at 10^{14} W/cm^2	207
III-22	Optimization of the Small Signal Gain as a Function of Neon Density at $5 \times 10^{13} \text{ W/cm}^2$. . .	208
III-23	Optimization of the Small Signal Gain as a Function of Neon Density at 10^{14} W/cm^2	209
III-24	Sensitivity of Gain to Electron Temperature at $5 \times 10^{13} \text{ W/cm}^2$	211
III-25	Sensitivity of Gain to Electron Temperature at 10^{14} W/cm^2	212
III-26	Monte Carlo Calculation of the Ne II Ion Temperature.	213
III-27	Fractional Isoelectronic Populations for LiH Target	216
III-28a	Gain on the Ne II $2S_{1/2}-2P_{3/2}$ Line for the LiH Target	218
III-28b	Gain on the Ne III $1S_0-2sp^5 1P_1$ Transition at 23.47 eV.	219

III-28c	Gain on the Ne III $2s^5 1p_1 - 2s^4 1d_2$ Transition at 32.69 eV	220
III-28d	Gain on the Ne IV $2p^5 2p_{3/2} - 2s^4 2p_{3/2}$ Transition at 20.45 eV	221

IV. Resonantly Pumped Photoexcitation Schemes

IV-1	Resonant Photoexcitation Pumping of Helium-like Ne IX.	237
IV-2	Line Positions of the K-, L- and M-shell Lines of the Ions.	246
IV-3	Spectra in the Vicinity of the Helium-like Neon K-shell Lines	247
IV-4	Spectra in the Vicinity of the Helium-like Fluorine K-shell Lines	249
IV-5	Flashlamp and Filter System for the Line Pumping Schemes.	255
IV-6	Temperature and Density Profiles 10 psec After the Pumping Pulse Maximum.	256
IV-7	Modal Photon Densities Produced by an Optically Thick Plasma in LTE.	259
IV-8	Electron Collisional Excitation Cross Sections of the n-n Transitions of the Singlet States in Ne IX.	271
IV-9	Two Dimensional Mesh for Hydrodynamics and Line Transfer Calculations	278
IV-10a	Density Profiles of the Inner Wall of the Flashlamp Filter System Expanding into Vacuum	293
IV-10b	Expansion of the Inner Filter Wall into the Laser Medium	295
IV-10c	Shock Position and Velocity as a Function of Time.	297

IV-10d	Width of the Unshocked Laser Medium as a Function of Time	298
IV-11a	Electron Density Versus Time at the Target Center	302
IV-11b	Fractional Populations of the Neor-like Through Helium-like Sequences	303
IV-12	Fractional Ion Populations of the Final Sequences	304
IV-13a	Ion Temperature as a Function of Time	305
IV-13b	Ion Temperature Due to Proximity Heating	306
IV-14a	Fractional Level Population of the Helium-like Ne IX M & N Singlet Levels	307
IV-14b	Fractional Inversion Densities on the Strong 3-4 Lines	308
IV-14c	Gain on the Strong M-N Lines as a Function of Time	309
IV-14d	Stimulated Emission Cross Sections on the Strong 4-3 Transitions	310
IV-14e	Voigt Parameters for the 3-4 Transitions	311
IV-15	Peak Absorption Cross Sections on the Resonance Lines	312
IV-16	Gain on the Strong 3-4 Lines of Ne IX as a Function of Neon Density	320
IV-17	Small Signal Gain on the 4^1P-3^1S Transition as a Function of Time at Different Ion Densities	321
IV-18a	Ion Temperatures as a Function of Time for Different Initial Neon Densities	323
IV-18b	Electron Densities as a Function of Time for Different Initial Neon Densities	324

IV-19	Variation of the Small Signal Gain on the M-N Lines as a Function of Resonant Line Pump Strength	325
IV-20	Electron Density in Space for a Symmetric Target.	327
IV-21	Pressure in Space for a Symmetric Target	329
IV-22	X-directed Velocities for the Symmetric Target	330
IV-23	Y-directed Velocities for the Symmetric Target	331
IV-24	Y-directed Velocities Normalized to the Local Ion Thermal Velocity.	332
IV-25a	Ion Thermal Temperature in Space (Neglecting Proximity Heating).	334
IV-25b	Ion Thermal Temperature in Space	335
IV-26	Gain on the 4^1F-3^1D Transition of Ne IX for the Symmetric Target.	337
IV-27	Electron Density for the Asymmetric Target	339
IV-28	Velocity Flow for the Asymmetric Target.	341
IV-29	Isocontours of Velocity Normalized to the Local Thermal Ion Velocity.	342
IV-30a	Small Signal Gain as a Function of Initial Neon Density for the 151 eV Line.	345
IV-30b	Time Dependent Absorption Coefficients of the $1s^2-1s2p^1P$ Transition.	347
IV-31a	Fractional Singlet L-shell and M-shell Levels in the Case of $1s^2-1s3p^1P$ Pumping	350
IV-31b	Small Signal Gain for the Ne IX 3-2 Laser Scheme.	351
IV-31c	Stimulated Emission Cross Section for the $1s3d^1D-1s2p^1P$ Line.	353

IV-31d	Resonant Absorption Coefficient of the Helium-like Ne IX K-shell Resonance Lines.	354
IV-32	Gain and Fractional Level Populations for the 151 eV Laser in the Case of 1s3p ¹ P-1s3d ¹ D Pumping.	356
IV-33	The Hydrogen-like Ne X 3-2 Resonant, Pumped Laser Scheme.	358
IV-34a	Parameter Study of the Gain at 189 eV With Initial Neon Density as a Parameter.	360
IV-34b	Absorption Coefficients of the 1s _{1/2} -2p _{1/2} Transition at Different Neon Densities	361
IV-35	Fractional Population of the Final Isoelectronic Sequences for the 189 eV Laser	363
IV-36	Fractional Populations of the L-shell and M-shell States of Ne X	364
IV-37	Small Signal Gain on the 2-3 Lines of Ne X at 189 eV.	367
IV-38	Absorption Coefficients for the 1s-2p and 1s-3p Lines of Ne X.	370

LIST OF TABLES

I.	Introduction	
II.	Neon Atomic Processes	
II-1	Neon Energy Levels	93
II-2	Neon Oscillator Strengths and Wavelengths.	98
II-3	Comparison of Semi-classical Results with Other Results for Ne VIII	106
II-4	Example of Rate Coefficient Calculation and Least Squares Fit	110
II-5	Fitting Parameters for 2s-2p Rate Coefficients for Neon Ions.	112
II-6	Fitting Coefficients for Neon Ion L-shell Single Electron Photoionization Cross Sections.	123
II-7	Fitting Coefficients for Neon Collisional Data	137
III.	Neon Laser Design	
III-1	Radiative Energy Balance	186
IV.	Resonantly Pumped Photoexcitation Schemes	
IV-1a	Resonances with He-like $1s^2-1s3p\ ^1P$ Lines.	250
IV-1b	Resonances with H-like $1s-3p$ Lines	251
IV-1c	Resonances with He-like $1s^2-1s4p\ ^1P$ Lines.	252
IV-2	Neon Levels Included Explicitly In Present Calculations.	265

I. Introduction

The advent of the laser in the early 1960's presented a cornucopia of possibilities for technological developments in many areas, few of which were anticipated by its developers. Speculations of possibilities for applications of VUV and X-ray lasers are at present limited to the blurred and dull foresight of researchers in the field, the vivid imaginations of the technological dreamers and the wondrous and exciting universes of science fiction writers. At the time of writing of this section, the author is unaware of any X-ray lasers which have demonstrated the capability of vaporizing intruding spacecraft as in Niven (1970), taking holographic movies of macromolecules as they squirm in their in vivo environments as described by Chapline and Wood (1975) or even to diagnose a 10^{26} electron/cm³ plasma in a compressed inertial confinement pellet as mentioned in the review of Waynant and Elton (1976). In the present research most of the effort has been channelled into the less exciting and perhaps mundane pursuit of a laser design, or more precisely, mundane construction of numerical tools with which a short wavelength laser can be designed. Originally, we had hoped to operate a laboratory EUV laser pumped by ARGUS or SHIVA (in the early days it was CYCLOPS), but such has not come to pass.

But I digress. The general approach to short wavelength laser design which was taken deserves some comment. At the Lawrence Livermore National Laboratory where the bulk of the work was done there is

an energetic laser fusion effort, replete with the worlds' most powerful laser, superb target fabrication facilities, and excellent target design and diagnostics groups. As experiments are costly both in terms of money as well as time and effort, most experiments are carefully planned and calculated in detail before being carried out. The X-ray laser effort was to be carried out in much the same way. The target design would be well thought out and calculated in detail before the experiment. The problem initially was that there were no general purpose X-ray laser design codes available on which to carry out the calculations in preparation for an experiment. One goal of the present effort was to fill in the gap, namely to develop a general purpose X-ray laser design code patterned in some respects after the laser fusion design code LASNEX. A discussion of the laser fusion effort, the lasers and LASNEX are given in the annual reports of the laser program which are referenced at the end of the chapter.

There are many approaches to the problem of X-ray laser design, that is, there are several pumping mechanisms by which a population inversion in the VUV and soft X-ray spectral regions could be pumped. In principle, laser designs derived from most of these schemes could be evaluated numerically with the design code and to date about a dozen schemes have been examined. A discussion of approaches suggested in the literature is given in the following section. As to which approach is best, cheapest, simplest, most efficient, or can get to the highest transition energies are questions which can probably not be answered yet as the field is still in its infancy, or at best its

early childhood. In any case, we shall not attempt categorical evaluation here: although a certain prejudice for photo-pumping schemes may become apparent.

In summary, X-ray lasers would be exciting if they could be made. What we would do with them if they were available is a question which we shall not address. What is of concern here are questions more of how to design them.

A. Background

The field of short wavelength lasers has been reviewed frequently (see the references at the end of the chapter). The review of Waynant and Elton (1976) is one of the more extensive although by now it may be somewhat dated. The NRL group was preparing another more recent review (D. Nagel - Private communication (1979)), however, a copy has not yet become ensconced in our archives. In this section we shall briefly discuss some of the approaches which are currently being pursued.

Electron collisional pumping stirred considerable interest when the Lebedev group reported experimental results involving a 3s-3p transition in neon-like calcium near 600 Å (Il'yuknin et al (1979)). In this experiment a laser pulse of width 2.5-5.0 nsec and intensity $2.0-7.5 \times 10^{10} \text{ W/cm}^2$ was focussed onto a calcium strip whose width was 0.03 to 0.125 mm and whose length was 10-40 mm. A resonator was constructed using lossy mirrors, one a gold coated mirror, the other a ruthenium coated mirror with reflectivities 0.13 and 0.25 at the lasing wavelength. Vinogradov et al (1977) and Vinogradov et al (1978) (see the review section) reports a theoretical analysis. The strong monopole excitation from the neon-like $2s^2p^6 \ ^1S_0$ ground state populates the $2s^2p^53p \ ^1S_0$ state which is inverted relative to the $2s^2p^53s \ ^1P_1$ state. Gains of order 30 cm^{-1} are calculated on this transition using a model which includes $2s^2p^6 \ ^1S_0 - 2s^2p^53s \ ^1P_1$ trapping. Other lines of the 3s-3p manifold achieve gain although the detailed results do not appear. Vainshtein et al (1978) give energy levels for the neon-like excited states for many low Z atoms.

This approach has also been examined in other isoelectronic sequences and for plasmas other than the laser produced variety. Gain in the carbon-like sequence has been examined by Elton (1974), Palumbo and Elton (1977) and Whitney et al (1977). Electron collisional pumping of the 3-2 transitions in the helium-like sequence is discussed by Palumbo and Elton (1977) and Bristow et al (1972). Jha was concerned with the 4s-3p transition in neon-like ions. Although the Lebedev work involved a laser plasma, the electron collisionally pumped schemes may be suitable for Z-pinch plasmas as discussed by Whitney and Davis (1975). Some Z-pinch references are found at the end of the chapter. Our group has examined the 3s-3p transitions in neon-like Krypton as occurs in Z-pinch discharges with the result that the gain at 83 eV may reach gains of several cm^{-1} . Some of the major uncertainties in the estimates lie in the assumptions made about the electron and ion temperatures, densities and plasma radius. This work is still in progress and will be documented in the near future. The 3s-3p approach appears promising at present and it is likely that the future holds promise for interesting experiments along these lines. The extension of the scheme to 4-4 transitions may be straightforward although at present the atomic physics are not as well known for the sequences near the nickel isoelectronic sequence. The use of suprathermal electrons to pump a 3s-3p transition in a laser plasma may be a fruitful approach, although no work has yet been reported on it. The advantage lies in that the suprathermal electron distribution can be better matched to 2p-3p excitation

(under transient stripping conditions) than thermal electrons at the temperature where the ionization stage is dominant.

Another approach involves the removal of inner shell electrons through electron collisional ionization (Lax and Guenther (1972)). In this case it is suggested that K-shell electrons can be preferentially ionized from beryllium or carbon leading to a self terminating (on a few femtosecond timescale) $1s-2p$ transition laser. The approach of using electron collisional ionization as a pumping mechanism is interesting and has not been greatly explored. The application to the K shell self terminating laser problem is not yet within reach of technology.

Recombination as a pumping mechanism was proposed by Gudzenko and Shelepin (1964) who considered inversions in hydrogen. Other work along these lines is referenced in the bibliography. Although transitions in hydrogen are not X-ray transitions, this effort was the precursor of approaches presently being examined involving multiply ionized systems. The problem in hydrogen is still under consideration (Furukane and Yakota (1979)).

Recombination in hydrogenic ions has been a very active area since early on. Gudzenko et al (1970) examines the $3-2$ transition of hydrogenic helium and gives references to earlier work. Bohn (1974) reports a steady state analysis of $2-3$ inversions in hydrogenic ions $1 \leq Z \leq 8$, and Pert and Ransden (1974) examine the same transition for low Z ions in an expanding laser plasma.

Experimental results were reported by Irons and Peacock (1974) in the case of hydrogenic carbon. Hydrogenic carbon has been a major focus of effort over the years and almost constitutes an area of study in itself (we have reserved a section in the references for this problem). The result given in the above mentioned paper is that the gain on the 3-2 transition in C VI a few mm from an irradiated polyethylene target was $2 \times 10^{-5} \text{ cm}^{-1}$. A numerical study by Pert (1976) (following an extensive earlier work) indicated that gain of about 10 cm^{-1} should be attainable in an expanding cylindrical plasma. Dewhurst et al (1976) reported spectrographic observation of hydrogenic carbon spectra taken from a laser heated thin (5.3μ) carbon fiber from which a gain-length product of order 10^{-2} was inferred through comparison with a numerical model. Tallents (1977) reports extensive calculations of inversions for hydrogenic carbon assuming exponential plasma cooling rates. This work was criticized by Jones and Ali (1978) on several ground (Jones and Ali (1975) and (1977) had previously analyzed the hydrogenic recombination scheme for low Z) and Tallents (1978) responded.

More experimental results were reported by Key et al (1979) who irradiated a carbon slab and monitored spectrographically the expanding plume. From line ratio data a gain of 0.2 cm^{-1} was inferred on the 3-2 line of the hydrogenic ion. Pert (1979) reported further numerical examination of the expanding carbon fiber problem, optimizing fiber radius and laser prepulse energy and reporting scaling laws. Gain of several cm^{-1} were calculated. Jacoby et al (1980) have

reported direct observation of a gain length product in the range 1-5 for 2 mm long thin carbon fibers on the 3-2 transition. In this experiment carbon fibers of diameters 2-5 μ were irradiated with pulses of energies 5-8 J. Spectroscopic observations were made from axial and transverse positions relative to the carbon fiber. Gain length products were estimated from the different spectrographic images. The group reports that 10 gain lengths may be achieved straightforwardly by scaling up the experiment to a 1 cm long fiber and 40J of cylindrically focused laser energy, with which a direct demonstration of laser action will be obtained.

Other work on recombination in hydrogenic plasmas is given in the references. Most of the work reported is theoretical - the experimental results of the British groups will constitute some of the more exciting results in this part of the field if they can be reproduced.

The extension of recombination to other sequences is straightforward. Theoretical results for helium are given by Gudzenko et al (1974). The recombination in helium-like aluminum has been studied both theoretically and experimentally by Bhagavatula and Yaakobi (1977). Theoretical calculations were performed using a one dimensional hydrodynamics code coupled with a kinetics package. The experiment is qualitatively different from the British experiment in that a nearby Mg plate was included to help increase the cooling rate. The presence of a population inversion was inferred from the ratios of the $1s^2-1snp$ lines ($n=2,3,4,5$) and a gain of 10 cm^{-1} was computed for the 3^3D-4^3F array.

Inversions in lithium-like ions have also received attention. Kononov and Koshelev (1975) examine 4-3 transitions in lithium-like aluminum and calculate gains of $0.1-1.0 \text{ cm}^{-1}$. Lacour et al (1975) proposed a scheme involving metastable lithium-like ions. Kononov et al (1976) inferred population inversions of the 3-4 and 3-5 transitions in Al XI from spectrographic observations and deduce gains of $0.05 - 0.2 \text{ cm}^{-1}$. Gaponov et al (1978) quote a theoretical result of 25 cm^{-1} on a 3s-2p transition in lithium-like neon in a rapidly cooling plasma. Theoretical studies of other systems have been carried out by Cacciatore and Capitelli (1980) (and references therein) although this work involves lower energy transitions.

Photoionization was suggested as a pumping mechanism for short wavelength lasers by Duguay and Rentzepis (1967). Two examples were given, including a low energy (33 eV) scheme involving 2p ionization from Na I leading to a 3s-2p inversion in Na II, and a high energy (8 keV) scheme involving 1s photoionization from Cu I leading to a 2p-1s inversion in Cu II. The Na scheme was analysed by Jones and Ali (1974) and again by Duguay (1976). A gain of 0.45 cm^{-1} is estimated under conditions where 2% of the initial Na atoms undergo 2p photoionization. A modification of the scheme was suggested by McGuire (1975) in which the 3s electron is collisionally ionized and the K shell electron is photoionized. The resulting Auger decay leads to a 2s-2p inversion in the oxygen-like ion. The present author extended the scheme to obtain a 2s-2p inversion in fluorine-like neon through 2s ionization of Ne I (Hagelstein (1977) in the review section). This

scheme is analysed in the Chapters II and III of this thesis.

Some experimental work was reported by McGuire and Duguay (1977) on the sodium vapor laser although they were unable to measure gain. Experimental work is planned at LLL on the Ne II laser; however, results have not yet been obtained.

The K-shell scheme was examined by Stankevich (1970) who pointed out that the shift in $1s-2p$ transition energy with loss of outer shell electrons is less than the linewidth. Arrechi et al (1974) estimates threshold conditions for the approach for several elements, and Elton (1975) reported a quasi-steady state analysis of many Z. In the case of silicon a gain of about 150 cm^{-1} is estimated for a 1.75 keV transition at an atom density of $7 \times 10^{21} / \text{cm}^3$.

Time dependent results for sulfur were given in Axelrod (1976) with the result that gains in excess of 2000 cm^{-1} could be reached, however the inversion endured only a few femtoseconds. Hagelstein (1976) examined the K-shell ionization approach for low Z and found that inversions could be obtained with slower risetimes at low Z, in the case of boron a risetime of about half a picosecond is required. Multi-stage pumping schemes were suggested in which continued ionization of K-shell electrons lead to strong $2p-1s$ inversions in the helium-like and hydrogen-like ions of atoms with $Z = 5, 7, \text{ and } 9$. Hagelstein and Chapline (1976) discuss the multi-stage scheme and the use of free radicals as starting material for the laser scheme. Axelrod (1977) reports the analysis of a laser design implementing a variant of the multi-stage scheme in boron. Inversion densities of

order 10% are calculated on the 1s-2p transitions of the hydrogenic ion with accompanying gains approaching 10^3 cm^{-1} . The scheme was found to be power hungry, requiring swept pumping of at least 20 TW absorbed infrared laser pumping power.

Several types of schemes have been proposed where photoexcitation is involved, including schemes using X-ray excitation, schemes using optical excitation and schemes employing combinations of the two. Vinogradov et al (1975) suggested photoexcitation pumping using near resonant line pairs, listing several candidates among which were included the sodium-neon pair (Na I $2^1P - 1^1S$ and Ne IX $4^1P - 1^1S$) with $\Delta E = 0.45 \text{ eV}$ and the silicon aluminum pair (Si XII $2^1P - 1^1S$ and Al XII $3^1P - 1^1S$) with $\Delta E = 3.9 \text{ eV}$. Gain in excess of 10^3 cm^{-1} is estimated for the 3-2 transitions in helium-like Al XII. Norton and Peacock (1975) suggest the use of opacity broadening in optically thick lines to make up defects in resonances and consider target design using the 1s - 2p transitions in C VI to pump the $1s^2 - 1s4p$ transition in C V. The use of satellite lines for pumping was also suggested in order to obtain more resonant pairs.

Bhagavatula (1976) pursued the resonant line approach suggesting the use of 1s - 2p transitions in hydrogenic and helium-like ions of charge Z to pump 2-4 transitions in ions of charge 2Z. Several resonant pairs of this type were given as well as a static analysis of the inversions obtained. Experimental results on this approach are reported in Bhagavatula (1978) in the case of carbon and magnesium resonances. Targets composed of carbon next to magnesium were irradiated with a 200 ps pulse of energy 10J. Spectra were taken from the

expanding plasma and magnesium intensity inversions were observed from which it was concluded that a population inversion on the 4-3 transition was obtained. This work was extended in Bhagavatula (1980), in which a very complete discussion of many of the issues is given. Gains of $5-10 \text{ cm}^{-1}$ at $130 \text{ \AA}$ is estimated. Apruzese et al (1978) examined the 3-2 inversion in helium-like aluminum as pumped by helium-like silicon. In this case a cylindrical plasma composed of an aluminum core and a silicon exterior as might be obtained in a laser plasma or exploding wire discharge was examined. Gain in excess of 10^3 cm^{-1} was estimated.

In the present thesis this approach is applied to laser target design involving resonances from L-M transitions in higher Z ions, for example, oxygen-like Ni pumping the helium-like neon $1s^2-1s3p$ singlet transition.

Photoexcitation schemes involving optical photoexcitation pumping were discussed in Freund (1974), Mahr and Roeder (1974) and Gudzenko et al (1974), (1975). Freund considers one and two photon processes assuming a $2s-1s$ inversion in hydrogenic oxygen. The idea is that it may be easier to obtain a $1s-2s$ population inversion than a $1s-2p$ inversion due to the metastability of the $2s$ state. One may take a long time to collect the metastable ions and then produce gain through application of an optical field. Mahr and Roeder (1974) propose creating a $1s2s^1S-1s^2^1S$ inversion in helium-like ions (Li II is the primary example) and optically stimulating the $1s2s^1S-1s2p^1P$ transition to produce $1s^2^1S-1s2p^1P$ gain. Gudzenko et al (1974), (1975)

considers storage in the $1s2s^3S$ of metastable helium-like ions (C was given as an example) and the production of gain through optical stimulation of the $1s2s^3S-1s2p^1P$ and $1s2s^3S-1s2p^3P$ transitions.

Experimental production of a strong flux of ionized lithium atoms was reported by Vekhov et al (1975), who optically pumped the 2s-2p transitions and detected the $1s2p^1P_1-1s^2s^1S_0$ radiation. Mani et al (1976) discuss a proposal for implementing the metastables produced by photoionization) and calculated optically induced Raman gain. The analysis, however, had neglected the effect of autoionizing states as discussed in Hyman and Mani (1977). The new analysis predicted a gain of 1.2 cm^{-1} .

Another approach to the problem has been discussed by Csonka (1976), (1978). In this case optical excitation is used to promote an outer shell electron to an excited state, for example a 2s-2p transition in Li I. Next, narrow band photoionizing radiation removes a 1s electron resulting in a 1s-2p self-terminating inversion. One novel aspect of the scheme is that photoionizing radiation would be provided by a synchrotron source. The analysis of the Raman scheme under these pumping conditions has not yet been presented, but may prove to be interesting.

McGuire and Duguay (1977) have suggested schemes involving photoionization, photoexcitation and metastable storage states. For example, photoexcitation of magnesium to the $2s^2p^63s^1P$ state followed by 2p photoionization leads to $2s^2p^53s$ states, one of which is metastable against Auger decay. Gain may occur on the $2s^2p^63s^2S-2s^2p^53s^2P$

transitions, and optical pumping can help dump the lower laser state.

Harris (1980) has proposed a scheme involving a $1s-2p$ transition in neutral lithium. Electron collisional excitation of the $1s^2 2s$ state populates the $1s2s\ ^4P^o$ metastable state, and optical stimulation of the $2s-2p$ spin-flip transition leads to gain on the $1s^2 2p^2 P - 1s2p^2\ ^2P$ transitions if the lower laser state is dumped (Harris suggests optical photoionization). Some of the metastable states have been observed in a microwave heated Li plasma (Willison et al (1980)).

Ion-atom resonant charge exchange was proposed in Vinogradov and Sobel'man (1973) wherein the problem of exchange between neutral helium and hydrogenic ions (neon is given as an example) is discussed. Scully et al (1973) consider charge pickup when, for example bare helium ions pass through hydrogen. In this case the electron transfer populates the 2p state of He II resulting in a 1s-2p inversion. An analysis of this scheme is given in Louisell et al (1975) and a gain coefficient of 0.1 cm^{-1} is estimated. Anderson et al (1976) solve numerically the population rate equations and the photon transport equation for the $584 \text{ \AA}$ radiation for the charge exchange of bare helium ions incident on cesium vapor target. A gain of 1.7 cm^{-1} is estimated. Charge exchange between Xe and hydrogenic helium and lithium leading to 3-2 inversions is discussed by McKnight and Seely (1977). Gains of 40 cm^{-1} for helium ($1640 \text{ \AA}$) and 6 cm^{-1} for lithium ($729 \text{ \AA}$) are calculated.

Charge transfer between neutral carbon and highly stripped carbon explains the spectra observed by Dixon and Elton (1977). In this experiment a carbon slab was irradiated with a 5J, 16 nsec pulse (intensity of $2 \times 10^{11} \text{ W/cm}^2$) and plasma expansion into a gaseous atmosphere was monitored. Charge exchange between neutral carbon and C VI and C V ions preferentially populates the $n = 4$ levels and enhanced emission of the 1-4 transitions relative to the 1-3 transitions was observed. Sizable 5-3 and 4-3 inversions are inferred. This work was extended in Dixon et al (1978) where further experimental spectral results are compared with numerical estimates of the

1s-np transitions of C VI, and charge exchange is shown to be necessary in explaining the observed spectra. The use of a channel to guide laser produced plasmas in order to create an elongated plasma when the laser is focused to a spot is discussed in Reintjes et al (1978).

Fukuda and Suemitsu (1977) observed 2s-2p transitions of C-, N- and O- impurity ions in a linear Z-pinch and found increased emission on some of the lines, possibly indicating the presence of charge exchange (from He I) induced population inversions.

Charge transfer onto protons is being studied by a group at Cornell (Haueisen et al (1978) and Dutfa et al (1979)). Tkach et al (1980) have reported observation of gain on the 1s-2p transition of H I. In this experiment protons in a plasma gun pulse penetrate a Na cloud at the 2p-1s signal is monitored. An 80° reflecting curved mirror reflects the 2p-1s radiation back through the charge exchange region, and is also monitored providing a direct measurement of the transition gain. A peak gain of 1.4 cm^{-1} was reported.

Other approaches and experiments are discussed in the reviews. Proposals for higher energy lasers involving nuclear transitions or exotic physical processes have appeared in the literature (see the graser section of the references); however, aside from providing pointers to this work we shall not consider the field further.

The use of mirrors in x-ray lasers would greatly decrease the gain requirements of many of the schemes discussed. At low energies, non-negligible reflectivities are observed in some metals (see Hunter et al

(1979) for the case of platinum as an example). Near 100eV the bulk reflectivities are dismally small ($\leq 1\%$) for normal incidence reflection. One approach to the problem is to use soap crystals (see Charles (1971), Toporkov and Sulabe (1976) and Henke et al (1978)) whose wavelength is matched to the frequency of the laser transition. Normal incidence reflection of lead stearate near $100\text{\AA}$ is about 10%. Another approach is to construct artificial multilayered structures which behave as lossy dielectric mirrors in the soft x-ray regime, as in for example, Dinklage (1967). Spiller (1972) discusses lossy mirror design in the UV (see also Spiller (1973), (1974), (1976)). Huelbich and Kunz (1976) report experimental results in the case of multilayer Au/C and Cu/C mirrors which were designed to maximize reflectance at 15° (90° is normal). Reflectivities between 5% and 10% were measured for the Au/C mirror between 60° and 15° near 100 eV. Vinogradov and Zel'dovich (1977) present theoretical results for multilayer mirrors and find theoretical normal incidence reflectivities which are quite high (35% on a Au/C mirror at 100 eV). Haelbich et al (1979) present experimental results for ReW/C multilayer films and find near 8% reflectivity at 15° near $200\text{\AA}$. Their earlier results for reflectivities involving the Au/C mirror were less than the theoretical results and the cause was apparently void regions in the gold layers. ReW films were found to be smoother.

B. Scope of Present Research

The short wavelength design effort which is discussed here is part of a larger and at present loosely coordinated effort at LLL to develop classes of short wavelength lasers. As was discussed in the previous section there are many approaches to short wavelength lasers, and interest locally is focused on schemes involving three types of sources, including high powered IR lasers (Argus, Shiva, and ICF class lasers elsewhere), Z-pinch discharges (such as Python at Physics International, a smaller version in E Division, and discharges elsewhere) and a pulsed nuclear source. The area of concern in this thesis includes short wavelength lasers pumped by laser fusion driver lasers.

The modelling effort which the author has been involved in follows on and is in some sense a continuation of some of the computational work of J. Weisheit and T. Axelrod. In the theoretical effort locally several types of schemes had been considered including rapid cooling and photoionization pumped schemes. In the literature, many more have been proposed. It is not immediately obvious which approach is best given a pump source, and theoretical efforts had been treating them one by one. It would be nice to have one code which is general which can evaluate classes of schemes and designs, and that was in part the goal of the work of the present author. The code in question is XRASER which is briefly described here and whose documentation may appear in the near future. Ultimately, many of the features or ideas of the code will be implemented in other codes around the laboratory (for example, LASNEX); however, for development purposes and in the

short term has been profitable to work with a smaller home-brewed code. (One can sometimes debug one's own code, but work with someone else's is much more difficult).

XRASER began simply as a rate equation code with some crude atomic physics recipes. It became clear that the field of atomic physics of ionized systems was expanding very rapidly and that it would be nice to include atomic physics results in the models as they became available. The format of the models became free form, namely, one could define a model in terms of its energy levels and rate coefficients and feed it to the rate equation solver. The idea was to have one file for each model and assemble a library of models with contributions from many sources and many individuals. Ultimately the models would become sufficiently complete that they could be used like, and as a supplement to, EOS tables at the laboratory. Not only would such a set of models be useful for short wavelength laser design, but for a wide range of plasma applications as well. Such an effort is of course not the work of one individual and interest has accordingly been narrowly focused to models directly applicable to the X-ray laser problem.

In order to provide atomic data for the models there are several approaches, including calculating it, looking it up in the literature or buying it. Calculating it in principle is possible given that one has the codes or has access to an atomic physics group and a fair sized budget, none of which was easily available. The route taken was to collect as much data as possible from the literature and to

restrict calculations to only the simplest of models. This approach has proven reasonable so far, and at present atomic physics parameters are drawn from a data base of several thousand papers. Code development has included simple atomic structure programs, photoionization and collisional cross section programs. In each of these areas massive work has been done elsewhere and continues to be done, and in no way can the part time effort of one graduate student approach state of the art levels in the different areas, and such was not attempted.

Similarly in the area of radiation transport, the field is very active and progresses on a month to month basis. The radiation transport in the present effort is not state of the art nor is it intended to be, instead it consists of algorithms which were simple, intuitively appealing and easy to program and debug. The same applies to the hydrodynamics.

The situation concerning experiments is simply that they are very expensive, extremely difficult to arrange and time consuming in terms of preparation and planning. At least six years of theoretical effort has been devoted locally toward the laser pumped approach and to date one related experiment has been attempted. The approach decided upon several years ago was to select one scheme with high gain and a high likelihood of success when tested experimentally. The experimental approach elsewhere has been to work the problem incrementally, namely, first producing a plasma, second measuring spectra from which the presence of the pumping mechanism can be inferred, third measuring directly gain of the order of unity and fourth,

building a laser. Such an option does not appear available in the present effort locally. The resources currently available have thus far been insufficient to successfully complete an experiment to measure the spectra produced from the flashlamp and filter system of the 27 eV laser design. In an environment where experiments are more precious by far than gold, one is forced to rely on theoretical results to an extent not usually encountered elsewhere.

The focus of the research devoted to the laser pumped approach recently has been on continued study of the 27 eV laser and on the development of a class of line pumped lasers at higher energies. The use of an additional optical laser directed axially down the length of the laser to assist in the inversion kinetics or to stimulate Raman gain would certainly lead to higher gain; however, the extra complexity and trouble which such an addition would make to an already complex experiment basically precludes it as an option, at least on SHIVA or ARGUS. Later on, if X-ray laser studies of this type become of interest, then use of an optical laser may be experimentally viable. For the present work we have not seriously considered designs requiring an optical laser. Similarly, mirrors would greatly improve matters; however, at present the expertise required to develop and implement a lossy mirror into a real target is not available.

The present design effort is therefore focused on laser targets which although not optimal are hopefully simple enough that they will require minimal effort on the part of the laser program. The

designs presented here do not represent finalized target specifications, since improvements in the physics and models which are constantly occurring as well as feedback from target fabrication and diagnostics groups lead to corresponding design modifications. With this as background we are ready to consider the next topic, namely an overview of the present document.

C. Thesis Overview

The design of a short wavelength laser pumped by a high intensity infrared laser is a complicated and multi-disciplinary venture. In many respects the problem is as challenging as that of inertial confinement pellet design. In both areas one faces the issues of laser plasma interaction and hot electron generation, hydrodynamic and shock physics of fluid models, equation of state and non-LTE physics, suprathermal electron and x-ray transport, and material opacity physics. Although short wavelength laser design does not involve neutron, gamma or charged particle transport, or burn physics, in their place one encounters resonantly trapped line radiation transport and laser light amplification problems. In fusion simulations two dimensional calculations are required and carried out on a routine basis. Similarly in the case of short wavelength laser systems two dimensional effects play important roles and a three dimensional calculation of the laser radiation transport appears to be necessary to predict the signal detected by an observing spectrograph.

One does not find a document in the area of inertial confinement target design by a single individual encompassing work in all the various disciplines involved as the problem is simply too large. Originally the present document had been intended to document the design procedure in short wavelength laser design and give a detailed account of all the issues involved (such intent originated back

when short wavelength laser design was simple). The present document was written and rewritten over a three year time span during which time the state of the art has advanced tremendously. The goals of the present document are severalfold, including:

1. Documentation of ideas and calculations relating to the target design implementing a photoionization scheme involving inversions in Ne II.
2. Documentation of ideas and calculations relating to target designs implementing line-pumping schemes.
3. Brief exposition of most of the major issues relating to laser-pumped target design.
4. Discussion of the design tools which have been used to date.
5. Proposal towards the commencement of experiments at LLL and elsewhere.

In the photoionization pumping scheme, X-rays near 500 eV preferentially remove 2s electrons from Ne I leading to a $2sp^6 2s-2s^2p^5 2p$ inversion near 27 eV. The photoionization cross sections for the different subshells of neon are shown in Figure I-1 and the scheme is illustrated in Figure I-2. This scheme is closely related to Duguay's photoionization scheme involving sodium (Duguay and Rentzepis (1967)), and has the advantages that neon gas is the laser medium rather than sodium vapor (leading to a simpler experiment) and that the collisional quenching of the inversion is greatly reduced due to the smaller collisional rates involved. The end result is a scheme which lends itself readily to an implementation in a simple target design. In Figure I-3 is shown a possible SHIVA

Figure I-1. Single electron photoionization cross sections for neon.

These cross sections were determined through photoelectron spectrometry by Francoise Wuilleumier (1973). The uncertainties in the cross sections were given to be 3%, 7% and 5% for the 1s, 2s, and 2p cross sections, respectively. One observes that X-rays in the 300 eV-800 eV region causes 2s ionization more often than 2p photoionization, the effect becoming greater at higher photon energies. It is this effect on which the Ne II laser scheme discussed in Chapters II and III is based.

Figure I-2. Ne II photoionization-pumped lasing scheme.

This lasing scheme is a modification of the photoionization-pumped lasing scheme of Duguay (1967) which involved the removal of a 2p electron from Na I. In the present case, selective photoionization of a 2s electron from Ne I leads to a transient population inversion in Ne II if the pumping occurs on a timescale shorter than the radiative decay time. The X-rays causing 2s photoionization to create the inversion also depletes the lasing transition states, which is an important effect if the pumping rates are high enough. As a practical matter it is difficult to produce sufficient X-ray fluxes in a pumping window near 500 eV to cause photoionization in 10% of the Ne I atoms.

Figure I-3 SHIVA Beam Positioning For Short Wavelength Laser Targets

The SHIVA laser facility was constructed for inertial confinement fusion experiments at the Lawrence Livermore Laboratory in California. Twenty arms of high-power Nd glass amplifiers provide multi-terawatt laser pulses to fusion pellet targets in order to study implosion and burn characteristics. The overall research effort is devoted towards the development of "inertial confinement fusion (ICF) technology for both military applications, in the near term, and civilian energy production, in the long term." (Emmett, Laser Program Ann. Rpt.-1979) High-power laser pulses have been used elsewhere (for example, NRL, U. of Rochester, Rutherford Labs, Sandia Labs, Hull University and the Lebedev Institute) in EUV and X-ray laser studies. In the present work we examine EUV and soft X-ray laser schemes which could be tested on SHIVA or on systems elsewhere. Shown here (schematically) is a possible beam configuration wherein 10 beams are focused one next to another along the length of the laser target on both top and bottom. No line focus capability is available or is planned. Due to the design of the lens mounts the maximum length which can be irradiated is 4 mm.

beam configuration on a short wavelength laser target. In the case of the Ne II laser low pressure neon gas flows through a brass channel, sandwiched from above and below by thin films which serve as flashlamps and filters upon irradiation.

In Chapter II we consider some of the atomic processes occurring in neon and its ions, including photoionization and photoexcitation processes, electron collisional processes and Auger and shake off rates. As the inversion mechanism is photoionization one is interested in the accuracy of the photoionization cross sections for the 2s and 2p sublevels, and as may be expected the cross sections are known to within a few percent from photoelectron spectrometry experiments. The dominant mechanism causing destruction of the inversions is collisional ionization of 2p electrons from Ne I. Perusal of the theoretical and experimental results available in the literature shows that 2p electrons are ionized much more easily than 2s electrons and that the relevant cross sections are known accurately from electron beam experiments.

Early calculations suggested that inversions could be obtained not only in Ne II but in Ne III and Ne IV if sufficient X-ray pumping flux could be made available. For this reason the neon model was extended to include atomic models for all of the low-lying 2s and 2p states through the lithium-like ion. Indeed, if one could arrange for a short (25 psec rather than 90 psec pulse which is currently available on SHIVA) and intense pulse and use specially designed

targets then inversions should be obtainable in the more highly stripped ions. Later on in the discussions of the line pumping schemes we will be stripping neon considerably further and some of the atomic data is applicable there.

Chapter III is devoted to target design issues for the Ne II laser. The basic problem is how to arrange for a sufficiently intense flux of pumping X-rays near 500 eV while filtering out the unwanted low energy spectrum. Molybdenum was selected for flashlamp material by virtue of its strong M-N line radiation which is emitted when heated. LiF was selected as filter material due primarily to its low absorption near 500 eV relative to the high absorption near 50 eV. The time-, energy-, and angle- dependent spectrum was calculated using LASNEX in conjunction with a post processor and a conversion efficiency of a few percent from the infrared pulse to pumping X-rays is estimated. No experimental confirmation of this result is yet available. The suprathermal electrons generated near the critical surface are assumed to wander into the laser medium and displace the thermal electrons, leading to a multi-kilovolt ambient electron temperature. Under this assumption (and others) the rate equations describing the population kinetics are integrated and the resulting inversion densities and gains are discussed. Gains in excess of 50 cm^{-1} are calculated.

Chapter IV is concerned with target design in the use of line pumping schemes. In this case the transition energies can be higher (calculations on the photoionization approach in silicon and argon suggest that the approach does not easily scale to higher energies; the 4-3 transitions in Ne IX are near 55 eV and the 3-2 are near 150 eV, which is considerably higher than the 27 eV of the Ne II scheme); however, the design issues are much more complex. Neon was chosen as an example and we consider line pumping schemes in helium-like and hydrogenic neon (Figure I-4 illustrates the 55 eV laser scheme).

One major problem encountered in resonantly-pumped laser schemes is finding sufficiently precise resonances. For example, line widths are near 0.1 eV and line energies are close to 1000 eV, hence, resonances should be accurate to within a few parts in 10^4 . Previously workers have examined K-shell/K-shell resonances and Bhagavatula (1976) has suggested 1-2 resonances with 2-4 transitions. In the case of Ne IX, the $1s^2-1s4p\ ^1P$ transition can be resonantly pumped with the Na X $1s^2-1s2p\ ^1P$ line (leading to a K-shell/K-shell resonance); however, this is a unique case for low Z elements. In general there are not precise resonances available of this type. The use of L-shell lines to provide resonances is explored and one finds that many resonances are available from classified lines.

Stripping the neon plasma to the one or two electron sequence requires several dozens of picoseconds and we discuss physical mechanisms through which this may occur. At the electron density

Figure I-4 Resonantly-pumped Photoexcitation Scheme Involving Helium-like Ne IX

In this scheme Ne IX is resonantly pumped to populate the N-shell, leading to a population inversion as the M-shell states can decay rapidly by radiative decay. Under the conditions of present interest, electron collisional processes distribute population between the N-shell singlet sublevels. Transitions which are calculated to attain usefully high gain include the $1s3s \ ^1S$ - $1s4p \ ^1P$, $1s3p \ ^1P$ - $1s4d \ ^1D$ and $1s3d \ ^1D$ - $1s4f \ ^1F$ transitions occurring between 53 and 55 eV. Resonant radiation may be provided by the helium-like Na X $1s^2$ - $1s2p \ ^1P$ resonance line or by lines in other systems which through chance occurrence are resonant to within a few parts in 10^4 . This scheme was published by Vinogradov et al (1975) and has received little attention since then. Pumping the M-shell rather than the N-shell may lead to a substantial $1s2p \ ^1P$ - $1s3d \ ^1D$ inversion at 151 eV, and in Chapter IV we examine both schemes.

chosen ($2 \times 10^{18}/\text{cm}^3$) for the Ne IX 4-3 scheme both photo- and electron collisional processes are important in the transient ionization process.

The transfer of line radiation in and out of the laser medium is an important issue in the 1-4 line pumped schemes. The problem is in getting the pumping radiation into the laser medium. In the 1-3 line pumped schemes the problem is in getting the trapped 1s-2p radiation out of the laser medium. One solution to this is to set up a graded ion velocity flow within the laser medium. This approach is used in recombination schemes where significant trapping would be lethal to the population inversion, and its use in photexcitation-pumped schemes has not yet been explored. The use of suprathermal electrons in conjunction with shadows to provide the requisite momentum source for a strongly graded fluid velocity flow is explored.

We examine the kinetics and explore conditions under which gains in excess of 50 cm^{-1} can be obtained. Such high gains are required for SHIVA since irradiation is limited to 4 mm and the development and implementation of mirrors given present resources is impossible. The requirements on gain translates roughly into a requirement on local infrared laser pump intensity, the idea being that the harder one drives the laser target the higher the gain can be since there will be more photons to resonantly pump an inversion. Within limits calculations suggest that this is true. Currently the amount of line radiation produced by the candidate pump lines is

uncertain, and the constraint on the gain translates into a constraint on the output of the pumping line--that is, if theory or experiment can demonstrate that in excess of 0.015 photons/mode of resonant pumping X-rays can be produced, then it is likely that the designs can be made to work on SHIVA. There is another problem which has not been examined in much detail, and that has to do with a factor of four or five variation in focused beam intensity present in the output of lasers at all of the large inertial confinement fusion laser facilities. Modulo these issues, it is the general conclusion of this research that short wavelength lasers of the variety discussed above and within can be made to work in the laboratory.

In the final chapter we summarize the major results and conclusions of the thesis and indicate areas where future efforts might profitably be spent.

The large number of references given reflects the literature intensive nature of the research as it has progressed so far. It may be possible to give an account of the work with a very minimal amount of bibliographic material; in any event the addition of it will certainly not hinder a worker in the field who is interested in pursuing some area of research.

This thesis is intended in part to give a general discussion of the physics involved in the design of short wavelength laser systems, and this implies that parts of the thesis are descriptive of work which does not represent a scientific contribution to a

field of research made by this author. For example, in Chapter III we discuss briefly what happens when intense infrared laser light is incident on a thin film flashlamp and filter system, and present the results of a calculation. The physics of laser plasma interaction is the scientific property of the laser fusion community, and the results and physics involved were obtained earlier by theorists and experimentalists in that field. The description given here is intended only to help this work to be accessible to a wider audience, some of whom are not familiar with what happens to matter as it is blasted into oblivion at 10^{14} W/cm². The calculations of the shock hydrodynamics and other disk properties when irradiated were carried out using LASNEX, which is currently supported by G. Zimmerman and his group in the laser fusion theory group locally. The contribution of the author in this area is in running the code (with aid from G. Zimmerman and others) and processing the data found in the dump files produced to arrive at the spectral data used elsewhere in the kinetics studies. The second chapter discusses some aspects of neon atomic physics and a neon kinetics model, and the intent here is to give a discussion of some of the physics involved in the Ne II laser kinetics as well as to describe what our models look like. The contributions made to the field of atomic physics are minimal (although there may be some; for example the problem of the Ne II laser is one of the first in which the shake-off cross section for two 2s electrons could be of practical interest, and will hopefully

prompt further theoretical work); however, readers not familiar with the general magnitude and energy dependences of the various cross sections involved will hopefully benefit from the discussion given.

It is of interest here to summarize what some of the major contributions of this research are. The primary contribution is the development of a general non-LTE and X-ray laser code, XRASER. The descriptions of what is in it are brief here, and further documentation may appear elsewhere. Most of the kinetics calculations in this document were carried out using XRASER, and in the course of the research the code has played a key role in analyzing different schemes and targets. The development of XRASER is largely the work of the present author, with some technical support kindly provided by B. Wilson. The Ne II laser scheme represents a contribution to the field of short wavelength lasers and was the brainchild of the present author. The scheme is related to the one described by Duguay (1967); however, the implementation of the self-terminating photoionization pumping approach in neon rather than sodium gives the scheme much higher gain. The line pumping schemes analyzed in Chapter IV are not original in the sense that Vinogradov (1975) described the approach earlier; however, the systematic exploration of L-M transitions of L-shell ions in the $Z=20$ to $Z=32$ range for resonances, though obvious, does not appear elsewhere. Many of the resonances suggested in Chapter IV have not appeared previously in the X-ray laser literature.

The use of a cylindrical laser target where a gas is surrounded by a thin film flashlamp and filter system has been described previously (see Axelrod (1977)). Here we examine two-dimensional targets where shadows allow coupling of the suprathermal momentum into the ion momentum via ionization gradients to establish very large velocity gradients. This type of target represents a contribution to the field of short wavelength laser design, as this approach has not been previously described. The idea itself was sparked by comments made long ago regarding differences between one-dimensional and two-dimensional effects by J. Nuckolls, and aside from such inspiration the idea is original. The existence and effects of the suprathermal electron driven shock wave which penetrates into the laser medium is an idea which does not appear in the X-ray laser literature (the author is unaware of literature on it in the laser fusion literature, although it must exist as it would play an important role in vacuum insulation target designs for laser fusion). The penetration of the laser medium by suprathermal electrons is a conclusion resulting from this research and does not appear elsewhere. The use of suprathermal electrons to pump 3s-3p inversions in Krypton follows immediately from trying to implement the scheme in a target similar to the Ne II target, and this was an idea suggested by the author. Even now the idea does not appear in the literature. Earlier, this approach constituted an additional chapter to this already too weighty document, and after having been deleted, all that remains is a brief paragraph in the conclusions.

Many of the contributions made are of a rather technical nature, as about five years of effort have been put in towards the development of modeling tools and towards the development and analysis of several target designs. This work has been carried out almost completely by the author. Much progress has been made towards the development of a laser-pumped X-ray laser, although upon perusal of Chapter IV, the reader will notice that there is yet work remaining, especially with regard to the analysis of pump radiation produced by the flashlamp. There is much room for contributions by others in this area, and it is the hope of this author that further work will lead to the rapid development of laser-pumped short wavelength lasers in the near future.

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II. Neon Atomic Processes

In this chapter we examine some of the physics involved in the inversion kinetics of 2s-2p self-terminating laser transitions in Ne II and more highly stripped ions. The scheme was discussed briefly in Chapter I. The basic idea is to use photoionizing X-rays in the vicinity of 500 eV to selectively remove 2s electrons from neon and its ions, leading directly to 2s-2p inversions between 15 eV and 35 eV. In the designs discussed in the next chapter we find that it is difficult to pump a usefully large inversion in ions other than Ne II; however, this has to do with the specifics of a Mo/LiF target when irradiated by a relatively long (90 psec) 1.06μ infrared laser pulse. If the pulse were shorter and more intense, then there begins to be the possibility of pumping inversions in the more highly stripped sequences. Also, if shorter wavelength light were used for the irradiation, the suprathermal electron contribution to the heating of the inner wall would be reduced, and the target would be able to produce more 500 eV X-rays before the thermal radiation of the inner wall caused excessive 2p ionization. Although the next chapter is devoted largely to the Ne II problem, this by no means indicates that the general problem of 2s-2p inversions in the more highly stripped sequences is not of interest. For this reason we have examined the physics of the L-shell states of the first eight ionization stages here. Some of the physical processes discussed here are applicable to the analysis given in

Chapter IV, wherein we consider neon stripped ultimately to the point of being bare of electrons in an examination of resonantly-pumped short wavelength laser schemes.

The most important process involved in the inversion scheme is photoionization, and the accuracy of the cross sections is of interest, especially for neutral neon. Fortunately very accurate experimental results are available in Ne I, and we examine the problem in section D. The second most important process is collisional ionization, and excitation followed by ionization, of $2p$ electrons, especially from Ne I. Collisional ionization of $2p$ electrons leads to population of the lower of potential laser states, and as the Ne II scheme is self-terminating, no recovery is possible after the lower laser state becomes significantly populated. The relevant rates involved in the case of Ne I are well-known from experimental work and we examine this issue in section E.

The gain of the self-terminating transitions is linearly dependent on the oscillator strength, and in section B we discuss the issue briefly. In the case of Ne II the experimental and theoretical results are in close accord and the oscillator strength is accurate to within 4%. In section A we discuss the problem of energy levels for a neon model. Higher order processes are considered in the final sections of the chapter. Significant among these are the large double photoionization cross sections for

radiation below the K edge (for which the double ionization cross section is in excess of 15% of the single ionization cross section) and the non-statistical branching ratio for Auger processes following removal of a K-shell electron, which can be used to pump a 2s-2p inversion in the oxygen-like ion.

One purpose of the chapter is to examine (mostly by example) what is involved in the construction of a kinetics model for the design code XRASER. Although it may not be of particular interest to the general reader unfamiliar with numerical simulation, there is a community locally to whom the present document is in part addressed who will be interested in using the code and developing kinetics models similar to the one described here. For this reason some of the tables used here were taken directly from a data file describing the neon model under discussion. The model includes configurations of the form $2s^m p^n$ in LSJ coupling of which there are 78. Over the course of the research several dozen models have been constructed for different systems, and many exist even for the Ne II kinetics problem. The main intent here was to give the interested reader some idea of what they looked like.

Non-LTE models of neon have been constructed by other workers, and some references are mentioned in Chapter IV. Low density neon kinetics is of interest in astrophysics; however, astrophysical problems are usually considered on much longer time-scales (the width of the pump laser pulse in the present context is 90 psec FWHM). Ionization balances which are of interest in astrophysics require careful

consideration of recombination processes for accurate calculation, and in the model described in this chapter we shall largely ignore the issue since recombination plays such an unimportant role in the transient inversion kinetics. Models which are closer in this regard to the one discussed here are reported in the short wavelength laser literature as listed in the references on photoionization- and photoexcitation-pumped laser schemes.

A. Energy Levels

The energy levels of the $2s^m p^n$ states of Ne I through Ne VIII have been tabulated by Fawcett (1975). Some errors appear in the work; for example, the $2p^4 \ ^1D$ and $\ ^1S$ states of Ne V are incorrect, and in these cases the energy levels of the same states in different elements were used to determine the required energy levels through least squares interpolation. The energy of the Ne III $2p^6 \ ^1S_0$ state was computed from the wavelength of the $2sp^5 \ ^1P_1 - 2p^6 \ ^1S_0$ transition which was recently observed (evidently for the first time) by Beyer et al. (1978). The ionization energies were taken from Bashkin and Stoner (1976). The energy levels, indices, degeneracies, and names are given in Table II-1. Woodyard and Altick (1975) have calculated Ne I energy levels and oscillator strengths. Although in this chapter we shall focus on a model including only L-shell states, we have in the course of the research examined more complex models and those results may be presented elsewhere. For present purposes we shall be content with a discussion of the model in Table I-1.

CM	ISO	i	NAME	G	ENERGY(11/CM)						
ENDT	2	HE-LIKE	1194.78								
CM	2	1	52-150	1.	0.						
						CM	5	5	P2-4P5/2	6.	101346.0
						CM	5	6	P2-203/2	4.	179020.6
						CM	5	7	P2-205/2	6.	178986.9
						CM	5	8	P2-251/2	2.	230853.0
						CM	5	9	P2-2P1/2	2.	249250.8
						CM	5	10	P2-2P3/2	4.	250112.5
ENDT	3	LI-LIKE	239.75								
CM	3	1	P0-251/2	2.	0.						
						CM	5	11	P3-453/2	4.	321575.4
						CM	5	12	P3-205/2	6.	359648.9
						CM	5	13	P3-203/2	4.	359711.9
						CM	5	14	P3-2P1/2	2.	405983.7
						CM	5	15	P3-2P3/2	4.	406051.1
CM	3	2	P1-2P1/2	2.	128151.9						
CM	3	3	P1-2P3/2	4.	129801.2						
ENDT	6	C-LIKE	125.38								
CM	6	1	P2-3P0	1.	0.						
						CM	6	2	P2-3P1	3.	412.4
						CM	6	3	P2-3P2	5.	1110.1
						CM	6	4	P2-102	5.	30291.5
						CM	6	5	P2-150	1.	63913.6
CM	6	6	P3-552	5.	88363.1						
CM	6	7	P3-303	7.	175832.6						
CM	6	8	P3-302	5.	175903.3						
CM	6	9	P3-301	3.	175926.9						
CM	6	10	P3-3P2	5.	208154.9						
CM	6	11	P3-3P1	3.	208156.6						
CM	6	12	P3-3P0	1.	208191.2						
CM	6	13	P3-102	5.	270561.8						
CM	6	14	P3-351	3.	279364.6						
CM	6	15	P3-1P1	3.	303818.9						
ENDT	5	B-LIKE	159.26								
CM	5	1	P1-2P1/2	2.	0.						
CM	5	2	P1-2P3/2	4.	1305.5						
						CM	6	16	P4-3P0	1.	413803.5
						CM	6	17	P4-3P1	3.	413460.6
						CM	6	18	P4-3P2	5.	412679.2
CM	5	3	P2-4P1/2	2.	100261.0						
CM	5	4	P2-4P3/2	4.	100705.0						
						CM	6	19	P4-102	5.	436549.1
						CM	6	20	P4-150	1.	500476.5

Table II-1. Neon energy levels

ENOT	7	N-LIKE	94.77		CM	9	3	P6-2S1/2	2.	217049.1
CM	7	1	P3-4S3/2	4.						0.
CM	7	2	P3-2D5/2	6.						41234.6
CM	7	3	P3-2D3/2	4.						41279.5
CM	7	4	P3-2P1/2	2.	ENOT	10	4L-LIKE		21.56	
CM	7	5	P3-2P3/2	4.	CM	10	1	P6-1S0	1.	0.
CM	7	6	P4-4P5/2	6.						183860.3
CM	7	7	P4-4P3/2	4.						184476.9
CM	7	8	P4-4P1/2	2.						184789.5
CM	7	9	P4-2D5/2	6.						254082.0
CM	7	10	P4-2D3/2	4.						254106.1
CM	7	11	P4-2S1/2	2.						299621.2
CM	7	12	P4-2P3/2	4.						320002.8
CM	7	13	P4-2P1/2	2.						320740.9
CM	7	14	P5-2P3/2	4.						484902.5
CM	7	15	P5-2P1/2	2.						485870.9
ENOT	8	O-LIKE	66.96							
CM	8	1	P4-3P2	5.						0.
CM	8	2	P4-3P1	3.						842.5
CM	8	3	P4-3P0	1.						920.5
CM	8	4	P4-1D2	5.						25840.8
CM	8	5	P4-1S0	1.						55748.4
CM	8	6	P5-3P2	5.						204289.6
CM	8	7	P5-3P1	3						204874.8
CM	8	8	P5-3P0	1.						205196.7
CM	8	9	P5-1P1	3.						289478.6
CM	8	10	P6-1S0	1.						478800.8
ENOT	9	F-LIKE	44.45							
CM	9	1	P5-2P3/2	4.						0.
CM	9	2	P5-2P1/2	2						780.6

Table II-1. Neon energy levels (continued)

Table II-1. Neon Energy Levels (legend)

In this table we have reproduced part of an XRASER data file which contains the neon model under discussion in the present chapter. There are cards describing the isoelectronic sequence (i.e., ENOT 2 HE-LIKE 1194.78, which describes the 2-electron sequence named HE-LIKE with ionization potential 1194.78 eV). Below these cards are ones describing the levels involved (i.e., CM 3 2 P1-2P1/2 2. 128151.9, which describes the second level in the 3-electron sequence, which in this case is $2p^1\ ^2P_{1/2}$, with statistical weight of 2 and an energy of 128151.9 cm^{-1} relative to the ground state). The purpose of this table is primarily to give the interested reader an idea of what is involved; for example, when the word model is used in this document, often it will refer to a data file part of which looks similar to the table given here. Over the course of the work we have studied several dozen models of different systems, and associated with each is a data file of similar construction.

The notation for a state is

$$(2S+1)_{L,J}$$

where S is the total spin, L is the angular momentum (given as a letter) and J is the total angular momentum. For example, the state $^2P_{1/2}$ has $S=1/2$, $L=1$ and $J=1/2$.

B. Wavelengths and Oscillator Strengths for 2s-2p Transitions

Perhaps the most important of the 2s-2p transitions for our purposes are the Ne II $2s^2p^5$ - $2sp^6$ transitions. The wavelengths are known accurately and have been given by Fawcett (1975) to be at $460.725\text{\AA}$ and $462.388\text{\AA}$ for the $^2p_{3/2}$ - $^2s_{1/2}$ and $^2p_{1/2}$ - $^2s_{1/2}$ transitions respectively. These wavelengths correspond to energies of 26.911 eV and 26.814 eV respectively, and because the transition energies are close to 27 eV, the laser resulting from the implementation of the scheme under discussion is sometimes referred to as the "27 eV laser." The oscillator strengths for the two lines probably have values close to one another (we deal with excitation oscillator strengths which are proportional to the statistical weight of the upper state--in this case the upper state is the same for both lines). In LS coupling the equality is precise, and it is spin orbit effects which would cause deviations from equality. Cheng et al (1979) gives results in j-j coupling for the two lines separately and finds about 1% difference between the two oscillator strength. Unfortunately the calculation is single configurational relativistic Hartree Fock for this case and the absolute values are not in agreement with other results. Westhaus and Sinanoglu (1969) quoted values of 0.073 and 0.091 for the length and velocity forms of the oscillator strength for the $2s^2p^5$ $2p$ - $2sp^6$ 2S transition (which should be the same for the two line in intermediate coupling modulo the few percent spin-orbit difference) using many body correlation theory. Beck and Nicolaides (1976) have reported a value

of 0.062 for a correlated perturbation theory result. Experimentally, the value has been quoted to be 0.070 by Curtiss and Smith (1974) and 0.076 by Knystautas and Drouin (1974). In the former case the accuracy is quoted to be $\pm 20\%$, and in the latter case to be $\pm 11\%$. Irwin et al (1973) have quoted a value of 0.077 ± 0.003 for their beam foil results.

In Table II-2 we give the wavelengths and oscillator strengths for 2s-2p transitions in the neon ions. Wavelengths were taken from Fawcett (1975), and where wavelengths were not tabulated, they were estimated from the differences in energy levels. Although precise values are not crucial for the calculations, they will be helpful in sorting out spectra in EUV laser experiments and from other experiments. Some of the line positions may be determined more accurately from experimental work on the laser. Oscillator strengths were taken from the literature, primarily from results of beam foil experiments and accurate atomic structure calculations. In the cases where an oscillator strength or a decay time was quoted for a set of lines, the individual transition oscillator strengths were determined from the tables given in Shore and Menzel (1965). The data marked "Kludge" was data estimated using the total decay rate of the upper state and using multiplet theory to get the individual ratios (see Shore and Menzel (1965)).

After the table was prepared, a compilation of 2s-2p oscillator strengths was published in Fawcett (1978). Comparison of the two

F-LIKE																													
9	1	9	3	0.070	460.725	P3/2-S1/2	KUR74			7	5	7	10	0.005	521.740	P3/2-D3/2	BEC76			6	4	6	13	0.181	416.198	D2-D2	HC178		
9	2	9	3	0.070	462.388	P1/2-S1/2	KUR74			7	4	7	11	0.077	421.609	P1/2-S1/2	KER77			6	4	6	15	0.135	385.594	D2-P1	HC178		
										7	5	7	11	0.077	421.616	P3/2-S1/2	KER77			6	5	6	15	0.245	416.034	S0-P1	LAU73		
G-LIKE										7	1	7	8	0.035	541.127	S3/2-P1/2	BEC76			6	7	6	18	0.012	422.214	D3-P2	KER72-KLU00E		
8	1	8	6	0.132	489.501	3P2-3P2	KER77			7	1	7	7	0.069	542.073	S3/2-P3/2	BEC76			6	8	5	17	0.009	420.951	D2-P1	KER72-KLU00E		
8	2	8	7	0.044	483.841	3P1-3P1	KER77			7	1	7	7	0.104	543.891	S3/2-P5/2	BEC76			6	9	5	16	0.020	420.386	D1-P0	KER72-KLU00E		
8	1	8	7	0.044	488.103	3P2-3P1	KER77			7	9	7	14	0.098	433.237	D5/2-P3/2	KER77			6	8	5	18	0.003	422.347	D2-P2	KER72-KLU00E		
8	2	8	6	0.074	481.050	3P1-3P2	KER77			7	10	7	15	0.081	431.472	D3/2-P1/2	KER77			6	9	5	17	0.005	420.983	D1-P1	KER72-KLU00E		
8	3	8	7	0.176	480.510	3P0-3P1	KER77			7	10	7	14	0.017	432.836	D3/2-P3/2	KER77			6	9	6	18	0.001	422.383	D1-P2	KER72-KLU00E		
8	2	8	8	0.058	488.898	3P1-3P0	KER77			7	12	7	14	0.183	606.429	P3/2-P3/2	KER77			6	10	6	18	0.013	488.594	P2-P2	KER72-KLU00E		
8	4	8	9	0.181	379.308	102-1P1	KER77			7	12	7	15	0.037	602.888	P3/2-P1/2	KER77			6	11	6	17	0.024	487.080	P1-P1	KER72-KLU00E		
8	5	8	9	0.058	427.840	150-1P1	KER77			7	13	7	15	0.147	605.595	P1/2-P1/2	KER77			6	11	6	16	0.008	486.270	P1-P0	KER72-KLU00E		
8	9	8	10	0.186	526.2	1P1-1S0	SAF75			7	13	7	14	0.073	609.168	P1/2-P3/2	KER77			6	10	6	17	0.004	487.070	P2-P1	KER72-KLU00E		
										7	11	7	14	0.497	539.731	S1/2-P3/2	SAF76			6	12	6	17	0.017	487.170	P0-P1	KER72-KLU00E		
										7	11	7	15	0.133	536.985	S1/2-P1/2	SAF76			6	11	6	18	0.007	488.594	P1-P2	KER72-KLU00E		
H-LIKE										6	14	6	16	0.002	743.830					5	1	5	6	0.078	558.585	P1/2-D3/2	1M73		
7	2	7	12	0.188	358.721	D2-P3/2	KER77			6	14	6	17	0.005	745.730					5	2	5	7	0.089	562.805	P3/2-D5/2	1M73		
7	3	7	13	0.138	357.831	D3/2-P1/2	KER77			6	14	6	18	0.008	750.100					5	2	5	6	0.013	562.735	P3/2-D3/2	1M73		
7	3	7	12	0.029	358.778	D3/2-P3/2	KER77			C-LIKE										6	13	6	19	0.238	802.480	D2-D2	1M73-KLU00E		
7	4	7	13	0.070	387.141	P1/2-P1/2	KER77			6	3	6	7	0.084	572.336	P2-D3	HC178			6	15	6	19	0.108	753.410	P1-D2	1M73-KLU00E		
7	5	7	13	0.018	387.147	P3/2-P1/2	KER77			6	2	6	8	0.057	588.830	P1-D2	HC178			6	16	6	20	0.048	508.500	P1-S0	1M73		
7	4	7	12	0.034	388.247	P1/2-P3/2	KER77			6	1	6	9	0.078	588.418	P0-D1	HC178			6	15	6	19	0.108	753.410	P1-D2	1M73-KLU00E		
7	5	7	12	0.088	388.218	P3/2-P3/2	KER77			6	3	6	8	0.011	572.106	P2-D2	HC178			6	15	6	20	0.048	508.500	P1-S0	1M73		
7	4	7	12	0.034	388.247	P1/2-P3/2	KER77			6	2	6	9	0.019	589.759	P1-D1	HC178			6	15	6	20	0.048	508.500	P1-S0	1M73		
7	5	7	12	0.088	388.218	P3/2-P3/2	KER77			6	3	6	10	0.088	482.987	P2-P2	HC178			B-LIKE									
7	2	7	10	0.009	488.774	D5/2-D3/2	BEC76			6	2	6	12	0.031	481.281	P1-P0	HC178			5	1	5	6	0.078	558.585	P1/2-D3/2	1M73		
7	2	7	9	0.123	488.880	D5/2-D5/2	BEC76			6	2	6	10	0.038	481.385	P1-P2	HC178			5	2	5	7	0.089	562.805	P3/2-D5/2	1M73		
7	3	7	10	0.119	488.885	D3/2-D3/2	BEC76			6	1	6	11	0.082	480.406	P0-P1	HC178			5	2	5	6	0.013	562.735	P3/2-D3/2	1M73		
7	3	7	9	0.013	488.882	D3/2-D5/2	BEC76			6	3	6	11	0.025	482.983	P2-P1	HC178			5	2	5	6	0.013	562.735	P3/2-D3/2	1M73		
7	4	7	10	0.037	521.741	P1/2-D3/2	BEC76			6	3	6	14	0.106	359.358	P2-S1	HC178			5	1	5	10	0.079	399.820	P1/2-P3/2	KER77		
7	5	7	9	0.031	521.820	P3/2-D5/2	BEC76			6	2	6	14	0.106	358.472	P1-S1	HC178			5	1	5	9	0.158	401.138	P1/2-P1/2	KER77		
										6	1	6	14	0.136	357.955	P0-S1	HC178												

Table II-2. Neon oscillator strengths and wavelengths

5 2 5 10 0.198 401.939 P3/2-P3/2	KERT77	4 3 4 7 0.036 561.378	P1-P1	VIC77
5 2 5 9 0.040 403.262 P3/2-P1/2	KERT77	4 3 4 6 0.048 562.992	P1-P0	VIC77
		4 4 4 7 0.036 564.529	P2-P1	VIC77
5 1 5 8 0.043 433.176 P1/2-S1/2	KERT77			
5 2 5 8 0.043 435.649 P3/2-S1/2	KERT77	4 5 4 9 0.134 973.100	P1-D2	VIC77
		4 5 4 10 0.063 561.210	P1-S0	VIC77
5 5 5 11 0.095 454.072 P5/2-S3/2	KERT77			
5 4 5 11 0.095 452.745 P3/2-S3/2	KERT77			
5 3 5 11 0.095 451.843 P1/2-S3/2	KERT77			
5 7 5 15 0.056 440.404 D5/2-P3/2	KERT77	L1-LIKE 3 1 3 2 0.095 780.324 S1/2-P1/2 ARW76 3 1 3 3 0.110 770.409 S1/2-P3/2 ARW76		
5 6 5 14 0.047 440.600 D3/2-P1/2	KERT77			
5 6 5 15 0.009 440.469 D3/2-P3/2	KERT77			
5 7 5 13 0.007 553.330 D5/2-D3/2	DANT77			
5 7 5 12 0.098 553.520 D5/2-D5/2	DANT77			
5 6 5 13 0.090 553.430 D3/2-D3/2	DANT77			
5 6 5 12 0.011 553.620 D3/2-D5/2	DANT77			
5 10 5 12 0.054 912.940 P3/2-D5/2	1PW73			
5 9 5 13 0.065 905.620 P1/2-D3/2	1PW73			
5 10 5 13 0.011 912.410 P3/2-D3/2	1PW73			
5 10 5 14 0.022 841.560 P3/2-P1/2	DANT77			
5 10 5 15 0.108 841.280 P3/2-P3/2	DANT77			
5 9 5 14 0.087 838.180 P1/2-P1/2	DANT77			
5 9 5 15 0.043 837.920 P1/2-P3/2	DANT77			
5 8 5 15 0.070 570.760 S1/2-P3/2	DANT77			
5 8 5 14 0.035 571.000 S1/2-P1/2	DANT77			
DE-LIKE				
4 1 4 5 0.381 465.221	S0-P1	VIC77		
4 4 4 8 0.109 561.728	P2-P2	VIC77		
4 3 4 8 0.060 558.610	P1-P2	VIC77		
4 2 4 7 0.145 558.947	P0-P1	VIC77		

Table II-2. Neon oscillator strengths and wavelengths (continued)

Table II-2. Neon Oscillator Strengths and Wavelengths (legend)

In this table we have reproduced part of the XRASER data file containing oscillator strength data. The form of the data is

IS01 I1 IS02 I2 F λ name reference

where IS01 and I1 denote the lower level of the transition and IS02 and I2 denote the upper level, where the levels are defined in Table I-1. F is the excitation oscillator strength and λ is the wavelength in Angstroms. Each transition was given a name which is related to the L and J values involved, and the last column contains the line reference. The references are given by

ARM76:	Armstrong et al (1976)
BEC76:	Beck and Nicolaides (1976)
CUR74:	Curtiss and Smith (1974)
DAN77:	Dankwort and Treffetz (1977)
IRW73:	Irwin et al (1973)
KER72:	Kernahan et al (1972)
KER77:	Kernahan et al (1977)
LAU73:	Laughlin and Dalgarno (1973)
MCI78:	McIntyre et al (1978)
SAF75:	Safranov (1975)
SAF76:	Safranov (1976)
VIC77:	Victrova and Safranov (1977)

Where "Kludge" was used we have used multiplet theory as mentioned in the text.

tables shows fair agreement in most cases, and in cases where good experimental results have been used, the present table will be more accurate. Through oversight the Curtiss and Smith (1974) value was used in the calculations of Chapter III, rather than the more accurate Irwin (1973) result. The major difference in the final answer would be an increase in the gain from what is given by a factor of 10%.

C. Collisional Excitation for 2s-2p Transitions

The collisional excitation cross section for dipole allowed transitions at energies much greater than threshold is dominated by the dipole contribution. In the neon plasma, the electron temperature will be determined by the suprathermal component of the electron distribution, and the resulting effective temperature will be in excess of 1 keV during pumping. Under these conditions, the dipole allowed collisional transitions will be dominant and cross sections can be calculated from standard classical path theory (see Burgess (1976)). The formula used is

$$\sigma(E) = 8\pi a_0^2 \left(\frac{I_H}{E} \right) \left(\frac{I_H}{\Delta E} \right) f \exp(\pi\xi) (\xi\epsilon) K_{i\xi}(\xi\epsilon) K'_{i\xi}(\xi\epsilon) \quad (\text{II-C-1})$$

where $\sigma(E)$ is the cross section as a function of energy, a_0 is the Bohr radius (5.29×10^{-9} cm), I_H is the ionization potential of hydrogen (13.605 eV), ΔE is the transition energy and f is the oscillator strength. The parameter ξ is defined by

$$\xi = \frac{Z}{a_0} \left| \frac{1}{k_i} - \frac{1}{k_f} \right| \quad (\text{II-C-2})$$

where Z is the charge seen by the incident electron and k_i and k_f are the initial and final electron wavenumbers. The parameter ϵ is

$$\epsilon = 1 + \frac{R_c |k_i - k_f|}{\xi} \quad (\text{II-C-3})$$

where R is the distance of closest approach, taken to be the smaller of the radii of the initial and final states. The functions involved include a modified Bessel function

$$K_{i\xi}(x) = \int_0^\infty \exp(-x \cosh(t)) \cos(\xi t) dt \quad (\text{II-C-4})$$

and its first derivative

$$K'_{i\xi}(x) = \int_0^{\infty} \exp(-x \cosh(t)) \cosh(t) \cos(\xi t) dt \quad (\text{II-C-5})$$

The "weak coupling" version of the semi-classical model assumes that the colliding electron follows a classical hyperbolic path, and that the target atom or ion behaves quantum mechanically, responding to the dipole term in the multipole expansion of the Coulomb interaction. The probability that a transition is made for a given choice of impact parameter can be computed to be

$$P = 4a_0^2 \frac{I_H}{E} \frac{I_H}{\Delta E} f \quad (\text{II-C-6})$$

$$\times (\xi \epsilon)^2 e^{\pi \xi} \left\{ \left| K_i(\xi \epsilon) \right|^2 + \frac{\epsilon^2 - 1}{\epsilon^2} \left| K_{i\xi}(\xi \epsilon) \right|^2 \right\}$$

This expression can be integrated analytically over impact parameters to give the result in II.C.1 where the limits of integration are from some smallest impact parameter corresponding to the distance of closest approach being equal to the radius of the smaller state, to some largest impact parameter beyond which the contributions are either screened or are ineffective due to state lifetimes. Under the present conditions, the cutoff at large impact parameters is completely negligible and we have the cutoff extended to infinite impact parameter. In order for the theory to be valid, we must ensure that the probability is much less than unity over the range of integration. The largest values are obtained at the cutoff radius and for the calculations here in all cases it is less than 10%.

In Table II-3 is compared the results of the semi-classical calculations with quantal calculations. The close coupling method is the most accurate method normally used for this type of calculation, and one observes that in the example considered that the accuracy is comparable to that of the Coulomb-Born approximation.

TABLE II-3. COMPARISON OF SEMI-CLASSICAL RESULTS WITH OTHER RESULTS
FOR Ne VIII

<u>E(I_H)</u>	<u>E(eV)</u>	<u>Ω(CC)</u>	<u>Ω(CBII)</u>	<u>Ω(CBI)</u>	<u>Ω(C.P.)</u>
1.6	21.77	3.34		3.50	2.92
3.2	43.54	3.59		3.76	3.31
8.0	108.85	4.23		4.47	4.02
14.0	190.48	4.56	4.63	4.90	4.60
18.0	244.90	4.86	4.93	5.17	4.89
30.0	408.17	5.60	5.66	5.72	5.50

Ω(CC) Close coupling results of Van Wyngaarden and Henry (1976)
 Ω(CBII) Unitarized Coulomb-Born results of Bely (1966).
 Ω(CBI) Coulomb-Born (ununitarized) results of Bely (1966)
 Ω(C.P.) Present classical path results

TABLE II-3 (continued)

The cross section σ is related to the collision strength Ω by

$$\sigma = \frac{\pi a_0^2}{g_i} \left(\frac{I_H}{E} \right) \Omega$$

where a_0 is the Bohr radius, g_i is the statistical weight of the initial state, I_H is the ionization energy of Hydrogen, and E is the incident electron energy. The collision strength Ω is dimensionless and σ has dimensions of length squared.

In Figure II-1 is shown the cross section for the $^2P_{1/2} - ^2S_{1/2}$ transition and the corresponding rate coefficient obtained by averaging the cross section over a Maxwellian. The expression II.C.1 was symmetrized using the procedure described in Burgess (1976) for this calculation, and below 1.25 eV the cross section is assumed to vary inversely with energy as the expression II.C.1 when symmetrized gives unreasonable near-threshold results. In Table II-4 is presented an example of the calculation, showing the cross section, the symmetrized probability at the cutoff radius, the rate coefficients, and the error in the least squares fit of the rate coefficient. The rate coefficients have been parameterized by

$$\langle \sigma v \rangle = \frac{1.58 \times 10^{-8} P(\beta) e^{-\beta}}{\beta T_{\text{eV}}^{3/2}} \quad (\text{II-C-7})$$

where

$$\beta = \frac{\Delta E}{kT} \quad (\text{II-C-8})$$

Figure 11-1. Cross-section and rate coefficient for the Ne II $2p_{1/2} - 2s_{1/2}$ transition.

This cross section was computed in the classical path approximation. The approximation considers the perturbing electron to follow a hyperbolic path around the target ion, and the effect is calculated to first order assuming the dipole contribution dominates the interaction. Near threshold the formula describing the cross section becomes inaccurate due to the inadequacy of the theory in removing the energy imparted to the ion from the perturbing electron. Inverse energy dependence has been assumed near threshold. One observes a weak temperature dependence of the rate coefficient from 10 eV to 10^4 eV electron temperature, and that the magnitude of the rate coefficient is small: ($10^{-9} \text{ cm}^3/\text{sec}$ at $10^{17}/\text{cm}^3$ yields a rate of 10^8 sec^{-1} which is slow compared to the target irradiation time of 10^{-10} sec).

Z = 1.000E+00
DE = 2.691E+01
f = 7.000E-02
RHO1 = 9.610E-03

ENERGY(EV) SIGMA(CM**2) PROBABILITY AT CUTOFF

35.00	5.5600E-18	5.83E-02
40.00	5.3962E-18	4.95E-02
45.00	5.1370E-18	4.36E-02
50.00	4.9431E-18	3.91E-02
70.00	4.2663E-18	2.83E-02
100.00	3.6677E-18	2.00E-02
150.00	3.2666E-18	1.70E-02
150.00	2.6755E-18	1.38E-02
170.00	2.6678E-18	1.22E-02
200.00	2.4112E-18	1.04E-02
225.00	2.2367E-18	9.31E-03
250.00	2.0863E-18	8.41E-03
275.00	1.9602E-18	7.66E-03
300.00	1.8601E-18	7.04E-03
325.00	1.7810E-18	6.50E-03
350.00	1.6833E-18	6.05E-03
375.00	1.5899E-18	5.65E-03
400.00	1.5197E-18	5.30E-03
425.00	1.4581E-18	4.98E-03
450.00	1.3941E-18	4.70E-03
475.00	1.3411E-18	4.47E-03
500.00	1.2964E-18	4.25E-03
600.00	1.1313E-18	3.59E-03
700.00	1.0066E-18	3.01E-03
800.00	9.1243E-19	2.66E-03
1000.00	7.3694E-19	2.13E-03
1500.00	5.6029E-19	1.61E-03
2000.00	4.4661E-19	1.26E-03
3000.00	3.2057E-19	9.18E-04
4000.00	2.6978E-19	7.68E-04
5000.00	2.3908E-19	6.70E-04
7500.00	1.6118E-19	4.64E-04
10000.00	1.1600E-19	3.34E-04
20000.00	6.4118E-20	1.80E-04
30000.00	4.5168E-20	1.26E-04
40000.00	3.5698E-20	9.70E-05
50000.00	2.9877E-20	8.16E-05
70000.00	2.1901E-20	5.96E-05
100000.00	1.5681E-20	4.27E-05

TEMP(EV) SIGMA(CM**2)/SEC

20.00	8.0600E-10
30.00	1.2700E-09
50.00	1.7217E-09
70.00	1.8958E-09

This example shows the details of what is involved in the calculation of a rate coefficient for the neon model and for other models which have been constructed. The example is concerned with the cross section and rate coefficient for the $2s^2p^5\ ^2P_{1/2}-2s^2p^5\ ^2S_{1/2}$ transition of Ne II. The parameters required for the computation of the classical path cross section are given first, including the charge seen by the incident electron ($Z=1$), the transition energy ($\Delta E=26.91$ eV), the oscillator strength ($f=0.07$) and the cutoff radius ($r_0=0.761$ a₀) where a₀ is a Bohr radius. The cross section was evaluated at a set of incident electron

energies between 35 eV and 100 keV and listed is the probability at cutoff discussed in the text. Since the probability at cutoff is much less than 0.5, the weak coupling approximation may be accurate for this cross section. The rate coefficient was calculated by integrating the cross section over Maxwellians at a set of electron temperatures. The data was then least squares fit to Eq. (II-C-7) and the results given in the final line in a form which can be used in a datafile. The error in the fit is seen not to exceed 5% over the range given.

TEMP(EV)	SIGMA	SIGMA1	ERROR
20.00	8.0600E-10	8.3457E-10	3.24E-02
30.00	1.2700E-09	1.2447E-09	2.63E-02
50.00	1.7217E-09	1.6662E-09	3.05E-02
70.00	1.8958E-09	1.9006E-09	1.36E-02
100.00	1.9600E-09	1.9037E-09	7.11E-03
150.00	1.9678E-09	1.8444E-09	2.18E-02
200.00	1.8958E-09	1.8267E-09	2.27E-02
300.00	1.6118E-09	1.6436E-09	1.26E-02
500.00	1.1618E-09	1.0882E-09	3.01E-03
700.00	1.0117E-09	1.7006E-09	8.37E-03
1000.00	1.7000E-09	1.6662E-09	1.50E-02
2000.00	1.4600E-09	1.4758E-09	5.43E-03
3000.00	1.2958E-09	1.3406E-09	3.40E-03
5000.00	1.1758E-09	1.1600E-09	1.23E-02

DE 9 1 5 3 1.5977E+00 3.7086E-01 8.1314E-03 6.6270E-04 2.6910E+01 20.0 5000.0

Table II-4. Example of rate coefficient calculation and least squares fit.

and the function $P(\beta)$ is defined by

$$\ln P(\beta) = a_0 + a_1 \ln \beta + a_2 (\ln \beta)^2 + a_3 (\ln \beta)^3 \quad (\text{II-C-9})$$

This parameterization is adopted in analogy with the semi-empirical approach of Van Regemorter (1962), and has been found to be useful in fitting cross sections over several orders of magnitude in temperature to within a few percent accuracy. In Table II-5 are given results for all the 2s-2p transitions in the neon ions considered.

	1501	11	1502	12	A0	A1	A2	A3	DE	THIN	THAT
F-LINE											
9	1	9	3	-3.5697E-20	-4.2460E-01	-4.3599E-02	-3.1467E-03	2.8619E-01	20.0	5000.0	
9	2	9	3	-3.7616E-20	-4.1586E-01	-4.2034E-02	-3.2054E-03	2.8619E-01	20.0	5000.0	
G-LINE											
8	1	8	6	-2.7421E-08	-3.5215E-01	-2.7440E-02	-1.8829E-03	2.5329E-01	20.0	5000.0	
8	2	8	7	-3.0460E-08	-3.5237E-01	-2.7513E-02	-1.8817E-03	2.5321E-01	20.0	5000.0	
8	1	8	7	-3.8340E-08	-3.5089E-01	-2.8794E-02	-1.7388E-03	2.5961E-01	20.0	5000.0	
8	2	8	6	-3.3081E-08	-3.5049E-01	-2.8807E-02	-1.6977E-03	2.5940E-01	20.0	5000.0	
8	3	8	7	-2.4251E-08	-3.5339E-01	-2.7085E-02	-1.6642E-03	2.5207E-01	20.0	5000.0	
8	2	8	8	-3.8940E-08	-3.5116E-01	-2.7120E-02	-1.7089E-03	2.5817E-01	20.0	5000.0	
8	4	8	9	-2.7057E-08	-3.5773E-01	-1.8538E-02	-1.6088E-03	2.6857E-01	20.0	5000.0	
8	5	8	9	-3.7463E-08	-3.4129E-01	-2.8329E-02	-1.9530E-03	2.6670E-01	20.0	5000.0	
8	9	8	10	-2.3940E-08	-3.7759E-01	-3.8509E-02	-2.8523E-03	2.7473E-01	20.0	5000.0	
H-LINE											
7	2	7	12	-2.7617E-08	-5.3889E-01	-8.3429E-02	-7.1673E-03	3.4583E-01	20.0	5000.0	
7	3	7	13	-2.6797E-08	-5.3899E-01	-8.3863E-02	-7.8311E-03	3.4499E-01	20.0	5000.0	
7	3	7	12	-4.5496E-08	-5.3869E-01	-8.3476E-02	-7.1381E-03	3.4503E-01	20.0	5000.0	
7	4	7	13	-3.7188E-08	-6.1998E-01	-1.0818E-01	-1.0897E-02	3.2459E-01	20.0	5000.0	
7	5	7	13	-5.6793E-08	-6.1586E-01	-1.0811E-01	-1.0829E-02	3.2459E-01	20.0	5000.0	
7	4	7	12	-4.1442E-08	-5.4116E-01	-2.7476E-02	-1.0922E-03	3.1464E-01	20.0	5000.0	
7	5	7	12	-3.2163E-08	-5.4117E-01	-2.7490E-02	-1.0818E-03	3.1457E-01	20.0	5000.0	
7	1	7	10	-5.4149E-08	-3.8782E-01	-3.8167E-02	-2.8613E-03	2.8388E-01	20.0	5000.0	
7	6	7	9	-2.7588E-08	-3.8718E-01	-3.8213E-02	-2.8627E-03	2.8388E-01	20.0	5000.0	
7	1	7	10	-2.8331E-08	-3.8718E-01	-3.8238E-02	-2.8618E-03	2.8387E-01	20.0	5000.0	
7	2	7	9	-5.8479E-08	-3.8787E-01	-3.8787E-02	-2.8697E-03	2.8388E-01	20.0	5000.0	
7	3	7	10	-3.8977E-08	-3.8476E-01	-2.8771E-02	-2.8189E-03	2.7537E-01	20.0	5000.0	
7	5	7	9	-4.0977E-08	-3.8438E-01	-2.8688E-02	-2.8029E-03	2.7508E-01	20.0	5000.0	
7	5	7	10	-5.8889E-08	-3.8428E-01	-2.8771E-02	-2.8191E-03	2.7505E-01	20.0	5000.0	
7	4	7	11	-3.8897E-08	-4.1258E-01	-7.7603E-02	-8.8277E-03	2.9477E-01	20.0	5000.0	
7	5	7	11	-3.8897E-08	-4.1257E-01	-7.7601E-02	-8.8320E-03	2.9477E-01	20.0	5000.0	
7	1	7	9	-3.8883E-08	-5.3248E-01	-7.7603E-02	-8.8140E-03	2.9476E-01	20.0	5000.0	
7	1	7	7	-3.8188E-08	-3.8948E-01	-3.8667E-02	-3.8521E-03	2.8736E-01	20.0	5000.0	
7	1	7	7	-2.8162E-08	-3.8789E-01	-3.8897E-02	-2.8794E-03	2.8788E-01	20.0	5000.0	
7	8	7	14	-5.1488E-08	-4.7797E-01	-5.2897E-02	-6.7773E-03	3.0818E-01	20.0	5000.0	
7	10	7	15	-3.2039E-08	-4.3381E-01	-5.4189E-02	-6.6888E-03	2.8735E-01	20.0	5000.0	
7	16	7	14	-4.8814E-08	-4.3882E-01	-6.3888E-02	-6.7438E-03	3.0804E-01	20.0	5000.0	
7	12	7	14	-3.8881E-08	-4.8881E-01	-7.7603E-02	-8.8140E-03	2.9476E-01	20.0	5000.0	
7	12	7	15	-5.8388E-08	-5.8488E-01	-7.7603E-02	-8.8140E-03	2.9476E-01	20.0	5000.0	
7	13	7	15	-2.4537E-08	-3.8798E-01	-2.8648E-02	-1.7738E-03	2.8733E-01	20.0	5000.0	
7	15	7	14	-3.1818E-08	-2.1181E-01	-1.0981E-02	-1.0981E-03	2.8388E-01	20.0	5000.0	
7	11	7	14	-1.3418E-08	-2.6141E-01	-3.8887E-02	-2.8588E-03	2.8471E-01	20.0	5000.0	
7	11	7	15	-2.8576E-08	-3.5767E-01	-3.5767E-02	-2.8598E-03	2.8498E-01	20.0	5000.0	
I-LINE											
6	1	6	7	-5.8883E-08	-3.8818E-01	-4.8881E-02	-4.1187E-03	2.1683E-01	20.0	5000.0	
6	2	6	8	-3.8818E-08	-3.8818E-01	-4.7948E-02	-4.8117E-03	2.1760E-01	20.0	5000.0	
6	1	6	8	-3.1818E-08	-3.8818E-01	-4.8887E-02	-5.8881E-03	2.1812E-01	20.0	5000.0	
6	3	6	8	-5.1281E-08	-3.8789E-01	-4.8887E-02	-4.1117E-03	2.1871E-01	20.0	5000.0	

Table II-5. Fitting parameters for 2s-2p rate coefficients for neon ions.

In this table the results of classical path calculations for 2s-2p transitions in the neon ions are presented. The first four columns of the table define the transition using the level definitions given in Table II-1. For example, a transition defined by 7 11 7 15 occurs between the eleventh and fifteenth levels of the nitrogen sequence, which means $2sp^4 4p_{3/2}-2p^5 2p_{1/2}$ of Ne IV. The rate coefficient is parametrized by

$$\langle\sigma v\rangle = \frac{1.58 \times 10^{-8} P(B) e^{-B}}{B T_e^{1/2}} \quad B = \frac{\Delta E}{kT}$$

6	2	6	9	-3.5790E-02	-3.0497E-01	-4.7319E-02	-4.6290E-03	2.1701E-01	20.0	5000.0
6	2	8	10	-3.3480E-02	-3.6172E-01	-3.5897E-02	-3.5870E-03	2.3870E-01	20.0	5000.0
6	2	8	12	-3.1190E-02	-3.8402E-01	-3.4690E-02	-3.4777E-03	2.5781E-01	20.0	5000.0
6	2	8	14	-3.0480E-02	-3.9420E-01	-3.4160E-02	-3.5017E-03	2.6757E-01	20.0	5000.0
6	2	10	12	-3.1520E-02	-3.8470E-01	-3.6160E-02	-3.5019E-03	2.5779E-01	20.0	5000.0
6	2	10	14	-3.2325E-02	-3.6700E-01	-3.6870E-02	-3.4943E-03	2.5694E-01	20.0	5000.0
6	3	6	11	-3.6250E-02	-3.6700E-01	-3.6900E-02	-3.5910E-03	2.5670E-01	20.0	5000.0
6	3	8	14	-3.1730E-02	-3.1590E-01	-3.1811E-02	-3.1891E-03	2.1691E-01	20.0	5000.0
6	2	6	14	-3.1790E-02	-3.1260E-01	-3.1811E-02	-3.1790E-03	2.16871E-01	20.0	5000.0
6	2	8	14	-3.1790E-02	-3.1141E-01	-3.2591E-02	-3.2700E-03	2.16871E-01	20.0	5000.0
6	2	13	15	-3.1680E-02	-3.1680E-01	-3.1680E-02	-3.1670E-03	2.1670E-01	20.0	5000.0
6	4	6	15	-2.8470E-02	-2.33170E-01	-3.7530E-02	-3.7670E-03	2.36734E-01	20.0	5000.0
6	5	6	15	-2.1280E-02	-3.7791E-01	-3.1887E-02	-3.1970E-03	2.8794E-01	20.0	5000.0
6	7	6	16	-3.1520E-02	-3.6880E-01	-3.5931E-02	-3.6080E-03	2.6790E-01	20.0	5000.0
6	8	6	17	-3.4430E-02	-3.6880E-01	-3.4729E-02	-3.4817E-03	2.6433E-01	20.0	5000.0
6	8	8	18	-3.0490E-02	-3.0371E-01	-3.1320E-02	-3.1390E-03	2.6433E-01	20.0	5000.0
6	8	8	19	-3.0490E-02	-3.0880E-01	-3.1000E-02	-3.1070E-03	2.6433E-01	20.0	5000.0
6	9	6	17	-3.0470E-02	-3.0790E-01	-3.1753E-02	-3.1823E-03	2.6433E-01	20.0	5000.0
6	9	6	18	-3.4381E-02	-3.0667E-01	-3.5790E-02	-3.1182E-03	2.6433E-01	20.0	5000.0
6	10	6	18	-3.0211E-02	-2.9167E-01	-3.1307E-02	-3.1370E-03	2.6433E-01	20.0	5000.0
6	11	6	17	-3.0780E-02	-2.8767E-01	-3.0416E-02	-3.0480E-03	2.6433E-01	20.0	5000.0
6	11	8	16	-3.0737E-02	-2.8890E-01	-3.1700E-02	-3.1441E-03	2.6433E-01	20.0	5000.0
6	10	5	17	-3.0791E-02	-2.8790E-01	-3.0490E-02	-3.0491E-03	2.6433E-01	20.0	5000.0
6	12	6	17	-3.0301E-02	-2.8881E-01	-3.0470E-02	-3.0480E-03	2.6433E-01	20.0	5000.0
6	11	10	17	-3.5321E-02	-2.8194E-01	-3.1887E-02	-3.1790E-03	2.6433E-01	20.0	5000.0
6	14	6	18	-3.0794E-02	-3.2100E-01	-3.1560E-02	-3.1630E-03	2.6433E-01	20.0	5000.0
6	14	8	17	-3.7610E-02	-3.2100E-01	-3.2961E-02	-3.2960E-03	2.6433E-01	20.0	5000.0
6	14	10	17	-3.2601E-02	-3.2790E-01	-3.1887E-02	-3.1880E-03	2.6433E-01	20.0	5000.0
6	13	6	19	-3.1080E-02	-3.1817E-01	-3.1300E-02	-3.1300E-03	2.6433E-01	20.0	5000.0
6	15	6	19	-2.7142E-02	-3.3077E-01	-3.0430E-02	-3.0430E-03	2.6433E-01	20.0	5000.0
6	15	8	20	-3.6180E-02	-3.1877E-01	-3.3290E-02	-3.3290E-03	2.6433E-01	20.0	5000.0

8-14E

5	1	5	8	-3.1030E-03	-3.5420E-01	-4.1130E-02	-3.8200E-03	2.8100E-01	20.0	5000.0
5	2	5	7	-3.19161E-03	-3.5880E-01	-4.20140E-02	-3.7670E-03	2.7600E-01	20.0	5000.0
5	2	5	8	-3.0840E-03	-3.5840E-01	-4.1080E-02	-3.7670E-03	2.7600E-01	20.0	5000.0
5	1	5	10	-3.17315E-03	-3.29710E-01	-3.8940E-02	-3.8881E-03	3.1018E-01	20.0	5000.0
5	1	5	9	-2.6080E-03	-3.27700E-01	-3.8940E-02	-3.7000E-03	3.0000E-01	20.0	5000.0
5	2	5	10	-2.87370E-03	-3.10810E-01	-3.8160E-02	-3.8160E-03	3.1080E-01	20.0	5000.0
5	2	5	9	-3.0880E-03	-3.0880E-01	-3.0887E-02	-3.0881E-03	3.0790E-01	20.0	5000.0
5	1	5	9	-3.8261E-03	-3.0887E-01	-3.1011E-02	-3.0817E-03	3.0880E-01	20.0	5000.0
5	2	5	9	-3.0790E-03	-3.1160E-01	-3.1070E-02	-3.1031E-03	3.1030E-01	20.0	5000.0
5	3	5	11	-2.9182E-03	-2.9080E-01	-2.9080E-02	-2.1495E-03	2.7308E-01	20.0	5000.0
5	4	5	11	-2.9080E-03	-3.0300E-01	-2.9430E-02	-2.8710E-03	2.7308E-01	20.0	5000.0
5	7	5	11	-2.9071E-03	-2.8110E-01	-2.8110E-02	-2.8110E-03	2.7308E-01	20.0	5000.0
5	7	5	15	-3.0817E-03	-4.1100E-01	-3.0890E-02	-3.1040E-03	3.0150E-01	20.0	5000.0
5	6	5	14	-3.7570E-03	-4.1173E-01	-3.0880E-02	-3.0910E-03	3.0140E-01	20.0	5000.0
5	6	5	15	-3.1180E-03	-4.1180E-01	-3.0790E-02	-3.1400E-03	3.0140E-01	20.0	5000.0
5	7	5	13	-3.0877E-03	-3.9414E-01	-3.0870E-02	-3.1400E-03	3.0140E-01	20.0	5000.0
5	7	5	12	-2.8881E-03	-3.9433E-01	-3.0891E-02	-3.1400E-03	3.0140E-01	20.0	5000.0
5	6	5	12	-3.0880E-03	-3.9440E-01	-3.0891E-02	-3.1400E-03	3.0140E-01	20.0	5000.0
5	6	5	12	-3.0811E-03	-3.9440E-01	-3.0887E-02	-3.1400E-03	3.0140E-01	20.0	5000.0
5	10	5	12	-3.0890E-03	-3.1180E-01	-3.0870E-02	-3.0870E-03	3.1180E-01	20.0	5000.0
5	10	5	12	-3.1180E-03	-3.1180E-01	-3.0870E-02	-3.0870E-03	3.1180E-01	20.0	5000.0
5	10	5	13	-3.0890E-03	-3.1081E-01	-3.1081E-02	-3.1081E-03	3.1080E-01	20.0	5000.0
5	10	5	14	-3.0890E-03	-3.1081E-01	-3.1081E-02	-3.1081E-03	3.1080E-01	20.0	5000.0
5	10	5	15	-2.8880E-03	-3.1017E-01	-3.1473E-02	-3.1473E-03	3.1020E-01	20.0	5000.0
5	8	5	14	-2.8870E-03	-3.1100E-01	-3.0415E-02	-3.0415E-03	3.0420E-01	20.0	5000.0

Table II-5. Fitting parameters for 2s-2p rate coefficients for neon ions.(Continued)

where $\ln P(\beta) = a_0 + a_1 \ln \beta + a_2 (\ln \beta)^2 + a_3 (\ln \beta)^3$.

The coefficients a_0 , a_1 , a_2 and a_3 are given in columns 5-8 and the energy difference ΔE used in the fit is given in column 9. The minimum and maximum temperatures used in the fits are given in the final two columns in eV. The data shown here was taken from a datafile for neon, and the purpose of the table is partly to give an idea of what is involved in the models. Collections of fitting coefficients for different transitions are tabulated in the datafiles, there being different sets for different types of transitions.

5	0	5	15	-3.61162E+00	-3.11823E-01	-3.02610E-02	-2.96289E-03	1.04399E+01	20.0	5000.0
5	0	5	15	-3.32702E+00	-3.76892E-01	-4.84229E-02	-4.35711E-03	2.17899E+01	20.0	5000.0
5	0	5	16	-3.38741E+00	-3.76892E-01	-4.83396E-02	-4.87943E-03	2.17137E+01	20.0	5000.0

BC-LINE										
4	1	4	5	-1.46847E+00	-3.39164E-01	-2.61388E-02	-2.89472E-03	2.88268E+01	20.0	5000.0
4	4	4	8	-2.73290E+00	-3.45387E-01	-4.15758E-02	-3.78538E-03	2.20721E+01	20.0	5000.0
4	3	4	6	-3.32479E+00	-3.45496E-01	-4.62610E-02	-3.67513E-03	2.21833E+01	20.0	5000.0
4	2	4	7	-2.444421E+00	-3.45613E-01	-4.60882E-02	-3.55892E-03	2.21443E+01	20.0	5000.0
4	3	4	7	-3.83396E+00	-3.44686E-01	-4.14201E-02	-3.77388E-03	2.73899E+01	20.0	5000.0
4	3	4	6	-3.55389E+00	-3.47333E-01	-4.21442E-02	-3.68773E-03	2.69299E+01	20.0	5000.0
4	4	4	7	-3.89177E+00	-3.46883E-01	-4.27468E-02	-3.68378E-03	2.71829E+01	20.0	5000.0
4	3	4	9	-2.35498E+00	-2.76224E-01	-2.27388E-02	-1.65192E-03	1.27413E+01	20.0	5000.0
4	5	4	10	-3.88389E+00	-3.44686E-01	-4.13578E-02	-3.76710E-03	2.73982E+01	20.0	5000.0

LI-LINE										
3	1	3	2	-3.21888E+00	-2.61447E-01	-2.6.687E-02	-1.67951E-03	1.58888E+01	20.0	5000.0
3	1	3	3	-2.58539E+00	-3.14788E-01	-3.37711E-02	-2.81231E-03	1.58894E+01	20.0	5000.0

Table II-5. Fitting parameters for 2s-2p rate coefficients for neon ions. (Continued)

D. Photoionization

The photoabsorption cross section of neon is one of the more accurately known of all the neutral elements (hydrogen and helium taking the first two places). We have used the results quoted by Wuilleumier (1973) for single electron ionization cross sections. The errors in this data are small; the 2p cross section is accurate to within 5% and the 2s cross section is accurate to within 7%. The quantities measured in Wuilleumier (1973) are the relative values of the 2s and 2p single electron cross sections, and with that information and knowledge of the shake-up and shake-off fractions one may compute the absolute 2s and 2p single electron cross sections given data for the absolute photoabsorption cross sections, which were determined elsewhere. More recent values for the total photoabsorption cross section of neon has been given in Marr and West (1976) and West and Marr (1976), in which the absolute error in the values quoted is probably 2% or 3%. With knowledge of the shake-off and shake-up cross sections as well as the relative values of the 2s and 2p cross sections, it is possible to estimate cross sections which are more accurate than the Wuilleumier data. In our case the major errors in the calculations described in Chapter III arise from the calculation of the pumping spectrum (the error may be a factor of two), and the errors associated with the single electron photoionization cross sections of neutral neon are small in comparison. We shall be content with the data of Wuilleumier.

Extensive theoretical work has been done on neon, and we refer the reader to the selection of papers given in the references. In the case of neon ions, there is no experimental data available, and so for these cross sections we require theoretical results. Some elaborate theoretical results have been reported at low energies for the first few ions by Pradham (1979). Ideally one would like data of similar quality for all the ions at higher energies; however, to date such studies have not been published. In the development of some atomic physics codes to do elaborate calculations one begins with the simplest steps, and here we discuss briefly some simple calculations which were performed using a distorted wave model. These calculations do not represent "state of the art" results; however, we have found similar results to be very useful here and in the study of many other models.

In the case of 2p photoionization, the 2p subshell cross sections were computed from single configuration bound state Hartree Fock wave functions, and from free electron wave functions satisfying the Schrodinger equation

$$\left(-\frac{1}{2} \frac{d^2}{dr^2} + \frac{\ell(\ell+1)}{2r^2} + V(r) - \epsilon \right) F(r) = 0 \quad (\text{II-D-1})$$

where the potential used was taken to be

$$V(r) = -\frac{Z}{r} + \sum_i q_i \int_0^\infty p_i^2(r') \frac{1}{r} dr' \quad (\text{II-D-2})$$

where $P_i(r)$ is a bound Hartree Fock wave function and $r_>$ is the greater of r and r' . Atomic units have been used. The Hartree Fock wave functions were generated by a program written by the author using standard numerical methods as described in Fischer (1978). The wave functions were computed on a logarithmic mesh using the Noumerov method without the deferred difference correction. Most of the numerical error in the determination of the wave function comes from the boundary condition at the origin, and is typically of the order of 10^{-6} or less. The free wave functions were computed from a WKB expansion (Belling (1968))

$$F(r) = n^{-1/2}(r) \sin\left(\int_0^r n(r') dr'\right) \quad (\text{II-D-3})$$

where the phase variable $n(r)$ satisfies

$$\begin{aligned} n^2(r) = & \left\{ 2\epsilon - 2V(r) - \frac{\ell(\ell+1)}{r^2} \right\} \\ & + n^{1/2}(r) \frac{d^2}{dr^2} n^{-1/2}(r) \end{aligned} \quad (\text{II-D-4})$$

This equation was solved by iteration for large r , and then the integration was carried out for the variable $n(r)$ for $\ell=0$ and $-2\ln(r)$ for $\ell \neq 0$.

As a result of working with these variables on a logarithmic mesh, the free electron wave function could be determined accurately on the same mesh as the bound state wave functions. As a practical matter, this means that the program can compute wave functions accurately

for low or high energies and for small or large angular momentum. Such properties are useful in distorted wave collisional cross section calculations for which the code was developed. Given the free and bound state wave functions, the cross section was computed in the length form (as opposed to the velocity form or acceleration form which gives slightly different results)

$$\sigma(E) = \frac{4\pi\alpha_0^2}{3} \left(\frac{E - \Delta E}{I_H} \right) \sum_{\ell' = \ell \pm 1} \left(\frac{\ell_{>}}{2\ell + 1} \right) \times \left| \int_0^\infty P_{n\ell}(r) r F_{E,\ell'}(r) dr \right|^2 \quad (II-D-5)$$

where $P_{n\ell}(r)$ and $F_{E,\ell}(r)$ are the bound and free wave functions, and where α is the fine structure constant and $\ell_{>}$ is the greater of the angular momenta of the two wave functions. The radial integral was done using a quadrature suitable for sinusoidal functions on a logarithmic mesh, and the error in the integration was typically of order 10^{-3} or less.

The cross sections computed in this approximation for the 2p shell of neutral neon were in rough agreement with the experimental results as shown in Figure II-2, and therefore the cross sections computed from this model for the neon ions were adopted (the results are shown in Figure II-3). In the case of 2s photoionization, the Hartree Fock 2s wave function for Ne I is in error sufficiently to cause the photoionization cross section to be overestimated by about

Figure 11-2. Comparison of theoretical and experimental 2p single electron photoionization cross sections for Ne I.

The experimental values are given by Wuilleumier (1973) and the theoretical value was computed in the distorted wave approximation from single configuration Hartree Fock wavefunctions. The point of the comparison is to get some idea of the accuracy of the theoretical predictions using a very simple theory (state of the art random phase and many body calculations are much more accurate, however, results do not yet appear for the neon ions in the literature). The agreement is reasonable in this case whereas in the case of 2s electrons the theoretical values are much higher ($\times 1.5$ - $\times 2.0$) than the experimental values. The distorted wave results for 2p electron cross sections of the neon ions may be of comparable accuracy to the results shown above.

Figure II-3. Theoretical photoionization cross sections (per electron) for 2p electrons for neon ions.

The distorted wave approximation using Hartree Fock bound state wave functions was used to compute the 2p cross sections. One observes an increase in the cross section per electron as more electrons are removed. This is similar to the trend noted by Missavage et al (1977) in their study of the photoionization of the oxygen ions. The distorted wave results for the lithium-like system was in excellent agreement with the accurate calculations of McDowell and Chang (1969).

50%. The cross sections for the neon ions were therefore estimated semi-empirically using the results of McDowell and Chang (1969) for the Li-like and the experimental result for the neon-like systems, as the distorted wave cross sections for the ions are probably too large. The results are shown in Figure II-4.

The photoionization cross sections were least squares fitted to the form

$$\sigma(E) = 10^{-18} n \left(\frac{I_H}{\Delta E} \right)^n P(E/\Delta E) \text{ cm}^2 \quad (\text{II-D-6})$$

where n is a multiplier, I_H is 13.606 eV, and ΔE is the ionization energy. The function $P(x)$ is defined in II-C-9.

Results for the L-shell are given in Table II-6. The K-shell cross section for Ne I was taken from Wulfeumier (1973), and as yet the more highly stripped systems do not include K-shell data as the Auger branching ratios in LSJ coupling do not appear in the literature.

Figure II-4. Single electron 2s photoionization cross sections for neon ions.

The ionization cross section for Ne I is known experimentally from the study of Wuilleumier (1973), and the distorted wave results for Ne VII are assumed to be accurate. Since a discrepancy exists between the Ne I experimental and theoretical results due to correlation effects on the 2s electron (Wuilleumier and Krause (1974)), I assume that similar effects will occur in Ne II and Ne III. The lithium-like system should be less sensitive to correlation, hence the lithium-like cross section is taken to be accurate. The remaining cross sections were obtained by reducing the distorted wave values by empirical factors to get agreement at Ne I.

ION	EO	A1	A2	A3	ΔE	EMIN	EMAX
2P ELECTRIC							
NE I	9.7949E-01	5.7233E-01	-1.0457E-02	0.7321E-02	2.4750E-01	2.0000E-01	1.0000E-03
NE II	1.0317E+00	-2.0002E-01	-1.1973E-00	1.7007E-01	4.4600E-01	4.0000E-01	1.0000E-03
NE III	2.1530E+00	-1.2000E-00	-0.6240E-01	1.1730E-01	0.0000E-01	0.0000E-01	1.0000E-03
NE IV	2.1261E+00	-2.0117E+00	-0.0762E-01	0.0000E-02	0.3000E-01	0.0000E-01	1.0000E-03
NE V	1.0070E+00	-0.3272E+00	-5.1500E-01	0.0000E-02	1.0140E-02	1.0000E-02	1.0000E-03
NE VI	1.0000E+00	-2.7037E+00	-3.0000E-01	5.0121E-00	1.5770E-02	1.0000E-02	1.0000E-03
NE VII	1.5012E+00	-2.0000E+00	-2.1000E-01	2.5000E-02	1.0000E-02	1.0000E-02	1.0000E-03
NE VIII	1.2010E+00	-2.0000E+00	-5.1000E-01	1.5000E-01	2.0000E-02	2.0000E-02	1.0000E-03
2S ELECTRIC							
NE I	-2.0049E-01	0.3302E-02	-6.1430E-01	4.3040E-02	4.0000E-01	5.0000E-01	1.0000E-03
NE II	3.4050E-01	-5.7004E-01	0.0047E-01	-5.0001E-03	7.0000E-01	7.1000E-01	1.0000E-03
NE III	0.0071E-01	-0.1271E-01	-5.0000E-01	0.0000E-02	0.0000E-01	0.0000E-01	1.0000E-03
NE IV	0.0100E-01	-1.0040E+00	-4.1277E-01	2.2001E-02	1.1000E-02	1.10-00E-02	1.0000E-03
NE V	1.0070E+00	-1.2240E+01	-4.0077E-01	0.1000E-02	1.4200E-02	1.0000E-02	1.0000E-03
NE VI	1.1401E+00	-1.5000E+00	-3.5023E-01	-1.1000E-02	1.7100E-02	1.7100E-02	1.0000E-03
NE VII	1.2010E+00	-1.0310E+00	-1.3000E-01	-2.0010E-02	2.0010E-02	2.0010E-02	1.0000E-03
NE VIII	1.2500E+00	-1.0040E+00	-5.2000E-01	7.0001E-02	2.0000E-02	2.0000E-02	1.0000E-03

Table II-6. Fitting coefficients for neon ion L-shell single electron photoionization cross sections.

These coefficients were obtained through least squares fitting of the photoionization cross section data described in the text. The cross sections are fit to the form

$$\sigma(E) = 10^{-18} \left(\frac{13.606}{\Delta E(\text{eV})} \right) P(E/\Delta E) \text{ cm}^2$$

$$P(x) = \exp(a_0 + a_1 \ln x + a_2 (\ln x)^2 + a_3 (\ln x)^3).$$

The fitting coefficients, a_0 , a_1 , a_2 , a_3 and ΔE are given in the above table and EMIN and EMAX refer to the smallest and largest photon energies used in the data which was fitted. This fitting procedure is more accurate than the form used by Henry (1970), and less general than the 10 parameter fits of Barfield et al (1972).

The next issue requiring attention is the branching ratios in intermediate coupling. The problem is the following: given an ion of quantum numbers $2s^m p^n LSJ$, and given that a 2s or 2p electron is photoionized, what fraction of the atoms go into a state $L'S'J'$ of the next isoelectronic sequence: In order to do the problem accurately, one should use accurate wave functions and include shake-up and shake-off processes for all transitions. For our purposes, we shall content ourselves for the time being with much cruder estimates, namely, we shall adopt the branching coefficients

$$\begin{aligned}
 \left| G_{L_0 S_0 J_0}^{LSJ} \right|^2 &= \left| G_{S_0 L_0}^{SL} \right|^2 (2L+1)(2S+1)(2J+1) \\
 &\times \sum_{\substack{M_{L_0} \\ M_{S_0} M_{L_0} \\ M_S M_L \\ m_S m_L}} \left(\begin{matrix} L_0 & l & L \\ M_{L_0} & m_L & -M_L \end{matrix} \right)^2 \left(\begin{matrix} S_0 & s & S \\ M_{S_0} & m_S & -M_S \end{matrix} \right)^2 \left(\begin{matrix} S & L & J \\ M_S & M_L & -M \end{matrix} \right)^2 \left(\begin{matrix} S_0 & L_0 & J_0 \\ M_{S_0} & M_{L_0} & -M_0 \end{matrix} \right)^2
 \end{aligned}
 \tag{II-D-8}$$

where the terms of the summation are composed of squares of 3-j symbols (see Sobel'man (1972)) and

where $\left| G_{S_0 L_0}^{SL} \right|^2$ are fractional parentage coefficients and $\left| G_{L_0 S_0 J_0}^{LSJ} \right|^2$ are the branching coefficients. The fractional parentage coefficients were taken from Sobel'man (1972) for ionization of p electrons in the p^n and $s^2 p^n$ configurations, and from Menzel and Goldberg (1936) for the sp^n configurations. These coefficients range between zero and unity and the sum over the initial or final quantum numbers yields unit.

The accuracy of this approximation can be examined in the case of 2p ionization from Ne I. For example, we have assumed that 2p photoionization leads to Ne II $2s^2 p^5 \ ^2P_{3/2}$ formation twice as often as $2s^2 p^5 \ ^2P_{1/2}$ formation, in accordance with the values of the statistical weights. Ong and Manson (1980) discuss calculations and experimental results; the experimental results place the $^2P_{3/2}$ state in a favored position relative to the $^2P_{1/2}$ state by about 10%, while relativistic random phase results predict a much smaller value.

E. Electron Collisional Excitation and Ionization Processes

The most important collisional process involved in the inversion kinetics of the Ne II laser scheme is the removal of 2p electrons from Ne I, which results in the production of Ne II ions in the lower laser state. Such a process is fatal to the population inversion as there is no mechanism to remove lower state population effectively once it builds up. Typically in the kinetics calculations which have so far been done, inversions can be produced readily if there is reasonable flux near 500 eV; however, once the electron density builds up, an electron collisional avalanche is triggered leading to complete destruction of the laser gain. This process ultimately limits the maximum gain achievable, and it is for this reason that one is interested in finding accurate data relevant to the calculation of 2p ionization and excitation processes for Ne I.

In the model under consideration in this chapter we have not included levels in the kinetics model corresponding to states with excited M-shell or N-shell electrons. Details with respect to more complex models may appear elsewhere. In this case the contribution of multi-step processes (such as collisional excitation followed by ionization) should be included in the calculation of the total 2p removal rate, even though the intermediate levels themselves are not included explicitly. For this reason we shall make the assumption that all Ne I atoms which are collisionally excited result in Ne II $2s^2p^5$ states immediately. This would be true if M-shell or N-shell

electrons were ionized very rapidly. In the laser target under consideration we expect the suprathermal electrons to play a large role in determining the local electron temperature, hence electron temperatures in excess of 1 keV may occur. Under these conditions, excitation of 2p electrons is less likely to occur by severalfold than ionization, and the excitation channel for removal of 2p electrons does not play a dominant role. Neglect of the M-shell states of Ne I and inclusion of the excitation rate for their production into an effective 2p removal rate leads to a model whose predictions may be considered to be conservative in terms of the Ne II gain which is calculated. Further discussion of these issues may appear elsewhere.

The situation in the more highly stripped ions is less important in terms of Ne II inversion kinetics. As a practical matter ionization does not occur past the nitrogen-like sequence under conditions at all realistic relevant to the target designs so far considered. Modeling of electron collisional ionization is straightforward; however, excitation rates become increasingly larger compared to ionization rates as the ions become more highly stripped. M-shell excitation will be followed by radiative decay most of the time at low densities in the lithium-like sequence, but the situation is not so clear in the first few ionization stages. For the present purposes we shall be content with including excitation and ionization into a single removal cross section and keep in mind the approximations

associated with the model. Such a model will be useful in providing conservative estimates of laser gain in the first few ionization sequences and in this work we shall ask nothing further of it.

The electron collisional ionization cross section of Ne I has been determined experimentally by Fletcher and Cowling (1973), Luyken et al (1972), Schram et al (1965) and Rapp and Englander-Golden (1965). The results discussed in this section were taken mostly from these works. More recently Nagy et al (1980) have reported a direct measurement of single and multiple ionization cross sections for the rare gases, and the accuracy of the experimental results is quoted to be about 10%. In Figure II-5 is shown the single electron ionization cross sections of the L-shell electron and the corresponding rate coefficient. One observes that at an electron temperature of 1 keV, 2p electrons are ionized more often than 2s electrons by roughly a factor of three on a per-electron basis. Since the inversion process is dependent upon removing 2s electrons preferentially, one finds that collisions always work against the inversion.

The excitation cross sections for 2p electrons are shown in Figure II-6. The 2p-3s cross section was calculated by Sawada, Purcell and Green (1971) in the distorted wave approximation for low energy, and the high energy result was taken from the Bethe coefficients given by Dayashankar (1979). The 2p-3p cross section has been measured by Sharpton et al (1970) at low energies, and the high energy data was again from the Bethe cross sections and sketching in a reasonable

Figure II-5. Ne I single electron collisional ionization cross sections and rate coefficients.

The cross sectional data is taken from the experimental works of Fletcher and Cowling (1973), Luyken et al (1972), Schram et al (1965) and Rapp and Englander-Golden (1965). One observes that the cross section for 2p ionization is about a factor of three larger per electron than the cross section for 2s ionization. In terms of laser kinetics, this means that electron collisional processes are destructive to the inversion process as removal of a 2p electron from Ne I produces a Ne II ion in one of the lower $2s2p^5$ states. As the inversion process is transient, the gain cannot recover after the lower laser state has become populated. The rate coefficient was computed by integrating the cross section over a Maxwellian at a set of electron temperatures. One observes that the rate coefficients are a slow function of electron temperature between 10^2 eV and 10^4 eV, falling slowly at the higher temperatures.

Figure II-6. Collisional excitation and ionization per electron for cross sections Ne I 2p electrons.

These cross sections were taken from the experimental data of Sharpton et al (1970) and from the theoretical data given by Sawada et al (1971) and Dayashankar (1979). The cross sections described by 2p-3s include the transitions from the $2s^2p^6$ ground state to all the multiplets with configuration $2s^2p^53s$. Similarly the 2p-3p and 2p-3d curves are for transitions to the $2s^2p^53p$ and $2s^2p^53d$ states. The curve denoted by 2p-c is the ionization cross section for the process $2s^2p^6-2s^2p^5+e^-$. All cross sections are per electron and the total cross sections for these processes are greater by a factor of six due to the six 2p electrons of Ne I. One observes that the ionization cross section is greater than all excitation cross sections above 30 eV. The different (non-dipole) asymptotic form of the 2p-3p cross section is observed. The rate coefficients were obtained by integrating the cross sections over Maxwellians at different electron temperatures.

low-energy behavior. The cross sections for 2p ionization and 2p excitation and ionization are shown in Figure II-7. One observes that at high electron energy the additional contribution to the cross section is small; however, at low energy the excitation process is dominant and results largely in excitation to the 3s states. Under laser conditions where the electron distribution is dominated by the suprathermal electrons, the differences should be small.

In Figure II-8 is shown the cross sections for 2s ionization and 2p excitation and ionization. Excitation of the 2s electron results in states which are unstable to Auger decay and are hence included in the 2p ionization cross section.

The next task is to estimate cross sections for excitation and ionization of all of the other isoelectronic sequences. For purposes of the present model, we have adapted the semi-empirical results shown in Figure II-9 for 2p excitation and ionization, and those shown in Figure II-10 for 2s excitation and ionization in all ions except for Ne I. The results shown in the two figures were obtained by the following procedure; the total excitations and ionization cross sections for the neutral case and for the lithium-like case were plotted. The lithium-like cross sections were estimated from the semi-empirical method of Golden and Sampson (1978) (eq. 21 is incorrect and the correct result is found in their earlier work - Sampson and Golden (1974) eq. (19)) for excitation and Sampson and Golden (1977) for ionization. The resulting single electron cross sections were

Figure II-7. Cross sections per electron for electron collisional ionization and for excitation and ionization.

The cross section for 2p ionization and all excitations was estimated from the experimental data (see text) for ionization and for some excitation processes. The excitation to the high Rydberg states is included using the Bethe coefficients given by Dayashanker (1979). One observes that the two cross sections differ by roughly 20% at high electron energies. The assumption implied through use of the total cross section is that if a Ne I $2s^2p^6$ ground state atom is excited to a Ne I $2s^2p^5n\lambda$ level, that ionization occurs to the Ne II $2s^2p^5$ states before any other process (for example radiative decay) takes place. (The notation $n\lambda$ refers to a highly excited orbital or principal quantum number n and angular momentum λ .)

Figure II-8. Collisional data (per electron) for 2s ionization and for 2p ionization excitation and ionization of Ne I.

The computation of the cross sections includes summation of ionization and excitation cross sections for the 2p result and only ionization data for the 2s result. Excitation of 2s electrons in Ne I leads to autoionization in most cases as the first $2sp^5$ excited state lies above the ionization limit for 2p electrons. The dominant electron collisional process is therefore removal of a 2p electron from Ne I producing a Ne II $2s2p^5$ (lower lasing state) ion. From the rate coefficient one finds that the electron densities at which collisional processes become important are of order $10^{17}/\text{cm}^3$ which corresponds to a rate of $10^9/\text{sec}$ if the rate coefficient is $10^{-8} \text{ cm}^3/\text{sec}$.

Figure II-9. Estimates of the total electron collisional excitation and ionization data for neon 2p electrons (per electron).

The cross sections for 2p excitation and ionization for Ne I and Ne VIII were computed from data given in the literature. The data for the intermediate ions is semi-empirical. Of interest at present are the rate coefficients at high (keV) electron temperatures in which case the empirical approach should be adequate. The assumption made in the use of these total excitation and ionization cross sections is that excitation is followed by further excitation and ionization rather than by deexcitation. In our case the model resulting will be conservative in terms of gain predictions.

Figure III-10. Estimates of the total electron collisional excitation and ionization data for neon 2s electrons (per electron).

The variation of the total 2s excitation and ionization cross sections is less than in the case of 2p electrons (see the preceding figure). In the case of Ne I we have estimated the 2s excitation cross sections (2s excitation usually results in Auger decay and is included in the 2p ionization process) and have added it to the 2s ionization cross section to yield the Ne I curve shown above. In the model we have used the true 2s ionization cross section for Ne I and for the neon ions we have used total excitation and ionization cross sectional results.

found to be roughly within a factor of two of each other and one might be tempted to guess that the cross sections for the intermediate states must be between the two limits. At high energies, this can probably be shown to be the case through relating the Bethe cross section to the photoionization cross section and bound-bound oscillator strengths, although this computation has not yet been carried through by the present author, nor does it seem to appear in the literature. The cross sections used here were obtained by picking curves which look reasonable between the two limits of neon-like and lithium-like cross sections. The results were tabulated and the resulting rate coefficients are given in Table II-7. We shall not attempt a justification of this procedure here. Subsequent work has been devoted toward models with M-shell states and results may appear elsewhere.

ION	A0	A1	A2	A3	DE	TMIN	TMAX
2P ELECTRONS							
NE I	-4.00087E+00	-1.27987E+00	-1.60043E-01	-7.44130E-03	1.67660E+01	1.00000E+01	1.00000E+04
NE II	-2.59233E+00	-8.19753E-01	-6.45053E-02	-3.29587E-05	2.79000E+01	1.00000E+01	1.00000E+04
NE III	-1.63025E+00	-6.70876E-01	-5.13769E-02	-7.17261E-05	4.01000E+01	1.00000E+01	1.00000E+04
NE IV	-1.29437E+00	-5.15519E-01	-1.72634E-02	2.84630E-03	5.39000E+01	1.00000E+01	1.00000E+04
NE V	-9.83160E-01	-4.29054E-01	-6.93099E-03	2.99205E-03	6.70000E+01	1.00000E+01	1.00000E+04
NE VI	-5.74007E-01	-4.17313E-01	-3.12022E-02	1.70203E-03	8.30000E+01	1.00000E+01	1.00000E+04
NE VII	-3.19116E-01	-3.86240E-01	-4.34068E-02	-5.18711E-03	1.00100E+02	1.00000E+01	1.00000E+04
NE VIII	-8.02339E-02	-3.41433E-01	-4.58275E-02	-7.09341E-03	1.20500E+02	1.00000E+01	1.00000E+04
2S ELECTRONS							
NE I	-3.13638E+00	-1.00046E+00	-1.39053E-01	-7.98989E-03	3.80000E+01	1.00000E+01	1.00000E+04
NE II	-2.43980E+00	-7.24786E-01	-6.97280E-02	-1.70713E-03	5.29000E+01	1.00000E+01	1.00000E+04
NE III	-2.01126E+00	-5.53369E-01	-2.09043E-02	3.43132E-03	6.60000E+01	1.00000E+01	1.00000E+04
NE IV	-1.76975E+00	-5.11406E-01	-3.03573E-02	8.15797E-04	7.92000E+01	1.00000E+01	1.00000E+04
NE V	-1.56763E+00	-4.29234E-01	-4.42754E-03	3.87537E-03	9.30000E+01	1.00000E+01	1.00000E+04
NE VI	-1.35931E+00	-4.22951E-01	-2.75327E-02	-1.02086E-03	1.07000E+02	1.00000E+01	1.00000E+04
NE VII	-1.15529E+00	-4.17997E-01	-4.27783E-02	-4.54512E-03	1.20100E+02	1.00000E+01	1.00000E+04
NE VIII	-1.24525E+00	-3.27059E-01	-2.14160E-03	2.15657E-03	1.36000E+02	1.00000E+01	1.00000E+04
NE I	-2.24872E+00	-1.04273E+00	-1.53116E-01	-8.80493E-03	4.85000E+01	1.00000E+01	1.00000E+04

Table II-7. Fitting coefficients for neon collisional data.

These coefficients were obtained through least squares fitting of the electron collisional rate coefficients describing collisional ionization and collisional excitation followed by ionizations as described in the text. The rate coefficients are fit to the form

$$\langle\sigma v\rangle = \frac{1.58 \times 10^{-8} P(\beta) e^{-B}}{\beta T_{\text{ev}}^{3/2}} \text{ cm}^3/\text{sec}$$

$$B = \frac{\Delta E(\text{eV})}{T_{\text{ev}}}$$

$$P(\beta) = \exp(a_0 + a_1 \ln \beta + a_2 (\ln \beta)^2 + a_3 (\ln \beta)^3).$$

The fitting coefficients a_0 , a_1 , a_2 , a_3 and $\Delta E(\text{eV})$ are given in the above table. TMIN and TMAX are the lowest and highest thermal electron temperatures used in the data which was fitted. The form of the fit was taken from Van Regemorter (1962) and has been useful in fitting rate coefficients to within a few percent accuracy over several orders of magnitude electron temperature. The last entry in the above table is for single electron ionization (only) of 2s electrons from Ne I. All other entries include excitation plus ionization.

F. L-Shell Double Photoionization Processes

When neon is photoionized below the K-edge, one observes that double ionization occurs some of the time. Carlson (1967) reported the results of charge spectrometry experiments wherein the ratio of double photoionization to single photoionization at 500 eV is 0.136 ± 0.007 . When an L-shell electron is removed suddenly, the potential seen by the electrons left behind is altered, forcing the remaining L-shell electron orbitals to relax to the new potential. In some cases this rearrangement leads to excitation (shake-up) and in other cases ionization (shake-off) occurs. In neon for L-shell ionization the latter process is the dominant one.

Experiments involving X-ray photoionization have been carried out by Carlson (1967), Lightner et al (1971), Samson and Haddad (1974) and Schmidt et al (1976) with reasonable agreement between 200 eV and 300 eV. Carlson (1967) reports a value of 13.7% at 225 eV and Schmidt et al gives 13.5% at the same energy. At 280 eV Lightner et al (1971) reports $17.1 \pm 1.7\%$ which is not far off from the other data. A recent experiment by Holland et al (1979) gives values which are lower (11.5% at 220 eV and 6.3% at 280 eV). The discrepancy between the experiments has not yet been accounted for and at present there is apparently uncertainties as to precise values for the ratios.

According to shake-off theory, the ratio considered above is nearly independent of the mode of ionization of the first electron. Equivalent

data concerning the ratio can be obtained from electron ionization experiments where the incident electron is energetic (many keV). Single electron ionization occurs with small energy transfer, after which shake-off may occur. If the energy loss of the initial fast electron is known, then the shake-off fraction should be the same for a photon of that known energy. The early data of Van der Wiel and Wiebes (1971) appears to be in error, and the later results of Wight and Van der Wiel (1976) appear to be closer to X-ray experiments.

Theoretical results have been reported by Chang and Poe (1975) and Carter and Kelly (1977), both of which involved many body perturbative techniques. The agreement between theory and experiment is fair (see Figure II-11). In the present work, we are interested in the cross sections for the shake-off process to the low-lying oxygen-like neon sublevels. The experimental results yield data on the ratio of singly ionized to doubly ionized ions, but say nothing about the final state of the doubly ionized ions. For the data of interest, S. Carter (1979) has kindly provided unpublished material containing the individual cross sections for transitions to the $2s^2p^4$ and $2sp^5$ states of the oxygen-like ion. No similar calculations have been reported for the ss processes leading to the $2p^6$ state. As was noted earlier the cross section for single 2s photoionization increases with increasing photon energy relative to the 2p; therefore, one might expect the ionization of two 2s electrons would be important at high energies (for example, near 500 eV where current interest is focused).

Figure II-11 Ratio of Two Electron Ionization Cross Section to Single Electron Ionization Cross Section for Low Energy Ne I Ionization

When an L-shell electron is removed from Ne I through photoionization or collisional ionization, the bound electrons which remain find themselves in a potential different than before ionization. Rearrangement occurs and some of the time a second electron is ionized. In principle, this process can remove several electrons, although the cross sections for the removal of each additional electron leads to a cross section smaller by severalfold. In the figure above we show results for X-ray ionization experiments, electron ionization experiments, and many body theoretical results. In the case of electron ionization experiments, high energy electrons are used to produce ionization (Van der Wiel and Wiebes (1971) used 10 keV electrons) and results are given in terms of energy transfer. Although ionization by photons and ionization by electrons are very different processes, the resulting atomic rearrangement is nearly the same for both.

Carlson and Nestor (1973) have published shake-off fractions for the different subshells, and this seems to be the only estimate available for the ss shake-off process. His calculation assumed that a vacancy occurred in a given subshell, independent of the energy of the ionized electron, and that single electron Hartree-Fock-Slater wavefunctions described the orbitals. The resulting shake-off probabilities are much lower than the observed values (factor of 3) and do not appear to be reliable. We have estimated the ss contribution by normalizing the Carlson and Nestor result to a value close to the midrange experimental values and filling in the low energy behavior with a curve of reasonable characteristics.

The results of Carter (1979) for the pp and sp processes and the above empirical result for the ss contribution were normalized to a curve which seemed to be closer to the X-ray experimental results (excluding the high energy data of Holland et al (1979) which may be suspect). The data of Carter (1979) extends only up to 280 eV, and for energies higher we have assumed that the ratios involved varied slowly and smoothly at higher energies. The data of Marr and West (1976) was used for converting ratios into absolute cross sections. The final results are shown in Figure II-12.

Figure II-12. Empirical shake-off cross sections for the L-shell photoionization of Ne I

Using the ratio of the two electron photoionization cross section to single electron cross section as shown in Figure II-11, and estimating the individual LS cross sections from theoretical data as discussed in the text, we have arrived at the empirical results shown above. The data given is in LS coupling, and to estimate the LSJ cross sections which are required by the model we have divided the cross section for the two 3P states between the different J sublevels statistically. Triple photoionization may occur where energetically allowed; however, the cross sections for this process are much smaller than for double ionization and we have therefore neglected it. Double photoionization of the neon ions is also possible, although data appears to be lacking completely for the L-shell case in the literature.

G. K-Shell Auger Processes

The Auger process following K-shell ionization in neon has been studied by many authors (see references), although so far no one has published data for branching ratios in LSJ coupling except for Ne I. Experimental results in the case of Ne I are reported by Krause et al. (1974) and the rates are given by Howat et al. (1978). The total Auger decay rate of the $1s2s^2p^6\ ^2S$ state is $3.5 \times 10^{14} \text{ sec}^{-1}$, and 60% of the time the decay is to the Ne III $2s^2p^4\ ^1D$ state. The $2s^2p^4\ ^1D$ state comprises only 26% of all the Ne III states, therefore it is populated preferentially by the Auger process. This means that inverting Ne III lines with respect to the 1D state is difficult in the presence of K-shell ionizing radiation is difficult. Auger decay to the Ne III $2s^2p^4\ ^3P$ state is forbidden by parity selection rules and the observed rate is very small (decay occurs about 0.3% of the time). Decay to the Ne III $2sp^5\ ^3P$ state occurs about 6% of the time, leading one to note that the $2s^2p^4\ ^3P-2sp^5\ ^3P$ transition at 25.3 eV can be inverted readily. The use of Auger decay to channel population into an upper state of a candidate laser transition is well known, and this approach is simply another application of the approach.

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III. Neon Laser Design

In the first chapter we have described the 27 eV EUV laser scheme involving Ne II, and in Chapter II we have reviewed some of the atomic physics involved in the inversion kinetics. In this chapter we consider a target design intended to implement the self-terminating photoionization scheme, and study the target operation under intense infrared laser irradiation as viewed through the porthole allowed by numerical simulation. The basic problem is to convert intense ($5 \times 10^{13} \text{ W/cm}^2 - 10^{14} \text{ W/cm}^2$) 1.06μ laser radiation into band passed limited X-radiation which is sufficiently bright and suitably filtered to effectively ionize 2s electrons selectively from neon gas.

Early on in the research we had considered the use of cylindrical short wavelength laser targets; for example, as is shown in Figure IV-1. Thin film cylindrical targets have the advantage of being straightforward to calculate; however, the actual construction of such a target, though not impossible, is sufficiently difficult to provide motivation for seriously considering other geometries. Figure IV-1 serves to illustrate some of the salient features of the target under consideration, which is why we pause to discuss it here. Neon gas flows through the center of the target at low pressure ($4 \times 10^{17}/\text{cm}^3$ translates into 12 torr at room temperature, which would be a reasonable density if the target were irradiated with a short 50 psec pulse), surrounded by a three-layer flashlamp and filter system. Outside the flashlamp and

Figure III-1. Ne II laser target design in cylindrical geometry.

The laser target shown here is a gas tight cylinder with Be/LiF/Mo walls and neon gas in the interior. The cylinder ends are capped with a thin layer of parylene which were to be heated to several eV and hence become transparent to 27 eV radiation during the course of laser irradiation. This type of design was to be a prototype to be followed by similar designs in cylindrical geometry for higher energy soft X-ray laser schemes. This target has the advantage of being easier to model, however, it has been deemed unbuildable by several target fabrication groups. Current laser designs are in alternate geometries, nevertheless, the cylindrical target design is useful for the issues of interest here.

filter system is the tenuous sub- 10^{-5} torr vacuum of the target chamber. Irradiation of the flashlamp causes heating and general ablation of the flashlamp layer, production of X-rays, and generation of suprathermal electrons through resonance absorption of the intense laser light at the critical surface. As the X-rays pass through the filter material, low energy photons are removed effectively, and some filtering occurs above the pumping window near 500 eV. The suprathermal electrons wander throughout the target, suffering inelastic collisions and direction-changing collisions as they pass through the filter layers. Penetration of the hot electrons into the laser medium occurs, and we assume that the thermal electrons which may be present are displaced by hot electrons through a strong electrostatic field which is set up in the interior of the laser medium. The resulting electron temperature is roughly equal to the hot electron temperature, which at an infrared laser pump intensity of 10^{14} W/cm² is several kilovolts. Meanwhile the ion temperature remains near its initial value of 300°K, since electron-ion collisional coupling is weak and there are no strong mechanisms which can raise the ion temperature very much higher. We have examined a novel process, termed "proximity" heating, which occurs when atoms close to one another are ionized and repel each other (Coulomb repulsion), carrying away kinetic energy much in excess of the thermal average energy. This effect results in about 30% increase in the ion temperature. The ion temperature is of interest because the dominant broadening mechanism in determining the small signal laser gain is Doppler broadening.

The problem of converting infrared laser radiation to broad band X-radiation for the purpose of pumping short wavelength laser systems has been of interest since the early days of X-ray laser research (Duguay and Rentzepis (1967)). In the case of the Na II laser, McGuire and Duguay (1977) have examined experimentally the production of and pumping with X-rays. In this case the flashlamp and filter system employed was not in such close proximity to the laser medium as discussed here. Axelrod (1977) simulated photoionization pumping in his analysis of the boron laser, although the intensity in the present case is lower by two orders of magnitude. The lithium laser examined by the Avco group required photoionizing radiation at lower energy (Mani et al (1978)).

The problem of X-ray production in the multi-keV range in laser produced plasmas is of interest to the spectroscopy community, although spectroscopic studies usually are not concerned with the absolute photon emission from a target, rather the relative intensities are noted. Molybdenum has recently been examined by spectroscopists as listed in the bibliography, of present interest is a study of the M-shell spectra by Mansfield et al (1979). The band structure of the M-N radiation produced by a molybdenum laser plasma has received comment by Kononov (1978). Molybdenum is important in the present context as it is the current choice for flashlamp material by virtue of its M-N radiation near 500 eV. Present interest in molybdenum is not directed towards spectroscopic issues as the main concern is the

production of strong X-ray fluxes over a relatively broad band of energy, namely several hundred electron volts. The model used in the LASNEX simulations is not a spectroscopic model (see Lokke and Grasberger (1977)), and the resulting calculations are by necessity somewhat crude. We have done some work towards the construction of a spectroscopic model; however, the amount of work involved is considerable. Details may appear elsewhere. Ideally one would irradiate the flashlamp and filter system and measure the spectrum produced as well as the broad band integrated flux in the 500 eV range. Much work has been done towards such an experiment, and one shot was performed in 1979; however, the experiment has yet to be successfully completed.

Much of this chapter is devoted to the problem of the generation of bandpassed X-radiation. In Section A we consider the selection of flashlamp and filter materials, including a discussion of the rationale behind the choice of molybdenum and LiF. Section B is devoted to the operation of the flashlamp and filter system while under intense infrared laser irradiation. This section includes a brief review of some of the issues involved in light absorption and discusses the irradiation as seen from numerical simulation studies. In Section C we touch on the suprathreshold electron problem as it relates to the gain medium. In the final section the laser kinetics and small signal laser gain are reviewed.

In the research conducted so far, the emphasis has been placed on the generation of population inversions and the creation of usefully large small signal gain. A more complete description would involve transport of the short wavelength laser radiation out of the laser medium and hopefully towards a suitable EUV or X-ray detector. The simulation of the laser transport involves a 3-D calculation, and to date a link between the two dimensional XRASER and a 3-D laser transport post processor code has yet to be completed. This work may be described elsewhere when finished. For purposes of this thesis we shall be content with a calculation of the small signal gain, and be pleased if the small signal gain is sufficiently high to lead to a gain-length product equal to or greater than 20. Under the conditions discussed here in fact the gain is calculated to be in excess of 50 cm^{-1} which satisfies the above criterion.

Some of the calculations discussed here were carried out before the 4 mm length constraint on SHIVA was appreciated. At the time this chapter was completed the benchmark target length was 5 mm, and so there is a minor problem with consistency. Most of the results would be nearly the same for a 4 mm target and so we have not altered the chapter.

One problem which is not addressed directly here has to do with beam inhomogeneity. All of the large high power laser systems which we have considered for experimental work (ARGUS, SHIVA, OMEGA, and the laser at NRL) have large (factor of 4 or 5) variations in beam

intensity at the focal plane. The present target design shows some sensitivity to pulse intensity variation, and we have considered two intensities ($5 \times 10^{13} \text{ W/cm}^2$ and 10^{14} W/cm^2) in the present chapter. If one specified a beam intensity of $7 \times 10^{13} \text{ W/cm}^2$, expecting excursions between $3 \times 10^{13} \text{ W/cm}^2$ and $1.5 \times 10^{14} \text{ W/cm}^2$, the present target designs may function sufficiently well to obtain lasing action. Excursions to $2 \times 10^{14} \text{ W/cm}^2$ would prove to make matters difficult since overdriving the flashlamp and filter system can cause the back wall to heat up prematurely, resulting in 2p ionization and subsequent pumping of loss instead of gain. Such issues may be discussed later on elsewhere.

A. Material Selection

The flashlamp material must radiate well in the 500 eV spectral region. The material selection is limited to elements with strong line emission in the desired range as bremsstrahlung and bound-free opacities are too small under conditions of interest. The strong $\Delta n=1$ emission seems to be most appropriate for the problem in which case there are K-, L-, M-, and N-shell lines from which to choose. Examination of the calculated X-ray emission from elements with $\Delta n=1$ emission close to 500 eV suggested that the M-shell was the best radiator; K-shell lines are too sparse to fill in the spectral region, L-shell lines occur over too small of a bandwidth, and N-shell radiation is spread over too large a bandwidth. Of the M-shell radiators with strong emission in the vicinity of 500 eV, molybdenum was chosen as it appeared to be the best known, although the optimum radiator probably has a nuclear charge less by two or three. Future work might include reexamination of this choice of flashlamp material. The question of whether a mixture of materials would radiate more brightly, or perhaps layering of elements of neighboring nuclear charge to optimize the radiative output was considered to be interesting but was deferred in view of the target fabrication difficulties such designs would entail.

Having selected a flashlamp material we turn to the problem of the bandpassed filter material. The most important function served by the filter is the absorption of the ruinous low energy radiation--the cross

section for 2p ionization is in excess of 10^{-18} cm^2 per electron near 50 eV while the cross section for 2s ionization is close to $2 \times 10^{-20} \text{ cm}^2$ per electron near 500 eV. The lasing scheme cannot tolerate low energy radiation passing into the neon gas and therefore some care must be taken in blocking it out. At the same time the filter must be optically thin near 500 eV lest the pumping radiation itself be absorbed. Pursual of the cold opacities of the elements suggests that lithium or lithium hydride are about the best one can do, however, these materials received vetoes from the target fabrication group who instead suggested Be or parylene. Parylene is unacceptable due to the K-shell absorption of one of its major constituents, carbon, and simulations showed beryllium to be marginal. An alternate material, LiF, was finally chosen and with the selection came the added advantage of fluorine K-shell filtering which is effective in blocking radiation which otherwise would cause K-shell ionization in neon. K-shell ionization in Ne I does not directly affect the Ne II inversion as the Auger process yields more highly stripped neon ions. Indirectly the Ne I reservoir is drained, thus less population is available for the Ne II inversion. The harmful effect of K-shell ionizing radiation is that it precludes an inversion on the 32.7 eV line in Ne III as the Auger decay preferentially populates the lower level ($2s^2 p^4 \text{ } ^1\text{D}$) of the transition. If the laser is pumped with sufficient intensity lasing may occur on the higher energy Ne III line if LiF is used as a filter.

The flashlamp thickness is only a few hundred Angstroms and the LiF thickness is about a micron in the flashlamp and filter system. The angle averaged transmission through one micron LiF is shown in Figure III-2, and examination of the transmission function shows that most of the radiation in the pumping band near 500 eV should transport through the filter. The optical depths of the various components of the filter are shown in Figure III-3, showing that at low energy, the lion's share of the absorption is done by the fluorine 2p electrons.

Neither member has sufficient mechanical strength to support an X-ray laser in cylindrical geometry, and so it was suggested by C. Hendricks to lay down a thin support layer of beryllium to add the requisite mechanical strength. The benchmark target design was therefore chosen to be three-layered, and such a target is considered here although in the course of time several variants on this design have been considered. In section E a target is considered in which the LiF filter is sandwiched between the parylene layers and the molybdenum flashlamp layer is deposited on the outer parylene layer.

The benchmark target dimensions were determined from a series of numerical simulations to be reasonable values. The design is illustrated in planar geometry in Figure III-4.

Figure III-2. Angle averaged transmission of 1μ LiF between 20 eV and 1 keV.

At one micron thickness the absorption of LiF is considerable below 1 keV. The window near 500 eV transmits nearly 40% of isotropic radiation passing through the LiF filter. The radiation produced by the molybdenum in the 500 eV range is due to M-N transitions of the molybdenum M-shell ions (11-28 electron systems). As much of the radiation is produced by optically thin lines, the angular distribution is not isotropic, rather the directions at large angles from normal are favored. The transmission of the filter to this radiation is less than shown in the figure, although the above figure is useful in showing the approximate transmission involved.

Figure III-3. Optical depth of constituents for a 1μ LiF filter.

At low photon energies the filter absorbs very strongly, with the main absorption being done by the fluorine. The weakly bound 2p electrons of fluorine absorb more at low energies than the 2s electrons due to the different characters of the absorption cross sections (similar to the difference in 2s and 2p cross sections of Ne I on which the laser scheme is based). The K-shell absorption of Li is seen to be weaker than the fluorine L-shell, due primarily to the numerical difference in electrons responsible for the absorption (2 1s electrons versus 7 2s and 2p electrons). The 2s electron of Li plays a very minor role in the filtering.

Figure III-4. Benchmark design for flashlamp and filter system displayed in planar geometry.

The benchmark target design consisted of three layers; molybdenum was the flashlamp material, LiF served as a bandpass filter and the beryllium was included for the mechanical strength which it would lend the target. The laser pulse is incident from the left. An experiment to check the X-ray spectra produced on the two sides of the target was suggested by A. Szoke. Such an experiment would be valuable as the X-rays which pass through the filter and backing in a cylindrical or "channel" target would be inaccessible to diagnostics. Unfortunately the planar target shown above may not be buildable as in the attempts at construction by the target fabrication group it was found that the beryllium would not stick to the LiF.

B. Flashlamp-Filter Operation

This section is concerned with the operation of the flashlamp-filter system while under irradiation by 1.06μ laser light at 10^{14} W/cm^2 . The absorption fraction for short (sub 100 psec) pulse irradiation should be close to the values measured in similar experiments--Maaswinkel et al (1979) report normal incidence reflection of about 66% from planar copper targets with 30 psec pulses at 10^{14} W/cm^2 . Shay et al (1978) report absorption fractions of 17%-40% on parylene and tungsten glass targets at 200 and 300 psec for 10^{14} W/cm^2 irradiation. Laser light absorption is dependent on incident angle, polarization, wavelength, target material, surface quality and pulsewidth as discussed in the references in the absorption section; however, for the purpose of X-ray laser design the figure 30% absorption at the critical surface (where most of the reflection occurs) has been adopted. This figure is the same as is discussed in Rosen et al (1979) in the discussion of 2-D ray tracing calculations.

At incident laser intensities near 10^{14} W/cm^2 most of the absorbed laser energy is imparted to the suprathermal electron distribution. The hot electrons are generated by plasma instabilities near the critical surface and at 10^{14} W/cm^2 the temperature which characterizes the distribution is about 2.8 keV for 1.06μ light (see Gitomer and Henderson (1979) or other references in the suprathermal section of the bibliography). Experiments monitoring K-shell fluorescence in

layered targets indicate that about 30% of this energy is deposited in the target (for 10^{15} W/cm² pulses of 100 psec duration as described by Hares et al (1979)), the remainder being observed as preheat and fast ions. Rosen et al (1979) gives similar figures in the case of gold disks. The fast ions are produced through acceleration by a strong electrostatic field set up by the outward streaming hot electrons. After the hot electrons are produced near the critical surface they wander throughout whatever space is energetically allowed to them. Nearly elastic scattering dominates the interaction of hot electrons with electrons and ions of the dense plasma of the target (hence hot electrons do not lose their energy in passing through or being scattered by the target). As the hot electrons stray into the coronal region they encounter the strong static field, losing kinetic energy as they climb up the potential well. Reaching the turning point they roll back down the well at the same velocity in the shifted frame (the outer coronal region expands outward at speeds in excess of 10^8 cm/sec and the electrons are Doppler shifted down in energy upon reflection). Hence the hot electrons can lose their energy through pushing against the electrostatic field. At 7.5×10^{14} W/cm² for a 70 psec incident pulse on a glass microsphere, Campbell et al (1977) report that 50% of the absorbed energy appears as fast ions.

There is another issue which deserves some comment here, namely, thermal electron transport inhibition. The gold disk analysis in Rosen et al (1979) includes a graphic illustration of the failure of classical electron transport in the ablation region at 8×10^{13} W/cm²

for a 1 nsec pulse; strong flux limiting is required to bring the simulated emission profiles into agreement with the experimental results. Similar results have been previously reported as listed in the references in the transport inhibition section. McClellan et al (1979) discuss a threshold for the onset of strong thermal transport inhibition as a function of target Z, and inspection of the results indicate that in the case of molybdenum the threshold is near 2×10^{14} W/cm². Suprathermal electron transport into a Mylar disk seems to be uninhibited at 10^{15} W/cm² as discussed by Hares et al (1979), indicating some controversy on this point. The calculations described here include transport inhibition where the thermal flux is limited to 0.02 of the mean electron thermal velocity.

Ponderomotive force causes additional profile steepening along with that produced by transport inhibition, however, as noted by Rosen et al (1979) it is negligible at 10^{14} W/cm². This effect begins to play a role at 3×10^{15} W/cm² and is therefore not of interest in the present discussion.

Non-LTE effects are important in the flashlamp, in particular at electron densities below about 10^{22} /cm³ the N-shell can no longer be kept in equilibrium against a half-filled M-shell core. It was observed in simulation that the degree of ionization can change appreciably from one spatial zone to the next where in the flashlamp region a spatial zone starts out as approximately 5 Å of solid matter. Fine zoning was required to obtain reasonable results in molybdenum, similar to the zoning issues described in Rose et al (1979) in the longer pulse gold disk simulations.

We now consider briefly the hydrodynamics of the flashlamp and filter system in planar geometry while under irradiation. The simulation was carried out using a laser fusion code (see the bibliography). The example under consideration is the benchmark target (see Figure III-2) which consists of a 200Å Mo flashlamp, a 1 μ LiF filter and a 3000Å Be backing. The incident laser pulse is a 90 psec FWHM gaussian pulse with peak intensity of 5×10^{13} W/cm² centered at 120 psec. The molybdenum is seen to expand and be pushed outward by the expansion of the LiF filter as is illustrated in Figure III-5. A shock wave is observed to drift majestically across the filter over the course of the problem, eventually entering the beryllium backing toward the close of the sequence.

In Figure III-6 is shown the electron temperatures of the spatial zones of the Lagrangian mesh of the three materials. The transient heating of the molybdenum zones as they travel from the cool dense inner region to the outer coronal region is observed to take a few dozen picoseconds when the incident laser pulse has reached its peak intensity. One might imagine that transient non-LTE effects would be important in determining the degree of ionization of the molybdenum as it is transported from dense 100 eV regions to under-dense 1500 eV regions. In the LiF filter a strong thermal temperature gradient is sustained over the course of the problem--the outer part of the filter serves to thermally insulate the inner part, allowing it to remain relatively cool (sub 10 eV) during the irradiation. The lithium K shell and fluorine L shell remain more or less intact

Figure III-5. Density of flashlamp and filter system as a function of space at six times.

The laser pulse is a 90 psec FWHM gaussian pulse of peak intensity $5 \times 10^{13} \text{ W/cm}^2$. The peak intensity occurs at 120 ps. The laser is incident on the right hand side. The spike at the right hand edge of the 10 ps figure is the molybdenum flashlamp, which is next to the LiF filter. On the left hand edge is the beryllium backing. The flashlamp is pushed outward over the duration of irradiation, and can be recognized as the spike at 3μ on the 190 ps figure. The expansion of the LiF region pushes the beryllium to the left, although the division between the regions is less obvious. A shock wave is generated at the right hand side of the target and can be observed moving slowly to the left over the course of the problem.

Figure III-6a. Electron temperatures of the molybdenum zones as a function of time.

The laser is incident on $200\text{\AA}$ of molybdenum which is the flashlamp part of the flashlamp and filter system discussed in the text. Below the flashlamp is 1μ LiF which acts as a heat sink for the inner edge of the molybdenum. The molybdenum region is finely zoned with each zone representing on the average about $6\text{\AA}$ of cold molybdenum. As the plasma heats up and expands, the Lagrangian mesh follows it such that the amount of mass between two mesh points is kept constant throughout the problem. The fluid and transport equations are integrated on the Lagrangian mesh and the result for the molybdenum flashlamp is shown above. One observes that different zones heat up from about 10 eV to 1 keV in a few dozen picoseconds. The calculation was performed on a laser fusion code (Zimmerman (1973)). The initial electron temperature under experimental conditions is room temperature (0.024 eV); however, in the calculation we have started the problem at 0.5 eV for numerical reasons. The difference in the hydrodynamics and heating rates due to this approximation are small.

Figure III-6b. Electron temperatures of the beryllium and LiF zones as a function of time.

The LiF filter behind the molybdenum flashlamp is 1μ thick and was zoned using 17 mesh points. The part of the filter near the flashlamp was finely zoned and the part near the backing was finely zoned. The filter is seen to sustain a large temperature gradient, with the outer part of the filter reaching a peak temperature of about 80 eV. The inner edge remains below 10 eV throughout the problem. The absorption of the filter is greatly reduced at high temperatures due to loss of the fluorine 2p electrons. Making the filter thicker would help to insulate the inner part of the filter allowing the filter system to remain operational for long pulses (~ 100 ps) at relatively high (10^{14} W/cm²) intensities. The beryllium remains cool (below 10 eV) for the duration of the problem.

in the cooler regions of the filter such that low energy filtering occurs. The beryllium remains below 10 eV during the problem.

The zone positions and velocities of the molybdenum flashlamp as a function of time are shown in the next figure (Figure III-7). The dense region of the molybdenum blows outward reaching a velocity of somewhat less than 10^7 cm/sec, while zones in the ablation region blow off much faster. From a hydrodynamics viewpoint, X-ray laser targets are boring in comparison with the high compression and shock physics of laser fusion targets.

We now turn to a discussion of the radiative properties of the flashlamp and filter system. In terms of the total energy balance, about 35-40% of the total energy absorbed is emitted in radiation at all frequencies. The conversion efficiency is similar to the 35% figure given for the gold disk analysis of Rosen et al (1979). Most of the energy is calculated to radiate in the direction of the ablation, however, about 15% of the total radiated energy transports through the LiF filter and Be backing and would pump neon gas if some were present. The calculation time integrated spectra in the reverse direction is shown in Figure III-8 and the temporal variation is shown in Figure III-9. The calculated values in the radiative energy balance are given in Table III-1.

Figure III-7. Positions and velocities of the molybdenum zones as a function of time.

The 35 zone positions and velocities of the molybdenum flashlamp in the calculation discussed in the text are shown here. One observes that the outermost zone blows out to about 50μ by 400 psec. The velocity of the outermost zone reaches $1.8 \times 10^8 \text{ cm/sec}$ although the dense part of the flashlamp moves much more slowly (near 10^7 cm/sec). As the laser hits a zone, the zone heats and moves rapidly outward into the outer coronal region. The zones of the dense region are so close together that they are not resolved on the figure, and the effect of the laser is to peel them off one by one.

Figure III-8. Calculated time-integrated spectra out the backside of the target.

One observes that the molybdenum M-N radiation near 500 eV is the dominant feature in the spectrum. Radiation below 350 eV has been filtered effectively by the LiF/Be filter system. The effect of filtering by the fluorine K-edge is seen in the depression of the spectrum near 750 eV. Around 1 keV is emission due to molybdenum M-O and higher shell transitions. Emission near 2.5 keV is due to L-shell transitions where structure is lost due to the gross zoning of the photon bins. The structure near 500 eV is an artifact of the spatial zoning and the hydrogenic approximations employed in the simulation.

Figure III-9. Power radiated out the backside of the target.

The emitted radiation follows closely the incident laser intensity at an overall efficiency of about 1.8% although the peak shown here is only about 1.5% of the $5 \times 10^{13} \text{ W/cm}^2$ peak laser intensity.

Table III-1 RADIATIVE ENERGY BALANCE

Peak Intensity W/cm^2	Incident Energy J/cm^2	Absorbed Energy J/cm^2	Energy Out (front) J/cm^2	Energy Out (back) J/cm^2	Fractions of Absorbed Energy		
					% (Front)	% (Back)	% (Total)
5×10^{13}	4790	1440	459	84	34.4	5.8	40.2
10^{14}	9570	2870	807	182	28.8	6.3	34.5

C. Suprathermal Electrons Within the Gain Medium and Ion Heating

Resonance absorption of intense laser radiation at the critical surface results in the production of hot electrons as discussed briefly in the previous section. At 10^{14} W/cm² we have used a Maxwellian at 2800 eV to describe the hot electron distribution. The hot electrons are attenuated (both elastic and inelastic scattering cut down the fast electron flux) as they transport through the flashlamp and filter system. Initially none of the hot electrons make it through due to the low intensity and hence low hot electron temperature. As the pulse reaches its peak value much attenuation still occurs; however, not many hot electrons are required to displace the 10^{16} /cm³ electrons which are created through photopumping the laser medium.

We have calculated the hot electron spectrum in the case of a flashlamp and filter system in close proximity to the laser medium using LASNEX, and find that at the time of maximum pump intensity the hot electron density at the backside of the target exceeds 10^{17} /cm³. Having transported through the flashlamp and filter material we assume that the hot electrons displace the thermal electrons completely. Electrons produced by photo- or collisional ionization leave the interior of the laser in favor of the hotter suprathermal electrons. The question may be raised as to what effect this process has on the laser gain, which is the topic of interest in what follows.

The hot electron temperature is near 2.8 keV at 10^{14} W/cm², and the mean thermal velocity is slightly higher than a tenth light speed, or about 35μ/psec. Traversal of the laser target takes about 10 psec. Thermal electrons travel more slowly; however, the factor is about two in speed for the photoejected electrons. Electrons produced collisionally may have trouble transporting out unless they are accelerated electrostatically. Under these conditions one might envision that electron-electron or electron-ion plasma instabilities might develop. The former process is not worrisome, as the laser scheme benefits from having a high electron temperature; however, electron-ion heating would cause problems if it occurred as ion Doppler broadening (due to room temperature thermal motion) is a dominant broadening mechanism. If the ions were to heat up to 1 eV, the gain would decrease by a factor of nearly $\sqrt{40}$. We have examined growth rates for electron-ion instabilities under conditions which we believe appropriate for the laser medium and have found that if the plasma were collisionless, the instability growth rate is too small by several orders of magnitude to be important. Under conditions which are unlikely although physically allowed (assuming electron densities as high as 10^{17} /cm³ and drift velocities of 10^9 cm/sec) the instabilities would probably not be important if the plasma were collisionless. Inclusion of ion-ion collisions into the formulation results in considerable damping of the plasma waves and the net result is that the laser medium is very stable against electron-ion instabilities, and the

ion temperature of the fluorine-like ions should not stray greatly from their initial room temperature value. Hydrodynamic work on the ions has been calculated and found to be negligible. The only mechanism which is likely to be effective is "proximity" heating, which may lead to about a 30% increase in ion temperature under the conditions of interest as will be described later.

D. Laser Kinetics and Gain

In previous sections we discussed the target design and simulation. Here we use the spectra from these calculations to drive kinetics calculations to obtain estimates of the laser gain. The photoionization of Ne I at 500 eV is a small effect under the conditions discussed so far, and the rates on a per electron basis for a neon atom sandwiched between two radiating planes of flashlamp and filter material of infinite extent are shown in Figure III-10 for incident 90 psec Gaussian pulses of $5 \times 10^{13} \text{ W/cm}^2$ and 10^{14} W/cm^2 intensity. The peak ionization rate of 2s electrons is seen to be $3.5 \times 10^8 \text{ sec}^{-1}$ and $8 \times 10^8 \text{ sec}^{-1}$ per electron at the two pumping intensities. One observes that early on in the problem the 2s electrons are ejected preferentially, and as time goes on 2p electron photoionization dominates. This effect is not due to radiation from the flashlamp leaking in below the filter edges; rather, it is caused by thermal emission from the inner edge of the flashlamp filter system. In the low intensity case the inner wall heats to about 9 eV and in the high intensity case the temperature reaches 16 eV (the temperatures of the Be zones are shown in Figure III-11). The 2p cross section is large at low photon energies and responds strongly to 10 eV and 15 eV radiation, while the 2s cross section is zero below the threshold at 49 eV (see Chapter II).

The population kinetics are analyzed using a rate equation code (XRASER). The atomic physics parameters described in Chapter II were used for the calculations discussed here. The overall design

Figure III-10. Photoionization rates per electron at $5 \times 10^{13} \text{ W/cm}^2$ and 10^{14} W/cm^2 .

The photoionization rates for the 2s electron is dominated by the 500 eV photon contribution and hence follows the pumping pulse closely. The 2p electron rate in the lower intensity case follows the 2s rate at half the value until 150 ps when the blackbody emission from the inner edge of the filter becomes important. The story is similar at the higher intensity except that since the filter heats up more the effect is more pronounced.

Figure III-11. Electron temperatures of the Be zones.

The temperatures of the zones of the Be region on the inner edge of the LiF filter are shown here. The edge near the LiF heats first and in the ensuing tens of picoseconds the entire layer is shocked to higher temperatures. The rise from 5 eV to 10 eV is the heating of interest at $5 \times 10^{13} \text{ W/cm}^2$. At 10^{14} W/cm^2 the shock heats the layer faster bringing it close to 16 eV. The initial electron temperature in the experiment is room temperature (0.024 eV); however, in the calculation we started the target at 0.5 eV for numerical reasons.

procedure works as shown in Figure III-12, namely, the target specification is analyzed using a laser plasma code to determine relevant spectral or hydrodynamic quantities. This data along with relevant atomic physics parameters are fed into the XRASER post processor for further simulation. XRASER computes populations, transfers line radiation and continuum radiation, estimates plasma parameters, and does 2-D hydrodynamics calculations of interest to x-ray laser problems. Here only the kinetics routines are being exercised. Detailed documentation of the numerical issues involved will appear elsewhere and we shall therefore limit our discussion to physics issues here.

We consider the results of a kinetics calculation of a laser design using the benchmark flashlamp and filter system as described in section B. The target has a rectangular cross section of dimensions $300\ \mu$ by $500\ \mu$, with a length of 5 mm. The target is irradiated uniformly along the top and bottom of the laser with no irradiation of the sides. Although this type of target is not optimal in many respects (for example, in cylindrical geometry a test atom in the center of the laser would see about twice as much pumping radiation as in the rectangular geometry) it is one which appears to be buildable. Uniform irradiation of the top and bottom of the target is not possible on the SHIVA laser system, rather ten circular spots of $500\ \mu$ diameter will be lined up along the length of the laser. The pumping strength will vary along the length of the laser, and therefore the assumption of uniform irradiation is an idealization.

Figure III-12. Schematic of code system used in X-ray laser target design.

This figure is intended to give some idea of the overall design system. The idea was to use the existing hydrodynamics and laser fusion codes (for example see Zimmerman (1973)) to model the flashlamp heating, shock waves, plasma expansion and supra-thermal electron generation. The atomic physics models in the hydro and laser fusion codes is not suitable for X-ray laser calculations, therefore the models used in the kinetics calculation were constructed using data from many references and from calculations where data was not available in the literature. One advantage of this approach is that as more is learned about the atomic processes in different ionic systems the incorporation of the data into the kinetics models is trivial. Models can be incorporated from many contributors in a convenient fashion.

Due to passive aberrations within SHIVA's optical system, the pumping pulses are nonuniform in space with factor of five variations in intensity occurring over the beam profile. It is therefore reasonable to examine the target operation at two laser pumping intensities in order to comment on differences in target performance which will accompany the spatially nonuniform irradiation. The pumping pulse is assumed to have a 90 psec width (FWHM) which is far from optimal, rather the figure was selected due to the physical limitations of SHIVA's pulse generation capability (90 psec is the shortest available pulse).

Under the conditions described above the population kinetics were computed assuming a neon density of $1.5 \times 10^{17} / \text{cm}^3$ (4.7 torr). This density is somewhat high of the optimum however it is useful for the purposes of the present discussion. The photon field was computed in a quadrant of the cross sectional rectangle (the problem has four-fold symmetry) using ray tracing on the results of the hydrodynamics calculations described in section B. The atomic processes of the neon-like, fluorine-like and oxygen-like isoelectronic sequences were included in the kinetics. The gain over the laser cross section at the time of peak gain (130 ps, where the peak of the laser pumping pulse occurs at 135 ps) is shown in Figure III-13. The electron temperature was pinned at 2 keV for the duration of the problem as an approximation (albeit crude) of the displacement of the thermal electrons by the suprathermal electrons. The variation of the gain as a function of this assumed fixed temperature is weak ($\sim 30\%$) as will

Gain (cm^{-1})

A : 2.16E+01
 B : 2.37E+01
 C : 2.58E+01
 D : 2.79E+01
 E : 2.99E+01
 F : 3.20E+01
 G : 3.41E+01
 H : 3.62E+01
 I : 3.83E+01
 J : 4.04E+01

Figure III-13. Gain on the 460.7 Å line at $5 \times 10^{13} \text{ W/cm}^2$.

In this figure and following figures we examine some of the kinetics issues as a function of space and time. The target considered has a rectangular cross section with dimensions 300μ by 500μ . The extension of the target in the third dimension was taken to be infinite in the calculations, however in reality the laser length is about 5 mm . Irradiation occurs on the top and bottom of the target, and the pulse used was a 90 ps Gaussian of peak intensity $5 \times 10^{13} \text{ W/cm}^2$. The flashlamp and filter system used was the benchmark target ($200 \text{ Å Mo/1 } \mu \text{ LiF/0.3 } \mu \text{ Be}$). The peak gain here is 40.4 cm^{-1} (this is the small signal value, assuming lasing does not occur) and one observes that the gain is nearly flat over a rectangular roughly 200μ by 250μ .

be examined later on. One observes a factor of roughly two variation of the gain over the cross section with the peak occurring at the laser center and minima occurring at the centers of the boundaries near the side walls which are not pumped.

The photoionization rates for ionization of Ne I to the Ne II $^2P_{3/2}$ and $^2S_{1/2}$ states are shown in Figure III-14. The Ne I L shell has eight electrons, six 2p and two 2s electrons. In terms of per electron rates, pumping into the $^2P_{3/2}$ state accounts for four of the 2p electrons and pumping into the $^2P_{1/2}$ state accounts for the other two. The rates shown in Figure III-14 are therefore results for four 2p electrons and two 2s electrons. One observes that 2s electrons are being photoejected at more than twice the rate of 2p electrons at the time of peak gain, as expected from perusal of the time dependent rates of Figure III-10. The fact that the two rates have nearly identical spatial dependences is attributable to the fact that nearly all of the radiation involved is high energy (~ 500 eV radiation) which is directed radiation by virtue of having been filtered. The pumping rates due to isotropic radiation have a very different spatial dependence, namely they would look like saddles in a 3-d projection whereas the present results would look more like mounds.

The electron density at the center of the target is illustrated in Figure III-15 as a function of time. Of interest primarily is the magnitude of the density (several times $10^{15}/\text{cm}^3$ at time of peak gain) and the temporal variation. The collisional ionization rate of 2p electrons from Ne I is about $4 \times 10^8 \text{ sec}^{-1}$ at 130 psec which is

Figure III-14. Photoionization rates (sec^{-1}) at the time of peak gain at $5 \times 10^{13} \text{ W/cm}^2$.

The photoionization rates are for Ne I $2s^2p^6$ -Ne II $2s^2p^5 \ 2p_{3/2}$ (top) and Ne I $2s^2p^6$ -Ne II $2s^2p^5 \ 2s_{1/2}$ (bottom) under the same conditions and at the same time as Figure III-13. The peak rates equate to per electron values of $6.9 \times 10^7 \text{ sec}^{-1}$ per 2p electron and $1.5 \times 10^8 \text{ sec}^{-1}$ per 2s electron. One is still ionizing 2s electrons more rapidly than 2p electrons by more than a factor of two, and yet after this time the gain begins to fall. This is due to electron collisions which are beginning to make their presence felt as the dominant electron collisional process is the removal of a 2p electron from Ne I as discussed in Chapter II. One observes that the peak pumping rates are close to half of the infinite sandwich results of Figure III-10 (the time is 130 ps).

Figure III-15. Electron density as a function of time at the target center.

The initial neon density is $1.5 \times 10^{17}/\text{cm}^3$ in this calculation and one observes that about 6% of the atoms have been ionized by 145 ps. The electron collisional avalanche is well underway, dumping population into the lower laser states through collisional ionization of Ne I. The electron temperature used was 2 keV, which is roughly the hot electron temperature at $5 \times 10^{13} \text{ W/cm}^2$, where the assumption has been made that suprathermal electrons completely displace the thermal electrons. Although the electrons will be cooler at early times, the use of an artificially high value for the electron temperature should not greatly affect the results as the electron densities are low, and as the collisional rate coefficient is a slow function of temperature as discussed in Chapter II.

comparable to the photopumping rates. In Figure III-16 are shown the fractional populations of the neon-like, fluorine-like and oxygen-like isoelectronic sequences and the fractional populations of the Ne II low-lying states. About 6% of the Ne I atoms have been ionized by 145 psec. The upper laser state Ne II $2S_{1/2}$ is seen to contain about 1% of the total population at the time of peak gain. The fractional inversion density is less by about a factor of seven.

The average energy of the ionizing X-rays which are effective in pumping an inversion can be estimated by examining Figure III-17, wherein is shown the average energy of photoejected electrons resulting from photoionization of Ne I. Since the ionization potential of 2s electrons is near 50 eV, one notes that 2s electrons respond on the average to photons of near 600 eV energy at 80 psec and 500 eV near 140 psec. The average energy of photoejected 2p electrons is very nearly the same, which is a statement about the narrow band nature of the pumping flux.

Figure III-18 shows the small signal gain as a function of time in the target center. The peak gain is near 40 cm^{-1} as calculated here (Due to the low oscillator strength the gain should be corrected upwards by 10%--see Chapter II, Section 8. This correction applies to all gains quoted in this chapter and we shall not mention it again). The peak gain is sufficient to permit 20 gain lengths in 5 mm.

Figure II-16. Fractional populations as a function of time at the target center.

The fractional populations of the neon-like, fluorine-like and oxygen-like sequences are shown at the left. One is reminded that photoionization of the Ne I L-shell by 500 eV radiation is at best a weak effect as most of the ($\sim 94\%$) remains in the Ne I ground state by 145 psec. On the right hand side is shown the low lying Ne II state fractional populations. The population of the upper laser state is about 1% of the total population (i.e., $1.5 \times 10^{15}/\text{cm}^3$) at the time of peak gain.

Figure III-17. Average energy of photoejected electrons as a function of time at the target center.

In this figure is shown the average energy (defined by

$$\langle \epsilon \rangle = \frac{\int_{\Delta\epsilon}^{\infty} S(\epsilon)\sigma(\epsilon)(\epsilon - \Delta\epsilon) d\epsilon}{\int_{\Delta\epsilon}^{\infty} S(\epsilon)\sigma(\epsilon) d\epsilon}$$

where $\langle \epsilon \rangle$ is the average energy, $\Delta\epsilon$ is the ionization energy, $S(\epsilon)$ is the photon flux (photons/cm²/sec/eV) and $\sigma(\epsilon)$ is the cross section) of photoejected electrons from Ne I. One notes that 2s and 2p electrons are ionized by radiation of nearly the same energy, indicating that the filter has remained cool relative to what is required for thermal emission to be a small effect. At 5×10^{13} W/cm² it is electron collisions which spoil the gain rather than thermal radiation from the inner filter edge.

Figure III-18. Small signal gain as a function of time at the target center.

One observes that the gain is high for about 35 psec, which is more than sufficient time for a pulse to travel the length of the laser (light can go 5 mm in 15 psec). The number of gain lengths at the time of peak gain is about 20 which should be sufficient with a margin of safety for one pass amplification of noise to reach saturation values. Strictly speaking, this laser does not saturate in the normal sense, rather it is quenched as stimulated emission dumps the upper state population into the self terminating lower laser levels.

In the next set of figures we examine the situation at 10^{14} W/cm². Photoionization rates are shown in Figure III-19 and in Figure III-20 we plot the fractional sequence and level populations in time at the target center. In this case the ionization proceeds further, by 155 psec, about 15% of the Ne I atoms are ionized. Figure III-21 displays the small signal gain, which is calculated to exceed 80 cm⁻¹. The kinetics are affected in two ways relative to the 5×10^{13} W/cm² example; on the one hand the stronger ionization causes more electron production which brings on the electron collisional avalanche somewhat earlier, and on the other hand heating of the inner wall causes increased 2p ionization by thermal emission. A plot of photoejected electron energy (analogous to Figure III-17 but for the higher intensity case) shows significant divergence for the 2s and 2p electrons as may be expected from the rates shown in Figure III-10, although we do not include the plot here.

In the following set of figures we examine the optimization of small signal gain as a function of initial neon gas density. At 5×10^{13} W/cm² the optimum density occurs at 10^{17} /cm³ (see Figure III-22), and at 10^{14} W/cm² the optimum is higher by 30% (see Figure III-23). At low neon density there is simply not enough neon to produce high gain, and at high neon density, collisional stripping quenches the gain early on. The optimum covers a range of about a factor of two in initial neon density. These studies are of interest in terms of providing accurate data on tolerances for the gas pressure in the

Rate (sec^{-1})

A : 3.71E+08
 B : 4.05E+08
 C : 4.39E+08
 D : 4.73E+08
 E : 5.07E+08
 F : 5.41E+08
 G : 5.75E+08
 H : 6.09E+08
 I : 6.43E+08
 J : 6.77E+08

A : 3.29E+08
 B : 3.60E+08
 C : 3.92E+08
 D : 4.24E+08
 E : 4.55E+08
 F : 4.87E+08
 G : 5.18E+08
 H : 5.50E+08
 I : 5.81E+08
 J : 6.13E+08

Figure III-19. Photoionization rates (sec^{-1}) at 10^{14} W/cm^2 .

The photoionization rates for 2p ionization (top) of Ne I to the $2p_{3/2}$ and for 2s ionization (bottom) into the $2s_{1/2}$ level are shown as a function of space at the time of peak gain (as in Figure III-14). The magnitudes of the rates are nearly twice that shown for the $5 \times 10^{13} \text{ W/cm}^2$ case, however, the spatial variation is different. The 500 eV radiation is directional in that rays travelling normal to the filter see less attenuation than rays travelling at wide angles to normal. The thermal radiation produced at the inner filter edge is much closer to being isotropic, and here accounts for a larger percentage of the pump rate than in the low intensity case shown in Figure III-14.

Figure III-20. Fractional level populations at the target center at 10^{14} W/cm².

The stripping of the neon proceeds further by roughly a factor of two from the 5×10^{13} W/cm² case (Figure III-16), and by 150 psec about 10% of the population is fluorine-like. Although the magnitude of the fluorine-like population is several percent ($\sim 5\%$) at 130 psec, the peak fractional inversion density at that time is much less ($\sim 0.3\%$). One notes from the figure on the right that by 140 psec the inversion is quenched, and examination of Figure III-10 will indicate the mechanism--namely thermal emission from the inner Be wall causes strong 2p ionization at this time.

Figure III-21. Small signal gain at the target center at 10^{14} W/cm².

The peak gain is about a factor of two larger than the corresponding case shown in Figure III-18 for 5×10^{13} W/cm² intensity. The duration of the gain in this case is about 20 psec as quenching drives it down near 140 psec. One unfortunate aspect of the scheme is thus illustrated, namely the potential for large negative gain. This could be a real problem experimentally, since if one of the beams irradiating the target is early and is too intense, it could pump a region to large negative gain and cause no laser emission to be observed.

Figure III-22. Optimization of the small signal gain as a function of neon density at $5 \times 10^{13} \text{ W/cm}^2$.

In this calculation the time histories of the populations at several values of the initial neon density were computed and the maximum gain occurring during the problem was plotted as a function of the initial ion density (left). The time histories of the gain are shown on the right for the different initial densities. One observes that the optimal density at $5 \times 10^{13} \text{ W/cm}^2$ is at $10^{17}/\text{cm}^3$ with the peak gain reaching nearly 45 cm^{-1} . From the plot on the right one notes that at high initial densities collisions quench the gain early precluding high values of gain. At low initial densities the peak gain occurs later on; however, the peak value reached is lower due to the small inversion densities achieved.

Figure III-23. Optimization of the small signal gain as a function of neon density at 10^{14} W/cm².

This calculation is similar to that of Figure III-23 except using the photon data from a calculation at 10^{14} W/cm². The peak gain reached is 87 cm^{-1} which is obtained at neon densities near $1.3 \times 10^{17}/\text{cm}^3$. The time histories show that the peak gains are achieved about 10 psec earlier than at the lower intensity. From this analysis it appears that a tolerance of about $\pm 30\%$ on the neon gas density in the target is a reasonable requirement.

laser system, which in turn places a constraint on the pressure measurement system to be implemented.

In Figures III-24 and III-25 we examine the sensitivity of peak laser gain to ambient electron temperature. In the earlier calculations we have assumed a fixed value for the electron temperature and in these figures we explore the sensitivity of the peak gain achieved to the hot electron temperature assumed. The sensitivity is about 20% to 30% over the range of parameters examined. The minimum gain occurs roughly at the temperature corresponding to the peak of the Ne I 2p electron collisional removal rate. As the local electron temperature increases, so does the gain. From this, one concludes that penetration of the hot electrons into the laser medium and displacement of the thermal electrons, if it occurs, is beneficial to the laser gain because the gain increases with electron temperatures above 400 eV.

In the next figure (III-26) we present the results of a calculation of the Ne II ion temperature due to "proximity" heating; that is, heating which occurs when two atoms near one another become ionized, and the Coulomb repulsion accelerates them away from one another at velocities much higher than thermal velocities. In this case the effect is responsible for a 30% increase in ion temperature. In the next chapter we consider neon at higher density, pumped considerably more vigorously, and find that the ions can reach 1 eV through this process alone. The effect is important in both cases

Figure III-24. Sensitivity of gain to electron temperature at $5 \times 10^{13} \text{ W/cm}^2$.

Time histories of level populations were computed at several values of an electron temperature fixed in time. The figure on the left shows the maximum gain achieved as a function of the electron temperature. One finds that the variation is weak ($\sim 30\%$) and that the lowest gain is obtained at the temperature where the 2p collisional ionization cross section is largest. At electron temperatures which are higher than 1 keV (which is hopefully the case if the suprathermal electrons can displace the thermal electrons effectively) the peak gain achievable rises rapidly. The time histories involved are shown on the right, indicating that little time shift is associated with uncertainties in the assumed electron temperature. These calculations were performed at an initial density of $10^{17}/\text{cm}^3$.

Figure III-25. Sensitivity of gain to electron temperature at 10^{14} W/cm².

The results here are similar to those of Figure III-25 with the photon data from a 10^{14} W/cm² pulse calculation used rather than 5×10^{13} W/cm². The basic result is the same although the sensitivity is less ($\sim 20\%$). Collisional ionization of 2p electrons plays a smaller role in the determination of the gain than at the lower intensity. These results were computed at 10^{17} /cm³ initial neon density.

Figure III-26. Monte Carlo calculation of the Ne II ion temperature.

Using the spectra from the 10^{14} W/cm² target and taking the initial neon density to be 1.5×10^{17} /cm³ we have computed the ion temperature of the Ne II ions. Five hundred particles were tracked inside a sphere with reflecting boundary conditions. The ionization rates were computed using XRASER and fed into a Monte Carlo particle dynamics code. The initial transients are due to the small number of ionized ions involved. As time proceeds and more particles are stripped, the calculation becomes smoother and more accurate. One finds an increase of about 30% in the ion temperature due to "proximity" heating, that is, heating caused by the ionization of a particle next to one which already charged (the Coulomb repulsion between the particles results in non-thermal ion velocities which contribute to a higher ion temperature when thermalized). This effect contributes to higher ion temperatures the higher the initial neon density and places additional limits on the maximum gain achievable in the scheme.

because Doppler broadening is important in determination of the small signal laser gain. Gain calculations in this chapter were done not including the effect, hence the corrected answers will be lower by about 15% (this does not include the oscillator strength correction which would shift the answers up by 10%--the overall accuracy of all of these estimates is limited by the calculation of the pumping spectra which may be in error by a factor of two; hence, 5% or 10% errors due to other effects are probably not important).

We conclude the section with an example of a laser design intended to show some of what might be possible if some of the constraints imposed in the design procedure were relaxed. The benchmark design discussed earlier had a basic limitation that the suprathermal electrons created at the critical surface deposited energy causing heating of the back surface of the beryllium backing. Thermal emission from the inner wall caused 2p ionization from Ne I and destroyed the 2s-2p inversion in Ne II. Because of this effect the target does not behave well when irradiated at high intensities, which includes intensities significantly above 10^{14} W/cm². Use of short wavelength radiation for pumping the flashlamp would certainly help matters, as the suprathermal electron temperature and flux would be reduced. We have not pursued the topic. As shorter wavelength laser drivers become available, the topic should be of much interest.

Another factor limiting the EUV laser gain is the risetime of the X-ray pumping radiation, which we have constrained because 90 psec pulses are the shortest available on SHIVA. Other systems are not limited in this manner and so it is of interest to examine briefly the ramifications. We assume for this example that LiH is available for filter material, and that the flashlamp and filter is composed of a $275\text{\AA}$ of molybdenum and a 2.7μ layer of LiH. We shall not be concerned with mechanical support in this example. The laser dimensions are assumed to be $400\mu \times 150\mu \times 4\text{mm}$, which allows a larger fraction of the solid angle to be covered by flashlamp and filter material as seen by the laser medium than in the benchmark case considered in previous pages. The pumping pulse is assumed to have a Gaussian leading edge which rises with a 15 psec risetime to a peak value of $2.5 \times 10^{14} \text{ W/cm}^2$ at 37.5 psec, and remains at that peak value for 100 psec. The peak photoionization rate for 2s electrons reaches $4 \times 10^9 \text{ sec}^{-1}$ per electron in the limit that the flashlamp and filter system are of infinite extent (this compares with the value of $8 \times 10^8 \text{ sec}^{-1}$ per electron from the 10^{14} W/cm^2 benchmark case of Figure III-10).

In this case the ionization proceeds further, and we have calculated the kinetics keeping levels up to and including the carbon isoelectronic sequence. In Figure III-27 we show the fractional populations of the isoelectronic sequences, from which it is seen that the initial Ne I population is nearly depleted. The initial

Figure III-27 Fractional Isoelectronic Populations for LiH Target

The calculation was carried out in much the same way as in the case of the benchmark target. LASNEX was used to calculate the flashlamp and filter system hydrodynamic and radiative properties. Time-, energy- and angle-dependent results were used to compute the spectra at the center of the target ($400\mu \times 150\mu$ cross section), and this spectra was used in the XRASER kinetics. The calculation included the 49 sublevels of the 6- to 10- electron L-shell systems. The initial neon density is $5 \times 10^{17}/\text{cm}^3$ and the electron temperature was kept constant at 5 keV.

neon density chosen here is $5 \times 10^{17}/\text{cm}^3$. In Figure III-28 we show the gain resulting on lines in the fluorine-like, oxygen-like and nitrogen-like sequences, computed neglecting proximity heating. The increase in ion temperature due to proximity heating will probably cause a decrease in gain of a factor of between two and three from what is shown; however, the basic point of the example remains, namely that much higher gain on the Ne II lines are possible in principle and that inversions seem likely to occur in the more highly stripped ions if sufficient X-ray pump intensity can be supplied.

Figure III-28a Gain on the Ne II $2S_{1/2}-2P_{3/2}$ Line for the LiH Target

One advantage of having a fast risetime of the X-ray pumping flux is that the neon density can be higher and the gain is correspondingly higher by virtue of a larger numerical inversion density. Although we have not included it here, proximity heating will cause significant ion heating leading to ion temperatures probably in the neighborhood of 0.1 eV. In the present case electron-ion coupling causes heating which is small compared to room temperature, and therefore the corrected peak value may be near 400 cm^{-1} .

Figure III-28b Gain on the Ne III $2p^6\ ^1S_0$ - $2sp^5\ ^1P_1$ Transition at 23.47 eV

This transition is pumped by the removal of two 2s electrons from Ne I, and the dominant route for this is through single electron photoionization. In this calculation lasing on the Ne II transitions is neglected. Its effect would have been to reduce the upper state $2p^6\ ^1S_0$ population as some of the Ne II $2sp^6\ ^2S_{1/2}$ population would have been stimulated to the lower Ne II level. The lower Ne III $2sp^5\ ^1P_1$ state decays radiatively to the 1S_0 and 1D_2 states (Ne III $2s^2p^4$) which helps maintain an inversion. Radiation trapping of these transitions may be a factor on these transitions at late times and we have not included the effect.

Figure III-28c Gain on the Ne III $2s^2\ ^1P_1-2s^2\ 2p^4\ ^1D_2$ Transition at 32.69 eV

This transition is pumped by the removal of a 2s and a 2p electron. The lower laser state does not radiatively decay further and hence this transition is self-terminating. The difference between the kinetics of the two interesting Ne III transitions can be viewed through comparison of this figure and the previous one. Collisional and radiative population of the lower state in the above case cuts off the gain earlier and more sharply. This transition is the highest energy 2s-2p line of the oxygen-like sequence and eighth highest energy line of all neon 2s-2p lines. The highest energy neon 2s-2p line is in the nitrogen sequence at 34.65 eV.

Figure III-28d Gain on the Ne IV $2p^5 2p_{3/2}-2sp^4 2p_{3/2}$ Transition at 20.45 eV

Although lower in energy, this transition had the highest gain of the nitrogen-like lines in the calculation being discussed. The gain peaks later than in the case of the oxygen-like lines by about a dozen picoseconds, and later than the fluorine-like lines by two dozen picoseconds. The gain is calculated to be less than 20 cm^{-1} , and for this reason the line would be at best weak if the calculation at all modeled the experiment. The estimate of gain in the nitrogen-like sequence is likely to be underestimated due to the assumptions made about electron collisional stripping, and so the results shown above are not too accurate. Countering this is the neglect of proximity heating which will lower the gain from the computed value by a factor of about two.

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IV. Resonantly-Pumped Photoexcitation Schemes

The two previous chapters were concerned with a laser scheme involving photoionization as a pumping mechanism. The motivation for considering the laser originated from the possibility of achieving high gain and hence the development of a solid "proof of principle" experiment. One drawback of the approach is that it does not readily scale to higher energies. In this chapter we pursue an alternate approach in which photoexcitation serves as a pumping mechanism. Line radiation produced by a strong transition in a hot dense flashlamp plasma is used to resonantly excite an appropriate transition in a laser ion. For example, in the case of helium-like ions, one might use resonant L-M transition radiation to pump the $1s^2-1s3p\ ^1P$ transition of the helium-like system, pinning the $1s3p\ ^1P$ population to a value consistent with the intensity of the incident radiation field. If the electron density is sufficiently great then the $1s3s\ ^1S$ and $1s3d\ ^1D$ levels will be populated through electron collisional deexcitation. If the ion temperature were in excess of 50 eV, then proton collisions would be effective; however, we shall find that the ion temperature remains quite low, being heated only to a few electron volts at most by proximity heating and electron-ion collisional coupling. An infrared laser might be tuned to the $1s3p\ ^1P-1s3d\ ^1D$ transition energy to stimulate the transition, as will be discussed at the end of section E. If

further the plasma is optically thin enough on the $1s^2-1s2p\ ^1P$ transition for radiation trapping to be a small effect, then population inversions are established on the $1s2p\ ^1P-1s3s\ ^1S$ and $1s2p\ ^1P-1s3d\ ^1D$ transitions. The scheme is illustrated in Figure IV-1 in the case of helium-like neon.

This approach can be used for systems other than the one described above. In section D we shall consider $1s^2-1s4p\ ^1P$ pumping in helium-like Ne IX, leading to 4-3 inversions near 55 eV (see Figure I-4), in section E we examine the scheme described above, and in section F we present results concerning $1s-3p$ pumping in hydrogenic Ne X. The one- and two-electron systems were selected here for their simplicity, and furthermore, the atomic physics involved is relatively well known (which is not the case, for example, in nickel-like systems).

The use of resonant line pairs for pumping short wavelength lasers was suggested by Vinogradov et al (1975) who presented a rough quasistatic analysis of the inversion kinetics. The silicon-aluminum resonance was considered as an example. The pump plasma region was assumed to be near 900eV at an ion density of $5 \times 10^{22}\text{ cm}^{-3}$. The pumping field produced by the silicon $1s^2-1s2p$ transition was described by a modal photon density of 0.4 which is what is estimated by balancing the collisional and radiative rates between the $n=1$ and $n=2$ states under the above conditions. The resulting gain in the

Figure IV-1. Resonant photoexcitation pumping of helium-like Ne IX.

Line radiation at $11.547 \text{ \AA}$ causes pumping of the $1s^2$ - $1s3p \ ^1P$ transition in helium-like neon. Once excited to the $1s3p \ ^1P$ state, an ion radiatively decays most of the time back down to the $1s^2 \ ^1S$ ground state, however, transfer to the $1s3d \ ^1D$ state can occur through either electron or ion collisional de-excitation (the transition energy is 0.46 eV). Electron collisions can transfer the population between the states, although due to the high electron temperatures involved this process is accompanied by electron excitation of the M-shell electrons to the N-shell and higher shells, a process which is detrimental to the laser gain. If the ion temperature is high enough, ion collisions can cause the required excitation transfer without accompanying Rydberg excitation. The $1s3d \ ^1D$ state can be inverted with respect to the $1s2p \ ^1P$ state if radiation trapping does not inhibit the $1s2p \ ^1P$ decay. Similarly, the $1s3s \ ^1S$ state can be inverted with respect to the $1s2p \ ^1P$ state, although the gain is less, due to the lower transition oscillator strength.

aluminum plasma was estimated to be 2400 cm^{-1} . In Norton and Peacock (1975) the pump field was taken to be 0.19 photons/mode leading to a gain near 100 cm^{-1} on the helium-like C V 3^1D-2^1P line. In Apruzese et al (1978) the pump flux corresponds to 0.18 photons/mode and gains in excess of 10^3 cm^{-1} are calculated in aluminum.

As may be expected the small signal gain is roughly linear in the strength of the pumping field. Part of the problem in designing a short wavelength laser system of this variety is to provide conditions in which large pumping fields are obtained. In the present case we are concerned with pumping fields which are smaller and perhaps best described in units of 0.01 photons/mode. Some explanation of this may be in order, as the pump conditions assumed here are considerably different from those assumed in earlier works. In the above mentioned works the focus has been on pumping with helium-like resonance lines. As is discussed in section A, there are very few K-shell/K-shell resonances at low Z, the only precise one being the sodium-neon resonance. The silicon-aluminum and carbon-carbon resonances described in the above mentioned works are not precise and require either opacity broadening or large bulk Doppler shift to cause radiation to spill sufficiently far out into the line wings to pump effectively. Bhagavatula (1980) has discussed a way to find more precise resonances using hydrogenic and helium-like 1-2 and 2-4 lines. In the present case we have examined L-shell/K-shell resonances and have discovered many precise resonances for which neither opacity

broadening or bulk Doppler shift are required to cause effective pumping. The conversion of infrared light to an L-shell pump line is less efficient than a K-shell pump line. The assumption in the current effort is that at present simply the development of the lasers is of more interest than the optimization of their efficiency, though this may be of interest later.

In the present case we are interested in L-shell/K-shell resonances. Irradiation of a thin film flashlamp with a short pulse (90 psec) of medium intensity ($6 \times 10^{14} \text{ W/cm}^2$) leads to pump radiation which may be near 0.01 photons/mode. Due to the lack of relevant experimental results and due to the difficulty in obtaining accurate theoretical estimates, we cannot say at present precisely how large the pump intensity produced will be. Some of the issues will be briefly described in section B. We have used a pump intensity of 0.015 photons/mode for the calculations presented in the final three chapters. This intensity appears to be sufficient for pumping the 55 eV laser; however, the 151 eV laser requires more like 0.025 photons/mode (if an intense infrared laser were used to assist in the kinetics, the requirement would be slightly below 0.01 photons/mode), and the 189 eV laser needs in excess of 0.03 photons/mode. These estimates are for reaching a gain of 50 cm^{-1} in the optically thin limit, where such a large gain would result in 20 small signal gain lengths if the laser were 4 mm long (in agreement with the constraint imposed by the SHIVA optics). Our general approach to

the problem is to determine what pump intensity is required, and to explore some of the kinetics and target design issues involved assuming that 0.015 photons/mode is available.

There is also the issue of the solid angle subtended by the flashlamp and filter system as seen by the laser medium. Norton and Peacock (1975) have a small geometrical factor due to the geometry selected. Vinogradov et al (1975) and Bhagavatula (1980) have pump radiation incident from one side and Apruzese et al (1978) are in cylindrical geometry with maximum angular coverage. In the present two dimensional geometry the angular factor is about 80% maximum (pump radiation comes from about 80% of the "sky" as seen by the laser region).

Another assumption made in the papers discussed above has to do with collisions between the M-sublevels of the lasing ion. In the above works, the M shell was assumed to be statistically populated, and such can occur either through electron or ion collisions. Ion collisions are important when the ions are hot ($\geq$ several dozen eV, but this will be discussed in section B). If the lasing material is initially a gas and if one works in the short pulse regime, the ions are heated neither by electron-ion collisions nor by nonisentropic shock heating (as can be arranged if the containing walls do not cave in on it) and remain cool (≈ 1 eV). Electron collisions are important at low electron temperatures (the rate coefficient varies as $T_e^{-1/2}$); however, in the present case the suprathermal electrons are assumed to

dominate the electron kinetics resulting in a multi-kilovolt electron temperature. These conditions impose themselves given the basic target design philosophy used here. Such conditions lead to strong nonthermal populations of the M-sublevels of highly stripped hydrogenic and helium-like ions of Z near 10. These conditions are significantly different from situations which have been previously examined.

In the case of 1-4 pumping schemes, one major limitation on the achievable small signal gain is that the pumping radiation is absorbed as it penetrates the laser medium. Hence for a given transverse dimension there is an upper limit on the laser ion density (and therefore the small signal gain) before the pump radiation fails to reach the interior of the laser region. In the case of 1-3 pumping schemes, radiation trapping of the 1-2 transition limits the maximum density and therefore the maximum gain achievable. In either case, the limitation would be significantly reduced if a velocity gradient could be set up to allow deeper pump penetration or to relieve the trapping. In the designs discussed here velocity gradients are set up through pressure gradients induced by non-uniform plasma stripping. For example, if the stripping of the laser plasma is driven in part through x-ray photo-pumping then different regions will strip down faster or slower depending on the local photon field. If part of the plasma lies in a shadowed region then the stripping is slower and the local electron density is less. Since the dominant contribution to the pressure is due to hot electrons, especially if they are suprathermal electrons, strong pressure gradients are obtained

which can drive velocity gradients which are quite large as measured in ion thermal velocity units. This effect has not been described previously in the short wavelength laser literature and appears to be quite useful in relieving trapping in laser designs of the type considered here.

The 55 eV laser scheme appears to have the least stringent requirements. The pump intensity required is near 0.015 photons/mode and the penetration length is about 25μ . The velocity gradients required to allow near uniform penetration are relatively mild. For the 151 eV laser, the requirements both on the resonant pump intensity (0.025 photons/mode) and on the magnitude of the velocity gradients (the $1s^2-1s2p\ ^1P$ trapping length is 3μ) are more stringent, and hence the laser design is more difficult. Optical stimulation of the $1s3p\ ^1P-1s3d\ ^1D$ transition improves matters, and can lead to a pump intensity requirement of less than 0.01 photon/mode. The 189 eV laser requires in excess of 0.03 photons/mode to reach the target gain of 50 cm^{-1} ; however, the requirements on the magnitude of the velocity gradients achieved are much reduced.

In section A we discuss briefly some of the spectroscopic issues involved, including selection of resonant lines and flashlamp performance. Section B is devoted to some of the atomic physics issues relevant to the modeling. In section C we review the hydrodynamics and radiation transfer problems encountered in design simulation. In the final three chapters we examine briefly target

design issues in the case of the 55 eV laser (section D), the 151 eV laser (section E) and the 189 eV laser (section F). In section D, we are concerned with two-dimensional target simulation as well as with the neon kinetics, and we explore the production of large velocity gradients. In the following sections, we use these results, and examine higher energy laser schemes assuming that a similar target configuration is used.

A. The Spectroscopy Problem

In this section we are concerned with the availability of resonant line pairs appropriate for X-ray laser schemes. Perusal of published line lists convinces one that there are few resonances between K-shell transitions of helium-like and hydrogen-like low Z systems, the exceptions being the pairs Na X $1s^2-1s2p\ ^1P$ (11.003Å)/Ne IX $1s^2-1s4p\ ^1P$ (11.000Å) and Si XIII $1s^2-1s2p$ (6.6481Å)/Al XII $1s^2-1s3p\ ^1P$ (6.6350Å) (Scofield (1975)). The sodium-neon resonance was reported by Vinogradov et al (1975) and represents a singular occurrence in the low Z K-shell spectra. The silicon-aluminum resonance is not so precise and the resonance defect must be made up either by broadening, Doppler shift or large number of optical depths. The use of opacity broadened lines for resonant pumping was suggested by Norton and Peacock (1975).

The sodium-neon resonance is likely to receive much attention in the coming years in the short-wavelength laser field due to the preciseness of the resonance and due to the systems involved. Pumping the $1s^2-1s4p\ ^1P$ transition in helium-like neon leads to 4-3 and 4-2 inversions, the former occurring near 55 eV and the latter occurring near 205eV.

Another approach to the problem of K-shell resonances is to use satellite lines arising from doubly excited states in resonant line excitation schemes (Norton and Peacock (1975)). This approach may provide

additional resonances of the K-shell variety although the details of a laser design using this approach have not yet been reported.

Having examined the K-shell transitions we next turn to the problem of L-shell transitions. Figure IV-2 shows the bands over which the K-, L-, and M-shell series are found, from which it is seen that there is a window of perhaps five ions in Z-space which might have an L-shell resonance with a line in a candidate lasing ion which is to be pumped. As an example we consider the situation of spectra in the vicinity of the helium-like K-shell lines of neon as shown in Figure IV-3. Examination of the figure shows that copper ($Z=29$) and zinc ($Z=30$) have no lines in the neighborhood of the $1s^2-1s2p^1P$ transition in neon at $13.447\text{\AA}$, and hence might be expected to be good candidates for flashlamp material if lines could be found at the $1s^2^1S-1s3p^1P$ wavelength ($11.547\text{\AA}$). As the $1s2p^1P$ state is the lower laser state, no good can come by resonantly pumping the $1s^2^1S-1s2p^1P$ transition. Unfortunately, no lines are reported in copper and zinc at $11.547\text{\AA}$. In the case of zinc the Ne-like $2s^2p^6-2s^2p^53s$ transitions are nearby, however, no resonance is observed. In the case of copper, the line of interest falls between the Ne-like $2p-3d$ lines, and below the $2p-3s$ lines of the fluorine sequence. It would appear that nickel would be a likely candidate as the $2p-3d$ arrays of the oxygen-like ion fall in the vicinity of the $11.547\text{\AA}$ lines. Recent spectroscopic work reported by Gordon et al (1980) has turned up a line at $11.539\text{\AA}$, although no lines appear to be at $11.547\text{\AA}$.

Figure IV-2. Line positions of the K-, L- and M-shell lines of the ions.

Examination of the figure shows that one reason for the relative lack of K-shell resonances is the narrow band in which the K-shell transition lie. Resonances can only occur between ions of neighboring charge at low Z . The L-shell spectrum is much richer and there are usually several systems which have lines in the vicinity of one which is to be pumped.

Figure IV-3. Spectra in the vicinity of the helium-like neon K-shell lines.

The illustration consists of the spectra of the elements between and including iron ($Z=26$) and zinc ($Z=30$) with some of the important arrays indicated. Not all lines have been included, rather the lines which had been reported as of late 1978 were used in construction of the spectra. The horizontal lines at $13.447\text{\AA}$, $11.547\text{\AA}$ and $11.000\text{\AA}$ mark the positions of the $1s^2-1s2p\ ^1P$, $1s^2-1s3p\ ^1P$ and $1s^2-1s4p\ ^1P$ lines of helium-like neon. Of interest in the selection of a pumping element for Ne IX is that a line be found resonant with the $1s^2-1s3p\ ^1P$ line at $11.547\text{\AA}$ and that no strong lines be found resonant with the $1s^2-1s2p\ ^1P$ line at $13.447\text{\AA}$. From the figure, one sees that use of $Z=20$ or $Z=30$ ensures the latter condition, however there are no strong lines found at $11.547\text{\AA}$ in these elements. The situation for $Z=26$, 27 and 28 is discussed in the text.

Other candidate radiator ions include iron and cobalt. In iron the spectrum near $11.547\text{\AA}$ begins to thin out as the boron sequence appears to be in the vicinity, however, a line is reported near $13.447\text{\AA}$, implying that pumping could occur on the $1s^2 1s-1s2p \ ^1P$ line in neon. For reasons to be discussed shortly, it would appear that cobalt would make the best radiator material for helium-like neon if a strong line is found within a few milli-angstroms of $11.547\text{\AA}$.

A similar analysis for fluorine yields a resonance with a fluorine-like iron line. The positions of the different spectra are shown in Figure IV-4. The spectra of chromium ($Z=24$) and manganese ($Z=25$) have not been reported in sufficient detail to determine whether good resonances exist. Examination of the spectra in the vicinity of the K-shell resonances of low Z H-like and He-like ions yields the resonances given in Table IV-1.

Before leaving the discussion of resonances, the question of resonances in higher isoelectronic sequences should be brought up. For example, why should the search for resonant lines be halted at the neon sequence, why, for example, could not nickel-like ions do just as well? At present in the field of spectroscopy the L-shell is being analyzed in the wavelength range of interest. Exploratory analyses of the simpler M-shell systems have been reported, and the N-shell appears virtually unknown. It may be that the M-shell will provide good resonances; however, much work remains to be done in this area.

Figure IV-4. Spectra in the vicinity of the helium-like fluorine K-shell lines.

The $1s^2-1s2p^1P$, $1s^2-1s3p^1P$ and $1s^2-1s4p^1P$ lines are indicated by the horizontal lines at 16.807Å, 14.458Å and 13.782Å, respectively. No lines have been reported in cobalt ($Z=27$) at 14.458Å, however, the fluorine-like $2p^5-2p^43d$ array of iron contains a weak line at 14.456Å. It seems likely that the nitrogen sequence of manganese ($Z=25$) will have a resonant line near 14.458Å, however, the spectra has not yet been reported in sufficient detail to know which lines are close. The spectra near the boron-like chromium arrays similarly is not yet known well enough to determine which if any lines occur at 14.458Å.

Table IV-1a Resonances with He-like $1s^2$ - $1s3p^1P$ Lines

<u>Lasant Line</u>	<u>Pumping Ion and Line</u>	<u>$\lambda(\text{\AA})$</u>	<u>gf</u>	<u>Reference</u>
CV 34.972				
N VI 24.900	F-like Ti $2s^2p^5 \ ^2P_{3/2} - 2s^2p^4(^2P)3s \ ^2P_{3/2}$	24.907	0.16	Bromage (1977)
	F-like Sc $2s^2p^5 \ ^2P_{1/2} - 2s^2p^4(^1D)3d \ ^2S_{1/2}$	24.899	0.58	Feldman (1973)
	N-like Ca $2s^2p^3 \ ^2P_{1/2} - 2s^2p^2(^3P)3d \ ^2D_{3/2}$	24.900	0.99	Feldman (1973)
	B-like K $2s^2p^3 \ ^2D_{5/2} - 2s^2p(^3P)3d \ ^2F_{7/2}$	24.889	4.30	Fawcett (1975)
O VII 18.628	O-like V $2s^2p^4 \ ^1D_2 - 2s^2p^3(^2D)3d \ ^1D_2$	18.630	0.60	Fawcett (1975)
	C-like Ti $2s^2p^2 \ ^1S_0 - 2s^2p \ 3d \ ^1P_1^0$	18.623	1.20	Bromage (1977)
	Li-like A $2p \ ^2P_{3/2} - 4d \ ^2D_{5/2}$	18.631	0.54	Edlen (1979)
F VIII 14.458	F-like Fe $2s^2p^5 \ ^2P_{1/2} - 2s^2p^4(^1D) \ ^2S_{1/2}$	14.456	0.17	Boiko (1978)
Ne IX 11.547	O-like Ni $2s^2p^4 \ ^3P_2 - 2s^2p^3(^4S)3d \ ^3D_3$	11.539	1.25	Gordon (1980)
	Li-like Mn $2s \ ^2S_{1/2} - 3p \ ^2P_{3/2}$	11.554	0.49	Edlen (1979)

Table IV-1b Resonances with H-like 1s-3p Lines

<u>Lasant Lines</u>	<u>Pumping Ion and Line</u>	<u>$\lambda(\text{\AA})$</u>	<u>gf</u>	<u>Reference</u>
C VI 28.465 28.466				
N VII 20.909 20.911	Li-like K $2s\ ^2S_{1/2}$ - $3p\ ^2P_{3/2}$	20.896	0.24	Edlen (1979)
O VIII 16.006 16.007	B-like V $2sp^2\ ^2D_{3/2}$ - $2sp\ (^3P)3d\ ^2F_{5/2}$	16.007	0.8	Fawcett (1974)
F IX 12.643 12.645	Ne-like Ni $2s^2p^6\ ^1S_0$ - $2s^2p^5(^2P_{1/2})3s\ ^1P_1$	12.654	0.83	Boiko (1978)
	Be-like Mn $2p^2\ ^1D_2$ - $2p3d\ ^1F_3$	12.643	4.60	Boiko (1977)
	Li-like Cr $2s\ ^2S_{1/2}$ - $3p\ ^2P_{1/2}$	12.656	0.25	Edlen (1979)
Ne X 10.239 10.240	F-like Ga $2s^2p^5\ ^2P_{1/2}$ - $2s^2p^4(^1D)3s^2D_{3/2}$	10.239	0.15	Hutcheon (1979)
	F-like Zn $2s^2p^5\ ^2P_{3/2}$ - $2s^2p^4(^3P)3d\ ^2D_{5/2}$	10.245	1.33	Hutcheon (1979)

TABLE IV-1c Resonances With He-like $1s^2-1s4p\ ^1P$ Lines

<u>Lasant Line</u>			<u>Pumping Ion and Line</u>	<u>$\lambda(\text{\AA})$</u>	<u>gf</u>	<u>Reference</u>
252	C V	33.426				
	N VI	23.771	O-like Sc $2s^2p^4\ ^3P_1 - 2s^2p^3\ 3d\ ^3D_2$	23.73	0.50	Fawcett (1975)
	O VII	17.768				
	F VIII	13.782	Be-like Cr $2p^2\ ^3P_{21} - 2p3d\ ^3D_3$	13.779	3.42	Spector (1980)
			Ne-like Ni $2s^2p^1\ ^1S_0 - 2s^2p^5(2p_{1/2})3s\ ^1P_1$	13.779	0.05	Gordon (1980)
	Ne IX	11.000	F-like Cu $2s^2p^5\ ^2P_{1/2} - 2s^2p^4(^1D)3d\ ^2P_{1/2}$	11.002	1.45	Gordon (1980)
			N-like Ni $2s^2p^3\ ^2D_{5/2} - 2s^2p^2(^3P)3d\ ^2F_{7/2}$	11.000	0.97	Fawcett (1975)
			He-like Na $1s^2\ ^1S_0 - 1s2p\ ^1P_1$	11.003	0.72	Scofield (1975)

In the reference section we have provided a short bibliography from which the line resonances were taken. We refer the reader to these for discussions regarding line measurements, accuracy and spectroscopic techniques. It should be noted that the literature has thus far been concerned with line positions and sometimes relative intensities, and absolute intensity information which is of interest here is not a quantity which is usually measured. This leads us to the next issue in our investigation, namely the line radiation intensity produced on the pumping lines.

If one pumps a laser target using SHIVA and employs 20 beams each with a peak power of one terawatt, illuminating upper and lower rectangular flashlamp regions with dimensions 4mm by 400 μ , then the peak intensity at the flashlamp is about 6×10^{14} W/cm². The optics on SHIVA do not allow pumping for a length greater than 4mm (K. Manes - private communication (1980)), and there are no cylindrical lenses available for line focusing. Target irradiation would occur through placement of ten circular illumination spots arranged with edges touching or slightly overlapping on each flashlamp and filter system. The shortest laser pulse available is 90 psec FWHM (P. M. Campbell - private communication (1980)) and there are currently no plans to provide shorter pulses during the remaining life of SHIVA. Shorter pulses or simply faster risetimes permits the flashlamp material to radiate with a higher modal photon density on a pumping line by virtue of a higher ion density, and would aid the target operation.

As the available parameter range is fairly narrowly defined, we next turn to the target itself. Perusal of Table IV-1 shows that most of the pump materials lie between $Z=20$ and $Z=30$ and we choose $Z=30$ (zinc) for the example discussed here. A reasonable flashlamp thickness given the pulse characteristics is $500\text{\AA}-1000\text{\AA}$; however, such a film would have minimal mechanical strength when faced with a few dozen torr of laser gas on one side and 10^{-5} torr vacuum of the target chamber on the other side. For mechanical strength we use parylene, the thickness being constrained by the requirement that the absorption near the pumping line not be too great. The resulting flashlamp and filter system is shown in Figure IV-5.

We have calculated the behavior of the target assuming irradiation by a 90 psec pulse with peak intensity of $6 \times 10^{14} \text{ W/cm}^2$ using LASNEX. The resulting temperature and density profiles are illustrated in Figure IV-6 at a time a few picoseconds after the peak of the pumping pulse. Initially the target was positioned to the right of the origin and the densities of the parylene and zinc were assumed to be 0.85 and 7.14 gm/cm^3 , respectively. The zinc is observed to have blown down a factor of about 30 from its initial solid density. There exists a substantial density and temperature gradient across the flashlamp and one notes that the inner part of the flashlamp is coolest which is an unfavorable condition as far as line pumping is concerned. The boundary between the parylene and the zinc is marked by the density

Figure IV-5 Flashlamp and Filter System for the Line Pumping Schemes

Zinc has a resonance with hydrogenic neon and was therefore chosen as flashlamp material for this target. Parylene is present as a mechanical aid as a thin zinc film would not by itself stand up against the differential pressures which will be applied. Although the analysis here is focused on zinc, the results apply approximately to neighboring Z. The target is to be irradiated with a 90ps FWHM pulse of peak intensity $6 \times 10^{14} \text{ W/cm}^2$. The thickness of 1000Å is a reasonable choice based on simulation results which show that there are probably fluorine-like ions present through most of the pulse. If the zinc film were thinner then the flashlamp would burn out too quickly, and if it were thicker then the heating would take longer.

Figure IV-6 Temperature and Density Profiles 10psec After the Pumping Pulse Maximum

The density of the zinc region is at best about 30 times below its initial solid density value (7.14 gm/cm^3), and the electron temperature varies between 150eV and a bit over 200eV where the densest part of the flashlamp is. The discontinuity in density near -1.7μ marks the parylene-zinc fluid interface, and arises since the pressure is continuous across the boundary. Within the following tens of picoseconds the flashlamp region warms up and blows down to the point where one no longer would expect to find fluorine-like ions. The operation of the flashlamp and filter system is a transient affair even as measured on the timescales of the pumping pulse width.

jump near -1.7μ . The temperature and density profiles vary in time and within a few tens of picoseconds the zinc will find itself well above 200 eV, and at monotonically decreasing densities.

We now turn to a brief discussion of the problem of pump line output from laser-irradiated flashlamp systems. The problem is difficult, and we have not yet completed enough of the solution in the case of the pump lines of interest here to determine conditions under which the requisite pump intensity can be produced. This aspect of the analysis is at present the weakest part of the design procedure. We have devoted much effort to this problem, and have carried out both LTE and crude non-LTE calculations of pump line emission (using the temperature and density profiles calculated by LASNEX, with kinetics and line transport done by XRASER). The problem in the present case is not accurately described by an LTE calculation, and an accurate determination of the emission intensity requires more elaborate kinetics models than those we have so far employed.

Given that the theoretical description used so far is not satisfactory, we may still discuss some of the relevant issues involved. In the final sections of the chapter, when we study the inversion kinetics in the cases of helium-like and hydrogen-like neon, we shall assume that the pump intensity is 0.015 photons/mode. This pump intensity is probably sufficient for pumping the 55 eV

laser; however, for the 151 eV laser, 0.025 photons/mode will be a minimum assuming that the velocity gradients produced are strong enough to relieve the radiation trapping problems which occur. If an infrared laser can assist in the kinetics, the pump requirement can drop by threefold. In the case of the 189 eV laser, the small signal gain is anemic at 0.015 photons/mode, and it is likely that in excess of 0.003 photons/mode will be necessary to reach our goal of 50 cm^{-1} . This goal for the small signal gain is based on the constraint that the laser length not exceed 4 mm, and were this constraint to be relaxed, the corresponding pump intensity requirement would be reduced (the reduction is approximately linear in inverse laser length).

In Figure IV-7 we show the modal photon densities which pump lines resonant with the hydrogen-like and helium-like neon K-shell lines considered in the following sections would have if they were in equilibrium at electron temperatures between 200 eV and 300 eV. In the case of a pump line at $11.547 \text{ \AA}$ (resonant with the Ne IX $1s^2 - 1s3p \text{ } ^1P$ line), 0.015 photons/mode would be produced by optically thick matter at 255 eV. If the temperature of the matter were less by only 10%, the resulting drop in photon intensity would be 40%. In the low temperature regime which we are considering, the equilibrium photon intensity is a strong function of the electron temperature. These arguments are not meant to imply that the flashlamp is in equilibrium for our problem (as it is not);

Figure IV-7 Modal Photon Densities of Pump Lines in Equilibrium as a Function of Electron Temperature

In the two limits of high electron density (sufficiently high to establish equilibrium between the upper and lower levels of the pump transition) and large optical depth (sufficiently large to provide a radiation field at the pump frequency whose temperature is the same as the emitting matter), a resonant pump line would produce $1/(\exp(\Delta E/kT)-1)$ photons/mode, where ΔE is the line energy and T is the electron temperature involved. We show above this function for pump lines resonant with three K-shell transitions in helium-like and hydrogenic neon. The equilibrium radiation field is observed to be a strong function of temperature in the interval shown, and that to achieve 0.015 photons/mode, temperatures of between 250 eV and 300 eV are required.

rather, that if the matter were in equilibrium, the requisite temperatures must be above 250 eV. For a non-equilibrium plasma the matter must be hotter in order to produce the same pump intensity.

Examination of the temperature and density profiles of the flashlamp (see Figure IV-6) reveals that the dense part of the flashlamp region is under 200 eV. At lower flashlamp densities (to the right of the dense region, which is to say, closer to the critical surface where laser absorption occurs), the electron temperature is higher. In the next few tens of picoseconds the dense region of the flashlamp becomes narrower as more of the zinc ablates outward, and the electron temperature rises. Ultimately, all of the flashlamp material ablates and there remains no dense region. The result is that for most of the problem there are regions where the electron temperature is sufficiently high that if a pump line were in equilibrium and optically thick, the requisite pump intensity could be produced. The problem is to get the flashlamp at both high density and high temperature at the same time, with the hope that the high density will bring about relative equilibrium between the pump levels.

There is another issue involved which is equally important. Simply having high temperature and high density does not assure one of obtaining strong pump radiation, even if the line is very strong (as measured by its oscillator strength). The ionization

balance of the different isoelectronic sequences is a major factor in the calculation, and that if there are strong temperature and density gradients present, the isoelectronic sequence of the pump line will occur over a relatively narrow parameter (and hence spatial) range. We have calculated pump line output under various assumptions, with the result that the isoelectronic sequence of a pump line is perhaps most important in determining how intense the pump radiation produced can be. The calculations were not of high accuracy; however, the results tended to suggest that pump lines of the neon-like sequence could produce about 0.003 photons/mode, that those of the fluorine-like sequence could produce 0.005-0.006 photons/mode and that those of the oxygen-like sequence could reach 0.01 photons/mode. The trend is toward lines of the more highly stripped sequences, given that one wants to maximize the modal photon density. The numbers quoted above are very approximate and depend on models and assumptions which require further examination. The accurate calculation of the non-LTE line emission in the more complex sequences is non-trivial, and unfortunately it is this problem which must be faced to solve the rather crucial problem of resonant pump intensity in target design (namely, whether there is enough resonant pump intensity given a target and infrared laser pump pulse). The calculation necessary appears to involve many levels (in excess of 100 for the L-shell and M-shell states of the sequences of interest), and although we have made some progress towards carrying out the

relevant modeling, we are not in a position to present results here. It is likely that $6 \times 10^{14} \text{ W/cm}^2$ is sufficient to achieve in excess of 0.01 photons/mode on some L-M lines of the oxygen-like and more highly ionized sequences, and probably represents a not unreasonable (and perhaps low) choice for incident laser intensity. We expect future work to shed further light on this problem.

The choice of neon as laser material was made for the analysis of this chapter as it represents what may be the high end of what seems reasonable to expect from this approach given present SHIVA capabilities and constraints. The major problem which limits the laser transition energy attainable is the production of sufficiently bright pump radiation (as measured in units of per cent modal photon densities). As discussed above, the use of pump lines in the more highly stripped sequences can lead to brighter pump radiation, given that the pump levels are in relative equilibrium. The relative degree of equilibrium to non-equilibrium is determined mostly by the ratio of the electron collisional excitation rate (which scales as $\bar{Z}^{-3}$, where $\bar{Z}$ is the effective charge of the pump ion) and the radiative decay rate (which scales as $\bar{Z}^4$ for strong resonance transitions). The condition for the establishment of relative equilibrium scales then as $\bar{Z}^7$, where the electron density is kept constant and the electron temperature is scaled as $\bar{Z}^2$. The ramifications of this for the present scheme is that the line pumping schemes involving lower energy pump lines will have much less severe pump laser constraints,

and that the constraints will be determined from an examination of the density achieved at the requisite electron temperature (to see if it is sufficient to establish approximate relative equilibrium between the pump levels). As matter tends to blow apart the more intensely it is irradiated, the production of hot and dense emitting regions becomes more difficult at higher laser intensities. Also, at the higher laser intensities, the temperature and density gradients become sharper, hence it becomes more difficult to have pump lines be optically thick.

The limitation of the present approach is due to the inability of laser irradiated flashlamps to produce bright pump radiation at high pump line energies. This limitation will be quantified more precisely in future work, and estimates we have made (though have not presented here) suggest that pump lines for hydrogenic and helium-like neon are close to the limit. For line pumping schemes involving higher energy pump lines, other schemes will be necessary to obtain bright pump radiation. Our interest has been directed here towards schemes involving direct near equilibrium pumping as the designs and experiments resulting are likely to be simpler than if, for example, radiation trapping were used to supply strong pump radiation.

B. Atomic Physics Parameters

Much work has been done on the development on non-LTE models describing plasma population kinetics. Early analyses of recombination in hydrogenic plasmas were published by Bates et al (1962) and McWhirter (1963), where the motivation was provided by astrophysics problems of interest at the time. The collisional rate coefficient used were prescriptions based on the scant theoretical and experimental results available. Since then the field has advanced considerably due to laboratory plasma experiments and astrophysical observations and due to the rapid development of theoretical and numerical expertise. In the bibliography at the chapter's end some of the relevant work has been referenced.

The kinetics model of interest here includes 82 levels running from neutral neon to the bare ion (see Table IV-2). The model is most detailed in the Li-like, He-like, and H-like sequences as is required for modeling the inversion kinetics. The energy levels were drawn from several sources: Erikson (1977) gives the hydrogenic energy levels and the uncertainties in the calculations for all Z . Accad et al (1971) have solved the two electron problems for S and P states of low Z ions, and Scofield (1975) has given accurate energy levels up to $Z=54$. Brown (1969) and Laughlin (1973) give results for D and F states. The energies

Table IV-2 Neon Levels Included Explicitly in Present Calculations

265	Ne-like	$2\ell^8$	B-like	$2\ell^3$	$3d_{3/2}$	$1s5\ell$	$1L$	$4s_{1/2}$
				$2\ell^2_{3\ell}$	$3d_{5/2}$	$1s6\ell$	$1L$	$4p_{1/2}$
				$2\ell^2_{4\ell}$	$4s_{1/2}$	$1s7\ell$	$1L$	$4p_{7/2}$
	F-like	$2\ell^7$		$2\ell^2_{5\ell}$	$4p_{1/2}$	$1s2s$	$3S$	$4d_{3/2}$
				$2\ell^2_{6\ell}$	$4p_{3/2}$	$1s2p$	$3P$	$4d_{5/2}$
				$2\ell^2_{7\ell}$	$4d_{3/2}$	$1s3s$	$3S$	$4f_{5/2}$
	O-like	$2\ell^6$			$4d_{5/2}$	$1s3p$	$3P$	$4f_{7/2}$
		$2\ell^5_{3\ell}$			$4f_{5/2}$	$1s3d$	$3D$	5ℓ
		$2\ell^5_{4\ell}$	Be-like	$2\ell^2$	$4f_{7/2}$	$1s4s$	$3S$	6ℓ
				$2\ell_{3\ell}$	5ℓ	$1s4p$	$3P$	7ℓ
				$2\ell_{4\ell}$	6ℓ	$1s4d$	$3D$	
				$2\ell_{5\ell}$	7ℓ	$1s4f$	$3F$	
	N-like	$2\ell^5$		$2\ell_{6\ell}$		$1s5\ell$	$3L$	
		$2\ell^4_{3\ell}$		$2\ell_{7\ell}$		$1s6\ell$	$3L$	
		$2\ell^4_{4\ell}$				$1s7\ell$	$3L$	
		$2\ell^4_{5\ell}$						
				He-like	$1s^2$	$1L$		
					$1s2s$	$1S$		
					$1s2p$	$1P$		
					$1s3s$	$1S$		
					$1s3p$	$1P$		
					$1s3d$	$1D$		
					$1s4s$	$1S$		
					$1s4p$	$1P$		
					$1s4d$	$1D$		
					$1s4f$	$1F$		
	C-like	$2\ell^4$	Li-like	$2s_{1/2}$		H-like	$1s_{1/2}$	
		$2\ell^3_{3\ell}$		$2p_{1/2}$			$2s_{1/2}$	
		$2\ell^3_{4\ell}$		$2p_{3/2}$			$2p_{1/2}$	
		$2\ell^3_{5\ell}$		$3s_{1/2}$			$3s_{1/2}$	
				$3p_{1/2}$			$3p_{1/2}$	
				$3p_{3/2}$			$3p_{3/2}$	
							$3d_{3/2}$	
							$3d_{5/2}$	

of Li-like ions are available from experiment and from calculations; results are given by Onello (1975), Martin and Wiese (1976).

Oscillator strengths for radiative transitions are known analytically for hydrogenic ions (see Bethe and Salpeter (1957)), and high accuracy results exist for the Li-like and He-like sequences; for example, Schiff et al (1971) give f values between S and P states. Other sources are listed in the references.

Electron collisional excitation of the hydrogenic sequence is an old problem which is still under active investigation. Some of the recent Coulomb-Born distorted wave and close coupling results have been listed in the references. The close coupling approximation involves the solution of what are sometimes termed the continuum Hartree Fock equations, and in the limit of an infinite number of states included in the 2-electron wavefunction expansion the result should become exact. The close coupling results should be the most accurate. The distorted wave approximation is simply a Born approximation using continuum wavefunctions which see the nuclear potential screened by the atomic electrons. The Coulomb-Born approximation is the Born approximation using Coulomb functions. The distorted wave approximation is oftentimes more accurate than the Coulomb-Born approximation. The last two methods are the workhorses of the field as they are relatively inexpensive to carry out and are often sufficiently accurate for

most purposes. Close coupling results are expensive as the computation involved is increased relative to the Born approximations and are not generally available. These approximations, as well as others, are discussed in much more detail in the review papers given in the references.

In the hydrogenic model under consideration we are interested in the collisional rate coefficients in j-j coupling. Electron collisional excitation from s states calculated in LS coupling can be adapted to j-j coupling trivially through multiplication by the appropriate statistical factor (the collisional cross sections are proportional to the statistical weight of the final state in this case as discussed in Burgess et al (1970)). Excitation from p and d states is not as simple and results do not appear in the literature for transitions other than the $2p_{1/2} - 2p_{3/2}$ (Burgess et al 1970)). The reason for this may be that for low Z systems under most conditions the LS description is adequate. In the present case there are two effects which can lead to selective population of the j sublevels. For example, the resonant photoexcitation process can pump a $1s_{1/2} - 3p_{1/2}$ transition more strongly than a $1s_{1/2} - 3p_{3/2}$ transition if the pumping line is narrow. The second effect is the ion collisional coupling which favors transitions between states of the same principal quantum number and the same j.

The j-j collisional cross sections were computed for neon in the Coulomb-Born approximation using a computer code written by

the author. The numerical techniques for the most part are standard and are described in the papers listed in the "numerical methods" section of the references. The mesh chosen for the integrations is of the form $x = \alpha r + \beta \ln r$, where r is the radial coordinate and α and β are constants chosen to optimize the accuracy of the integration. The Noumerov method was used for calculation of the free wave functions. The normalization of the continuum wave functions was carried out through analytic evaluation of the WKB expansion (a series solution in $1/r$ was assumed and the coefficient computed automatically). Belling's method (see Belling (1969)) for calculating the asymptotic contribution for large radius to the "direct" integral results in a divergent series and was found to fail in some of the cases of interest. The algorithm was replaced by an alternate and slower method involving quadrature on a logarithmic mesh appropriate for oscillatory functions. The transformation matrix for converting the LS matrix elements to j-j matrix elements was computed using the simplified formulae quoted in Sobelman (1972). The R-matrix elements were found for the most part to be much less than unity, and as the matrix inversions required for the CBO II approximation promised to be costly in terms of human time and computer time, it was replaced by a kludge which ensured that deviations from unitarity would be small. Matrix elements for large values of the total angular momentum were computed until they approached the

analytical results for the Coulomb-Bethe approximation in the dipole case and a corrected Born approximation for the higher multipoles (see Alder et al (1956) and Burgess et al (1970)). The numerical aspects of the code were checked through comparison with earlier published results and the agreement in the case of the H-like Coulomb-Born Oppenheimer I results was to all figures. The distorted wave approximation was examined and in the approximation used the improvement in the results were small. The primary effect that was missing was exchange in the computation of the free wave functions. The comparison of the results obtained for 1s-2s and 1s-2p excitation showed 10 percent agreement with the more accurate close coupling and DWPO results, which is similar to analyses reported previously (see Baluja and McDowell (1977) and Hayes and Seaton (1977)).

In the case of the helium-like ions there are results for many of the transitions of interest available, including some Coulomb-Born calculations in intermediate coupling (Sampson et al (1978)) and distorted wave results (Bhacia and Tempkin (1977)). In the present model we have used the accurate results where available and in the cases where no results have been given (primarily the M and N shells) the cross sections were computed in a distorted wave approximation, using single configuration Hartree Fock wave functions. The situation is similar in the case of the lithium-like ions.

For these three sequences (H-like, He-like and Li-like) the n-n dipole allowed cross sections were calculated using the classical path approximation (see the references in the classical path model section of the references). The model and calculations are similar to what was discussed in Chapter II. In Figure IV-8 are shown results for electron collisional excitation of the helium-like Ne IX singlet levels. The cross sections are observed to be monotonically decreasing in energy with approximate $\log(E)/E$ behavior, where E is the incident electron energy. The rate coefficients peak at low electron temperature and decrease asymptotically roughly as $T_e^{-1/2}$ where T_e is the electron temperature. As a practical matter, this behavior makes it hard to drive population from the $1s3p\ ^1P$ state to the $1s3d\ ^1D$ state (as is required to achieve high gain) if the electron temperature becomes very hot. For example, if the suprathermal electrons displace the thermal electrons, the resulting electron temperature can be several keV, in which case the n-n electron collision excitations rates are small.

In the case of the hydrogenic ion the coupling of the $3p_{3/2}$ state to the $3d_{3/2}$ and $3d_{5/2}$ states is of interest in the 3-2 line pumping schemes. The $3p_{3/2}$ state is very nearly degenerate with the $3d_{3/2}$ state ($\Delta E = 2.28 \times 10^{-4}$ eV), while the $3p_{3/2}$ state is relatively distinct from the $3d_{5/2}$ state ($\Delta E = 4.46 \times 10^{-2}$ eV). One might naively expect that with such a small energy difference, the electron collisional coupling between the $3p_{3/2}$ and $3d_{3/2}$ states

Figure IV-8. Electron collisional excitation cross sections of the n-n transitions of the singlet states in Ne IX.

The cross sections for the $1s2s\ ^1S-1s2p\ ^1P$, $1s3s\ ^1S-1s3p\ ^1P$ and $1s3d\ ^1D-1s3p\ ^1P$ transitions were calculated in the classical path approximation as discussed in Chapter II. The cross sections are found to be quite large at low energies and decrease monotonically with increasing electron energy (asymptotically they fall proportion to $1/nE/E$). The rate coefficients were obtained by integrating the cross sections over Maxwellians at different electron temperatures. Notable about these rate coefficients are that they decrease rapidly with increasing electron temperature. Hotter electrons couple the states of the same principle quantum number much less effectively than slow electrons. If suprathermal electrons dominate the electron distribution or heat the thermal electrons significantly, the electron collisional coupling between the M-shell states can be decreased dramatically.

might be favored. Calculations do not confirm this, rather the $3p_{3/2}$ to $3d_{5/2}$ cross section is larger by a factor of near 9, which is the ratio of the transition strengths involved. For electron collisions the energy difference between nearly degenerate levels does not play a large role, and the reasons for this have to do with a relatively weak unperturbed dependence as well as contributions from the cutoffs imposed by Debye screening and from the finite duration of the collisions (radiative decay and collisional de-excitation processes are dominantly responsible for the latter effect).

Ions can cause n-n dipole allowed transitions, and these are potentially very important for the 3-2 line pumping schemes. Ion collisions were computed for many of the transitions using the classical path approximation, and the general behavior is that the cross sections are quite small at low energies (due to Coulomb repulsion) and becomes finite when the time in which the perturbing passes close enough to the target ion to cause a transition approaches the inverse frequency of the transition involved. As a practical matter, the ion temperature must be above 50 eV to allow protons to help drive the $1s3p\ ^1P-1s3d\ ^1D$ transition in helium-like neon. The situation is better for the $3p_{3/2}-3d_{5/2}$ transition in the hydrogenic ion since the energy mismatch is smaller by a factor of 10. The case for $3p_{3/2}-3d_{3/2}$ transitions would be expected to be even better because of the relatively negligible energy difference involved. In the low temperature design approach discussed later in the chapter,

the existence of the $3p_{3/2}-3d_{3/2}$ degeneracy is extremely important in terms of the inversion kinetics since the ion temperature is so low (in the neighborhood of a few electron volts). The calculation of the $3p_{3/2}-3d_{3/2}$ ion collisional cross section through use of the classical path model yields very odd results, namely, that the probability for a transition is negligibly small for energies below a certain value, and above that value the transition probability becomes enormous and must be limited by the unitarity constraint. The impact parameter integration (following the standard recipe) goes out to the relevant cutoff, which in our case always turned out to be the Debye radius. Typically the cross section would be near zero for low energies and then jump to a large constant value ($\pi/2$ times the square of the Debye radius) for all energies above a threshold energy. The problem is that the classical path model is a binary collision model, and in the case of a cold (near 1 eV) highly stripped neon plasma at several times $10^{18}/\text{cm}^3$, the ion positions are strongly correlated. Neighboring ions occur closer to one another than would be allowed by their kinetic energies if they were isolated. The ions are much like molecules in a fluid, and there are associated fluid phenomena; for example, phonon waves. If the transition energy is sufficiently high that only ions coming much closer than a lattice spacing causes a collision, then the binary model may give reasonable results. Such is not the case for the $3p_{3/2}-3d_{3/2}$ transition.

In the situation described above, the $3p_{3/2}$ - $3d_{3/2}$ states are strongly mixed by local quasistatic ion microfields. Dynamic effects also appear to be important; namely, phonon collisions. For example, if a hydrogenic neon ion were pumped to the $3p_{3/2}$ state, some of the time it would make a transition to the $3d_{3/2}$ state as the local lattice might be undergoing a phase of a vibration wherein a neighboring ion is very close and moving. Such vibrations occur on a picosecond timescale. Stark splitting of the m -sublevels occurs; however, the splitting involved is within the linewidths of the 2-3 lines. The linewidths are sufficiently small such that the different 2-3 lines (j - j') are distinct. The reason for discussing the point is that the inversion kinetics depends on the rate coefficient between the $3p_{3/2}$ and $3d_{3/2}$ states, and there is some question as to the proper way to model the ion collisions. Dealing with the individual Stark sublevels is one approach; however, we have kept the nlj description and included the effect of ion collisions (or Stark mixing) as a rate. From a time-dependent analysis of a single ion, the rate is fast if a neighboring ion is near, and slow if the perturbing microfield is weak. On the other hand, unitarity requires that the probability of a transition be less than unity in a given event. We have chosen to use a rate of several times the phonon frequency in the kinetics calculation for the transition. Physically, an ion is unlikely to remain excited for a time longer than the inverse phonon frequency, and the first time a neighboring ion approaches, Stark

mixing will occur. The gain is rather insensitive to the precise value as long as it is faster than the $3d_{3/2}$ decay time (which it is), although uncertainties in the number could affect the gain through the cross section (homogeneous broadening) if it were faster than the $2p_{1/2}$ radiative decay time (0.16 psec).

Electron collisional coupling with the higher Rydberg levels is modeled using the Coulomb-Born results of Perez (1977). These results were calculations for individual angular momentum states that were averaged over all states of a shell (for example, the results give rate coefficients between the $n=2$ and $n=3$ states, rather than the $2p$ and $3p$ states, etc.). The models used for the neon-like through beryllium-like ions are not as detailed and the results of Perez (1977) were used for many of the excitation rates.

Electron collisional ionization is not as important as the excitation processes except in the first few ions (Ne I - Ne IV). Collisional ionization cross sections for the H-like, He-like, and Li-like sublevels were calculated using the hydrogenic results of Golden and Sampson (1980), Moores, Golden and Sampson (1980), Golden, Sampson and Omidvar (1978), and Sampson and Golden (1978). Cross sections in the case of low lying states of the Ne-like through Be-like sequences were taken from results described in Chapter II. For ionization from the excited levels of the Ne-like through Be-like states, some of the cross sections were computed in the classical binary encounter approximation.

Some discussion of the classical model is in order here. Gryzinski (1965) derived expressions for binary collision cross sections in the case of the target electron moving and then averaged the cross sections over velocity distributions of several neutral systems, and obtained agreement with experimental results which appeared impressive. McDowell (1966) and Catlow and McDowell (1967) and others at the time compared the classical results with several neutral systems and found agreement within a factor of 2 in low Z systems. Vriens (1966) obtained an approximate result for exchange. Thomas and Garcia (1969) modified the theory to take into account attraction by the nucleus and found better than factor of 2 agreement with experimental results for several low Z ions. Other authors (see references) have found agreement to within better than factor of 2 with experimental or theoretical results for many systems. In the case of $n \geq 3$ states of the Ne-like through Be-like systems the classical theory is a reasonable standby until quantum theory results become available.

Ion collisional ionization is not important. The references at the end of the chapter include a survey of some of the recent electron collisional processes and results of experiments and calculations.

C. Hydrodynamics and Line Transfer

In the following section we consider target designs which take advantage of hydrodynamic motion (in the form of velocity gradients) to allow greater penetration of resonance pump radiation and to relieve radiation trapping where it would otherwise be damaging. With this type of approach, targets are optimized by obtaining the sharpest gradients possible, and here that is accomplished by using suprathermal electrons to push on the ions through a macroscopic electric field set up when the electrons encounter strong ionization gradients. The production of areas of strong ionization in close proximity to areas of weak ionization initiates the hydrodynamic coupling between the electrons and ions, and is therefore the goal of the target design in what follows. We work in a regime where the degree of ionization of a given region is determined in part by how strong the local radiation field is, and then arrange for brightly lit regions to be close to shadowed regions. The purpose of this section is to discuss briefly some of the issues involved in the geometry of the laser target and to comment on the methods and problems involved in carrying out simulations.

In Figure IV-9 we show the two dimensional mesh used for the symmetric target simulation described in the next section. The mesh covers the interior of the laser medium, extending to the high density flashlamp and filter layers on the left and on the right. Above and below are regions of brass which serve as shields to

Figure IV-9 Two Dimensional Mesh for Hydrodynamics and Line Transfer Calculations

In constructing this mesh we started with 110 points in one quadrant and relaxed the mesh by hand until the zone size and spacing looked reasonable. A program then selected a set of triangles connecting the points which do not cross and which tries to connect nearest neighbors. The resulting mesh is symmetrized to fill four quadrants (although the integrations are only carried out in one) and the result is shown on the left. On the right is shown the zoning resulting from dividing each triangle into three parts and assigning each part to a zone centered at a vertex.

block X-rays and cast shadows. In the simulation we assume that the flashlamps on both sides of the target are irradiated uniformly and coincidentally (no pulse time delay between the two sides), and the result is that the target center sees radiation coming from 80% of the solid angle, whereas the regions along side the brass wedges see only one flashlamp (being in the shadow of the other) and hence radiation from only 40% of the solid angle. The target center is ionized more than the regions above and below, and through the coupling of the electrons (which are assumed to be suprathermals) to the ions via a large electrostatic field produced because of the ionization gradient, the ions are accelerated generally away from the target center. Implicit in this discussion is the fact that the laser medium is optically thin to the ionizing radiation produced by the flashlamp (except at the frequencies of the strong resonance lines).

Given the above rationale for the target geometry we now turn to a brief discussion of the mesh used. This type of mesh is similar to what is described in Fritts and Boris (1979), except that in this case the hydrodynamics is Eulerian rather than Lagrangian. In the present case we selected a set of points which filled the region of the laser target. A mesh program of the author selected triangles which connected neighboring points in such a way as to try to maximize the area to surface ratio of each triangle. The list of points and triangles was reviewed using a program which can display

regions of the mesh, add points, move points and correct poorly zoned triangles, and some mesh relaxation was done. The resulting triangles are shown on the left hand side of Figure IV-9. Irregular polyhedra were then constructed by dividing each triangle up into three sections (selecting the center of each triangle and connecting it to the midpoints of the three sides). Each of the original points selected is within a polyhedra and if the triangles were nearly equilateral, the points will be nearly in the middle of the associate polyhedra. Points on the boundary are associated with polyhedra which would look much like their neighbors if they were reflected at the boundaries. Polyhedra at the vertices are like normal polyhedra which have been sliced proportional to the angle at the vertex, and although the results are reasonable for concave vertices, the resulting zones at convex vertices wrap around the vertex and can give rise to non-intuitive results locally in a hydrodynamics calculation. In the mesh of Figure IV-9, the polyhedra are shown on the right, and one may view that the zoning around the vertices is fine relative to that of the rest of the mesh.

The reason for using zoning of this type (as opposed to the grid type zoning in LASNEX, for example, where an individual zone is quadrilateral) is that since a heavy price is paid for each zone that is included, that it would be desirable to work with as few points as possible and maximize accuracy at each point and through the careful positioning of each point. One example is that the

continuum radiation field at each point on the mesh is computed using fine angular integrations (necessary because of the shadows), resulting in good accuracy at each individual point. The overall philosophy of the routines involved is directed toward this, although the exploration of the different facets possible has only begun.

Sod (1978) has reviewed some of the numerical methods currently in use in hydrodynamics calculations. In the present case we have so far been concerned with isentropic fluid flow within the laser medium, which is to say that the laser medium is shock free until the walls come crashing in (see Wilkins (1980)). The velocities are sufficiently small that the changes in ion density relative to the initial ion density can usually be neglected. In this case only the momentum and energy equations are integrated (explicitly) which saves computer time (the momentum and energy are scalars whereas the density involves on the order of 100 quantities in each zone).

The two dimensional hydrodynamics calculations are intended to be valid only within the tranquil region of the laser medium. Resolving the suprathermal electron dominated shock which is driven inward from the flashlamp and filter regions would take more zones than there is room for on the CDC 7600 using the algorithms currently employed. The shock is likely to be one dimensional, and its dominant effect on the gain calculation is to effectively remove the shocked part of the laser medium from the problem in terms of resonance line radiation transport, as shocking causes strong ion heating and a large velocity shift, which greatly reduces the

absorption cross section and shifts it off resonance for radiation traveling in and out of the laser medium. In the figures shown in the next sections, we simply plot quantities as they would occur in the absence of the shock wave, keeping in mind that the solutions are only valid in a region whose width shrinks as the problem evolves in time. In future calculations a more sophisticated approach may be used; however, for the present purposes we shall be satisfied including the effects of the shock after the tranquil hydrodynamics calculations.

Having discussed briefly some of the hydrodynamics consideration, we turn to the problem of resonance line radiation transfer in the x-ray laser simulations. We differentiate between transport and transfer, the former implying that time retardation of the radiation field is included while the latter implies a steady state computation. In the target designs which have been considered to date the time retardation of the radiation field has been unimportant as light travels 300μ in a picosecond (spatial scales in the designs are 20μ - 100μ while the time scales of interest are dozens of picoseconds), hence the use of a transfer description is quite justified.

The problem of resonance line radiation transfer in the two dimensional geometry under consideration, where time-dependent effects are important in the kinetics as is the case here, and where the bulk velocity shifts are in excess of ten thermal Doppler units, is a very difficult problem. We have not solved the problem in a

manner which allows results for the present type of target to be presented here. We have solved the transfer equations in three one-dimensional geometries (planar, spherical and cylindrical) for time-dependent problems where bulk Doppler shift is important. Earlier versions of the 3-2 laser designs were in cylindrical geometry and were amenable to simulation where the suprathermal driven electron shock wave was included. The algorithms extended to the two dimensional case and designs with rectangular and triangular cross sections were examined, also where Doppler shift was important. In fact it was the results of these simulations that have led us to the "butterfly" type targets currently under investigation. There is not a great difference between the two dimensional results for simple geometries and similar problems using cylindrical targets with similar dimensions. With a triangular target, we had hoped that the fact that the center was closer to the edges would help relieve radiation trapping, especially on the 3-2 lasers where it is very important and ruinous if the optical depth is too great. This led to consideration of more complex geometries; for example, "star" patterns and the "butterfly" target of Figure IV-9.

The use of the more complex geometries suggested the exploitation of shadows to help create ionization gradients and ultimately to develop strong velocity gradients. The advantage of velocity gradients has been well-understood, and the use of suprathermal electrons to drive them has been known at least locally for a period of years.

The possibilities suggested by using shadows to create ionization gradients which can couple the momentum of the hot electrons to the ions of the laser medium has only been appreciated recently, and for this reason there has not yet been time to adapt the algorithms to the new situation. We do not present radiation transfer results in the following two sections, as the development of the requisite software is not complete.

The basic difference between the calculation here and our earlier work (which may be published elsewhere in the future), is that the velocity gradients currently obtained are much larger than what we have previously dealt with. In the next sections, much of the laser medium of interest has been shifted up to ten Doppler units of velocity (see, for example, the asymmetric target results of section D) and small regions in targets not discussed here have achieved numbers higher by three fold over small regions. It is much harder to transfer radiation in the presence of these strong velocity gradients, especially when laser gain is sensitive to radiation trapping, than in the case where only two or three Doppler widths are involved. Part of the problem is involved with supplying more energy bins, but part is also involved in the relative Doppler shift from one spatial zone to another, which can now be a linewidth. Simple interpolation does not do as well in the present case.

Another major problem has to do with the angular integrations necessary to get reasonable results in the presence of sharp absorbing

edges, such as occur in the targets considered in the following sections. In some spatial regions the radiation is very strongly non-isotropic on account of proximity to absorbing regions, and to get smooth results, much care must be taken when using an angle-dependent transfer algorithm. We have not yet developed an algorithm which satisfies the requirements discussed above.

We conclude the section with a brief discussion of some of the relevant literature. The field of radiation transfer is quite old and rich in approaches and techniques (the reviews give excellent accounts of the field; for example, Mihalas (1978), Ivanov (1973), and Hummer and Rybicki (1967)). The transfer equation in the one dimensional cases (planar and cylindrical geometries) were solved by integrating the radiation field along a set of rays whose angles are selected by virtue of their cosines being zeros of Legendre functions of the appropriate order (the S_N methods, of which this is a trivial example, has been standard for decades--see Chandrasekhar (1960)), using an integration of the form

$$I_{n+1} = I_n e^{-\Delta\tau} + S_n (1 - e^{-\Delta\tau}) \quad (\text{IV-C-1})$$

along a ray. Here I_n refers to the radiation field at the n th spatial zone, S_n is the source function and $\Delta\tau$ is the optical depth between the n th and $n+1$ st zones. This type of integration is similar to that used by Hearn (1963) and is cruder and less

efficient than the Feautrier methods which are used in state of the art codes (see references). On the other hand, in most of the problems encountered in the present research to date, the bulk of the time involved in the calculations has been in the kinetics, and the evaluation of an occasional exponential function in the transfer algorithm does not make a large difference in the total run time of the code. This would not be the case in a two dimensional calculation wherein many angles were required (on account of shadows) as well as many energy bins (on account of large bulk Doppler shifts encountered).

In cases of velocity fields where the bulk Doppler shift of the medium is important, the absorption and emission opacities are computed through linear interpolation from the values stored at the energy group boundaries. State of the art algorithms in astrophysics use difference equations in photon energy when velocity fields are important (see the series by the JILA group in the references); however, under shock conditions and with the limited number of zones used this approach would be inappropriate. In non-LTE problems where radiation trapping on optically thick lines is important, some care is required to converge on a self-consistent solution. In the present case, reasonable results have been obtained through explicit solution of the transfer equation

at each timestep, which mimics Lambda iterations (Mihalas (1978)). Although Lambda iterations in optically thick cases are known to converge only after large numbers of iterations, in the case of lines of small optical depth (of order unity to 5) where the albedo is not close to unity (i.e., where scattering is not important) reasonable results can be obtained through use of Lambda iterations. When needed, an implicit correction was added to the source function using a variant of the complete linearization method (Auer and Heasley (1976)). The approximation is termed CRD for complete redistribution (as opposed to the more accurate PRD, or partial redistribution) which is reasonable in cases where scattering is not important.

In the case of two-dimensional problems, an x-y grid is assumed on which a cross section of the laser is zoned. Lasing would occur along the z-axis. A set of usually equally spaced angles are chosen to span the different directions in the x-y plane (ϕ in polar coordinates) and then the angles whose cosines are zeros of Legendre functions are selected to span the angles in the $-z$ plane (θ in polar coordinates). Given the choice of angles, the transfer equation is integrated across the mesh starting with the boundary condition on the "sunny" side of the mesh and marching across in the ray direction toward the "shadow" side of the mesh.

There is not a great deal of literature on two-dimensional line transport of this S_N variant variety with which to compare (see the references).

The radiation transfer does not at present make use of information other than that stored at the points, and one weakness is that it is not conservative. The algorithm is useful for providing at least approximate estimates of the radiation field at the points where the populations are defined and is simple to program. As has been discussed above we have yet to obtain satisfactory results on the target designs of current interest.

D. Laser Design in the Case of the Ne IX 4-3 Laser

In this section we shall discuss some of the issues involved in the design of the helium-like Ne IX 4-3 (55 eV) laser. In the next section we shall examine the situation in the case of 3-2 lasers in helium-like Ne IX and hydrogenic Ne X. The laser design issues are quite similar between the different line pumping schemes and many of the parameters are nearly equal; for example, neon density, laser pump intensity and modal photon density of the resonant pump radiation. Our approach will be to examine most of the issues in the case of the 55 eV laser, which in many ways is a simpler laser having less tight constraints than the 3-2 lasers, and then in the next section review some of the issues which make the design more challenging.

In section A we considered briefly the issue of finding resonances suitable for pumping helium-like $1s^2-1s4p\ ^1P$ transitions in low-Z ions. In the case of helium-like neon, there is the well-known resonance with the Ne X $1s^2-1s2p\ ^1P$ transition. We have examined some of the issues for this pair and will report calculations elsewhere. For present purposes we shall work with an L-M line, two examples of which are listed in Table IV-1c. Some of the issues involved in estimating the pump output were discussed in section A, and it was noted that much uncertainty is associated with the theoretical determination of the pump strength. Over the course of future work in this area more precise values for the modal photon densities at the resonant frequencies will become available; however,

since such estimates are not available now, we shall examine the problem from the viewpoint of finding a constraint on the pump intensity required to drive a short wavelength laser. Although we do not know precisely how much pump strength is available, the analysis which follows will provide an estimate of how much is required. In the case of a laser design for SHIVA where the laser length is limited to 4 mm, one finds that about 0.015 photons/mode of resonant pump intensity is required to achieve 20 gain lengths. Whether 0.015 photons/mode is available at the intensity selected ($6 \times 10^{14} \text{ W/cm}^2$) is not known; however, it seems reasonable to expect that at some laser intensity or under some conditions such pump strength can be achieved. Our purpose in what follows is to examine some of the issues involved in the determination of pump intensity requirements.

Let us review for a moment the main issues involved. The neon must be stripped to the helium-like stage in a timely fashion. This does not appear to be difficult with $6 \times 10^{14} \text{ W/cm}^2$ incident on the flashlamp as will be discussed. Given helium-like ions, one must get the pumping radiation into the laser medium and this is a problem since the absorption length is so small. For example, in the case of $1s^2-1s4p \ ^1P$ pumping the absorption length can be 25μ , and the albedo is not close to unity. The target dimensions are roughly $4\text{mm} \times 400\mu$ along two axes, and in principle a target could be constructed that is of the order of 25μ along the third axis,

although there are target fabrication people who may cringe at such tight specifications. In any case the expansion of the walls into the laser region is likely to destroy the laser medium before it is ready to lase. We have chosen an alternate solution here, that is, the use of a graded velocity flow to relieve the absorption problem.

Under the conditions of interest the ion temperature is quite low (of the order of one eV) and the $1s^2-1s4p^1P$ transition is mostly Doppler broadened with a linewidth near 0.01 eV. Velocity flows with gradients of several times 10^6 cm/sec can be established due to the large suprathermal electron induced electrostatic fields which can be arranged. Such large velocity gradients can give rise to bulk Doppler shifts of the order of 0.1 eV. The width of the pumping line is a few tenths of an eV. The result is that the absorption length of the $1s^2-1s4p^1P$ transition can be greatly increased, maybe by as much as ten-fold. The target under consideration becomes one which is less taxing on the target fabrication resources available.

Due to the large velocity gradients one might naively assume that the neon density could be raised resulting in increased gain. Unfortunately the scaling does not behave in such a fashion, rather the gain decreases with increasing neon density due to increased ion heating and collisional destruction of the upper laser state. The peak occurs near 3×10^{18} ions/cm³ and roles over slowly in either

direction. We have selected an example to present here which seems to be reasonable given that large bulk Doppler shifts relative to the ion thermal motion is required, that is 2.0×10^{18} ions/cm³.

As the flashlamp and filter are heated, the inner edge of the filter expands into the laser medium. The problem is complicated by the presence of suprathermal electrons which can cause a low density fluid to be accelerated to quite large velocities as is shown in Figure IV-10a. Suprathermal electron driven plasma expansion into vacuum can yield ion velocities well in excess of 10^8 cm/sec in the present case where the peak pump intensity is 6×10^{14} W/cm². The hot electron temperature peaks near 11 keV and the primary uncertainty at the moment has to do with their density. If one assumes that all absorbed laser light is carried off by hot electrons at thermal velocities, ignoring the possibility that they might return, one arrives at a hot electron density of 1.5×10^{19} /cm³. Ferusal of the results of the LASNEX calculation yields an ambient hot electron density near 2.1×10^{19} /cm³ on the back side of the target and several times that value on the front side near the critical surface. Hot electrons can inelastically scatter from the dense plasma of the flashlamp and filter system or lose energy on reflection from the receding corona-vacuum boundary. If there were laser medium rather than vacuum against the back side of the target the results from a LASNEX calculation would probably yield a higher value for the hot electron density as the energy loss to a receding vacuum boundary would be replaced by a source of hot electrons generated on the other side of the target. This may result

Figure IV-10a. Density profiles of the inner wall of the flashlamp-filter system expanding into vacuum.

Here are shown snapshots of the parylene region expanding into vacuum. The target used is the Zn/CH target of Figure IV-5 and the laser pulse incident on the zinc is a 90 psec FWHM Gaussian pulse of peak intensity 6×10^{14} W/cm² centered at 135 psec. The edge of the target started at the left hand edge ($Z=0$) at the beginning of the problem ($t=0$). The expansion is driven by suprathermal electrons at low density and one observes matter velocities well in excess of 10^8 cm/sec. One observes that the kinetic energy involved per ion if converted into thermal energy would result in ion temperatures of several hundred thousand electron volts.

in a factor of two increase in the hot electron density relative to the values quoted above. Attempts to calculate this problem with LASNEX have so far proven unsuccessful.

Given the fact that suprathermal electrons are present, we shall briefly examine the consequences thereof. The electron density in the laser medium is about 1.6×10^{19} near the peak of the pumping pulse, which is very likely to be less than the number of hot electrons at the filter-laser medium boundary. As many hot electrons as are allowed by charge neutrality will venture forward into the laser medium, displacing all thermal electrons in their wake via large electrostatic fields. The hot electrons left over, that is, hot electrons which can not be neutralized by the ions of the laser medium, will be turned away at the boundary, depositing their momentum into a strong ion shock as they turn. The resulting density profile as a function of space will develop an odd looking low density shock whose density will be determined approximately by the number of ions required to neutralize excess suprathermal electrons, and whose velocity can be estimated from mass and momentum conservation. Figure IV-10b-d illustrates the situation. Although precise values are not yet known, for a hot electron density of twice the laser medium density, the shock velocity is near 7×10^7 cm/sec. As long as the hot electron density exceeds the electron density of the laser medium the shock velocity scales as the square root of the ratio of the electron densities on either side of the shock discontinuity (which is to say that in this region the shock velocity is a reasonably slow function thereof and hence not too sensitive to

Figure IV-10b.

Vertical axis:

Electron density ($1/\text{cm}^3$)

Horizontal axis:

Distance (cm)

The horizontal line which crosses each of the plots denotes the value of the suprathermal electron density present at the filter-laser medium boundary.

Figure IV-10b. Expansion of the inner filter wall into the laser medium.

Here are shown snapshots of the electron density as a function of space from a calculation involving a very simplified model of the expansion of the filter wall into the laser medium. In this model one dimensional mass and momentum equations (of the shock hydrodynamics equations) are solved assuming that the pressure is due exclusively to the electrons. The electron distribution includes a low temperature thermal contribution whose values were taken from the LASNEX calculation discussed in Section A, and a suprathermal electron contribution whose temperature is given by

$$T_H = 5000 (I/10^{14})^{0.43} \text{ eV}$$

and whose density was given by

$$N_H = 0.3 f I e^{-\Delta x / \lambda} / v_{th} k T_H$$

where I is the pump laser intensity in W/cm^2 , T_H is the electron temperature describing the suprathermal electrons, N_H is the suprathermal electron density, v_{th} is the mean thermal electron velocity of the hot electrons, k is the Boltzmann constant, Δx is the target thickness (initially 0.5μ of parylene), λ is the electron deflection length, the 0.3 factor is due to the assumption of 30% absorption and f is a multiplier (2.5 here) which accounts for the finite lifetime of the suprathermal electrons. The plasma was assumed to have Z and $\bar{A}$ of 8 and 20.17 respectively where Z is the average charge of the plasma and $\bar{A}$ is the average mass per ion in atomic units. These latter quantities enter the problem mostly through their ratio, and the ratio for the highly stripped laser medium is close to that of highly stripped parylene.

Figure IV-10c. Shock position and velocity as a function of time.

For the calculation described in the previous figure we show here the position (as defined by $|u| > 5 \times 10^5$ cm/sec where $|u|$ is the magnitude of the fluid velocity or else by $\rho > 1.1 \times \rho_L$, where ρ is the matter density and ρ_L is the initial laser medium density) and velocity of the shock front as a function of time. At early times one observes the slow thermal expansion of the filter wall, and when the suprathermal electron density exceeds the electron density of the laser medium (which occurs here near 80 psec) a strong shock forms. The velocity of the front exceeds 7×10^7 cm/sec. The effects of the transient ionization of the laser medium on the shock characteristics remains to be examined.

Figure IV-10d. Width of the unshocked laser medium as a function of time.

From the calculation on the previous two figures we have estimated the width of the tranquil region assuming an initial width of $120\ \mu$ and symmetric irradiation of both sides (see Figure IV-9, and assume dimensions of $120\ \mu \times 420\ \mu$ for the width and height). Lasing may occur near 125 psec (the pumping pulse is 90 psec FWHM centered at 135 psec) at which time the unshocked region is about $50\ \mu$ wide. During the remaining picoseconds the width of the tranquil laser medium decreases by about $1.4\ \mu$ per psec. This calculation is dependent on the choice of peak suprathreshold electron density which at present is not known precisely. As the problem becomes better understood the laser target width can be further optimized. In an experiment some portions of the pumping beam will be brighter than others leading to a nonuniform shock front. The ramifications of this remain to be explored.

uncertainties and variations in the hot electron density). In the case of the 3-2 lasers the ion density may be higher, in which case it is conceivable that the electron density in the laser medium exceeds the suprathermal electron density. If this were the case, the shock front would be very much altered from the profiles shown in Figure IV-10b, and the shock velocity would be considerably slower. This possibility is one which we hope to avoid, as the conditions under which this would be true are not conditions under which interestingly large velocity gradients would be obtained. The pressure within the laser medium is determined primarily from the hot electron contribution, and if the hot electrons were free to wander throughout the region uniformly no large pressure gradients would result.

The shock front discussed above is not necessarily a bad effect. On the one hand the laser medium, which is shocked to high density and velocity is heated as well. The basic reason for pursuing conditions under which the suprathermal electron density exceeds the laser medium density is to set up velocity gradients to increase the effective absorption length of the $1s^2-1s4p\ ^1P$ line of Ne IX. However, whatever part of the laser medium is shocked is effectively removed from the picture in terms of absorption on two counts, namely that the shock heating greatly reduces the absorption cross section (Doppler broadening is dominant) and the associated bulk velocity shift moves the line off-resonance for pump radiation

directed inward. As the shock moves inward, more and more of the laser medium is removed from the problem, and the remaining tranquil region is reduced in width. Penetration of the pump radiation becomes less and less of a problem as the shock pushes into the laser medium.

If it were of interest to slow down the incoming suprathermal electron driven shock wave, one approach which has been suggested is the use of an inner layer of high Z material. On the one hand, such an arrangement appears to be Rayleigh-Taylor unstable (see the references), and on the other hand there is little reason to expect it to work in the first place as most of the shock is composed of laser medium which has been compressed and which is piling up behind the shock front. One major problem involved with the approach discussed in this section is that spatial inhomogeneities in the pump intensity will tend to produce locally hot regions from which faster shocks will be driven. Some parts of the laser medium will be pinched off before others. As to whether such effects are disastrous or merely bad has not yet been investigated. Certainly the gain is removed in regions where the shock has pinched off a region, and the question is whether in their place large regions of loss are produced.

We are now ready to review some calculations of the kinetics and hydrodynamics of a laser design. The flashlamp and filter system of section A was used and time-, energy- and angle-dependent spectra were calculated using data from a LASNEX dump file. Although

the pump candidates for the helium-like Ne IX $1s^2-1s4p\ ^1P$ transition include nickel and copper, the spectra calculated by LASNEX for zinc is probably sufficiently close for present purposes. The absolute accuracy of the LASNEX calculation is not known for this problem and the present author would not be surprised at factors of two or more differences between the calculated spectra in the spectral region near the L-M transitions and the "right" answer if it could be determined. The target whose mesh is found as an illustration in section C (Figure IV-9) was used as a design. The target extends 420μ from bottom to top and 120μ from left to right. The central restriction is 84μ from vertex to vertex. The flashlamp and filter were placed a few microns to the right and left of the target, and spectra for each point in the laser medium were calculated by integrating contributions from the faces of both source slabs.

We have calculated the neon kinetics at the center of the laser target with the results shown in Figures IV-11-15. The incident laser pulse is 90 psec FWHM centered at 135 psec with peak intensity 6×10^{14} W/cm². The initial neon density is 2×10^{18} /cm³. In the calculations in Chapter III we kept the electron temperature pinned at a value close to the peak hot electron temperature. Here we have let the ambient electron temperature follow the hot electron temperature in time. The time dependence of the electron density is shown in Figure IV-11a. The stripping through the different isoelectronic sequences is shown in Figure IV-11b and Figure IV-12 depicts the

Figure IV-11a Electron Density Versus Time at the Target Center

The peak of the pumping pulse occurs at 135 psec in this problem and one observes that the stripping of the neon ions takes place mostly before the pulse reaches its peak intensity. The initial conditions for the laser medium were 2×10^{18} atoms/cm³ of atomic neon at room temperature. At higher initial neon densities (for example, near 10^{19} /cm³) the plasma strips down much more rapidly due to a supra-thermal electron driven collisional avalanche. Under the conditions shown here electron collisional stripping is a significant factor although not dominant in the transient ionization process.

Figure IV-11b Fractional Populations of the Neon-like Through Helium-like Sequences (A: Ne-like, B: F-like, C: O-like, D: N-like, E: C-like, F: B-like, G: Be-like, H: Li-like, I: He-like)

The production of the helium-like sequence is observed to be no problem when the flashlamp and filter system is irradiated with $6 \times 10^{14} \text{ W/cm}^2$. Below $3 \times 10^{14} \text{ W/cm}^2$ it is somewhat of a problem for neon, however, for lower Z ions the plasma can be brought to the helium-like sequence at lower intensities. The maximum fractional of the helium-like sequence is near 50%, and without the strong photoexcitation of the $1s^2-1s4p \text{ } ^1P$ transition the peak value would be closer to 70%. The helium sequence is seen to occur slightly ahead of the pumping pulse maxima.

Figure IV-12 Fractional Ion Populations of the Final Sequences

The stripping of the helium-like sequence leads to a 70% value for the hydrogenic sequence. Later when we examine inversions in the hydrogenic sequence, we shall find the situation reversed, that the helium-like stage peaks near 70% and the hydrogenic maximum is less than 50%. One observes that the sequences appear only transiently and that the pumping strongly affects the duration of the pumped sequence. This effect is due in the present case to the low density and hence small recombination rates. Furthermore, there is a tradeoff between the duration of gain and the maximum gain attainable given that one could pump with a source of arbitrary brightness.

Figure IV-13a. Ion Temperature as a Function of Time

Presented here is the ion temperature during the course of the problem where only electron-ion coupling has been included. The ions are strongly coupled (collisionally) among themselves and we therefore assign the sequences a single ion temperature. The hydrodynamic work on the ions is quite small. For example, isentropic heating or cooling depends on the fractional decrease or increase of a Lagrangian volume element, and in the present case where the velocity gradients are of the order of 10^6 cm/sec the relative expansion or compression of the fluid is only a few percent. The shock heating of the ions is also minimal, except when the inner walls of the filter region compresses the plasma. Neglected is electron-ion current driven and two stream instability heating. It is hard to imagine how instabilities could grow when ion-ion collisional damping is as strong as it is in the present case.

Figure IV-13b. Ion Temperature Due to Proximity Heating

Here is shown the result of a Monte Carlo calculation of the ion temperature using ionization rates from the XRASER calculation under consideration. The initial ion neon density was $2 \times 10^{18}/\text{cm}^3$ and the ions started out at room temperature. One observes the contribution due to proximity heating to be slightly in excess of one electron volt at the time of lasing. Electron-ion collisional coupling was not included in this calculation. The ion temperature when both effects are included should not be far from the sum of this figure and the previous figure. Electron Debye screening was not included however due to the large electron Debye length this approximation should be excellent.

Figure IV-14a. Fractional Level Populations of the Helium-like Ne IX M and N singlet levels.

We continue in this figure the analysis of the kinetics of the problem under consideration ($2 \times 10^{18}/\text{cm}^3$ initial neon density, $6 \times 10^{14} \text{ W/cm}^2$ on a Zn/CH flashlamp and filter system, and 2-D target of Figure IV-9). The fractional populations are given as compared to the total neon density, for example, a fractional value of 0.008 for the $1s4p \ ^1P$ level equates to $1.6 \times 10^{16}/\text{cm}^3$ for a level density. The population pumped to the $1s4p \ ^1P$ state is collisionally distributed (electron collisions) to the other singlet N-shell states, and on a relative scale the coupling is weak. The separation of the 4^1D and 4^1F levels is not known precisely, however, the energy difference may be less than 0.01 eV in which case ion collisions could play a role in the kinetics (ion collisional coupling of the 4^1D and 4^1F was not included in the present calculation). Either neon ion collisions, or else the addition of some protons may increase the 4^1D - 4^1F coupling and could result in a factor of two increase in gain if the levels are sufficiently close.

Figure IV-14b. Fractional Inversion Densities on the Strong 3-4 Lines

In this figure we have plotted the quantity

$$n_{\text{frac}} = (N_u - \frac{g_u}{g_L} N_L) / N_0$$

where N_u and N_L are the upper and lower state populations and N_0 is the total neon population. g_u and g_L are the upper and lower state statistical weights. The strongest inversion occurs on the 4^1P-3^1S line which has the lowest stimulated emission cross section of the lines considered. If the ion temperature were several electron volts higher and protons were introduced into the plasma, ion collisional coupling among the N-shell states could result in fractional inversion densities a factor of two higher on the 4^1D-3^1P line, and several-fold higher on the 4^1F-3^1D line.

Figure IV-14c. Gain on the Strong M-N Lines as a Function of Time

The small signal gains were computed assuming 0.015 photons/mode on the $1s^2-1s4p$ $1P$ line and 10^{-5} photons/mode on the $1s^2-1s2p$ $1P$ and $1s3p$ $1P$ transitions. The ion temperatures used in the calculation were those of Figure IV-13a, which neglected proximity heating. The inclusion of the effect would be most telling on the 4^1F-3^1D transition where a 20% decrease in gain may occur. Ion collisional coupling of the 4^1D-4^1F could result in a factor of two increase in gain if the energy defect is sufficiently small. The gain was computed using Voigt profiles for line shapes and taking the homogeneous broadening from the inverse of the state lifetime, and the values shown above are for the line center values. Static ion Stark broadening is a small contribution under present conditions and has been neglected in this calculation.

Figure IV-14d. Stimulated Emission Cross Sections on the Strong 4-3 Transitions

The stimulated emission cross sections were computed at line center, including homogeneous and inhomogeneous broadening. The gain is proportional to $g_l f / g_u$ where g_l and g_u are the statistical weights of the lower and upper states and f is the absorption oscillator strength. This product is equal to 0.14, 0.37 and 0.73 for the 4^1P-3^1S , 4^1D-3^1P and 4^1F-3^1D transitions respectively. The gain is also decreased by the total decay rates of the state involved. The $1s3p \rightarrow 1P-1s2 \rightarrow 1S$ decay provides the major component to the homogeneous broadening of the 4^1D-3^1P line and the $1s4p \rightarrow 1P-1s2 \rightarrow 1S$ provides most of the same for the 4^1P-3^1S line. The 4^1F-3^1D line has much less homogeneous broadening.

Figure IV-14e. Voigt parameters for the 3-4 transitions.

The Voigt parameter is defined (see Hummer (1965)) as

$$a = \frac{\gamma}{\frac{v_0}{c} \sqrt{\frac{2kT_i}{M_i}}}$$

where a is the Voigt parameter, γ is the transition damping rate, v_0 is the transition frequency, c is the speed of light, M_i and T_i are the ion mass and temperature and k is the Boltzmann constant. In the above figure one observes that the 4^1F-3^1D transition has a Voigt parameter much less than the other lines, indicating the reduced homogeneous broadening of the line. The 4^1D-3^1P transition has the largest Voigt parameter by virtue of the $1s3p \rightarrow 1p-1s^2 \rightarrow 1S$ radiative decay.

Figure IV-15. Peak Absorption Cross Sections on the Resonance Lines

The peak values reached correspond to absorption lengths of 2.5μ , 10μ and 25μ on the $1s^2-1s2p \ ^1p$, $1s^2-1s3p \ ^1p$ and $1s^2-1s4p \ ^1p$ transitions. Under the conditions described here these lines are primarily Doppler broadened. Although the gain does not increase monotonically with increasing ion density, the absorption coefficients of the resonance lines do, hence in the interests of designing a target which is larger than a few microns on a side it behooves one to consider low ion density targets. This calculation does not include proximity heating, the inclusion of which reduces the peak values some ($\approx 30\%$); however, the basic points regarding trapping in this chapter remain intact.

final three sequences. Photoexcitation of the $1s^2-1s4p\ ^1P$ transition was taken to be constant in time at 0.015 photons/mode, which causes an increased ionization rate of the helium-like sequence on account of the additional N-shell population present. The value of 0.015 photons/mode was determined by trying several different pump intensities, and selecting one which would result in gain near 50 cm^{-1} . It is from this exercise that the constraint discussed early in the section has been determined, and in the remaining part of the section we shall explore details relating to the operating point.

The ion temperature resulting from electron-ion collisional coupling is shown in Figure IV-13a (the ions are strongly coupled between themselves and only weakly coupled to the electrons, hence one ion temperature was used to describe the temperatures of all the sequences). The ion temperature resulting from proximity heating is shown in Figure IV-13b. Proximity heating is the heating which occurs when nearby ions become ionized and repel one another at nonthermal velocities. The ion temperature resulting when both effects are included should be nearly the sum of the temperatures shown in the two figures. One observes that proximity heating dominates at early times when the ion positions and velocities are least correlated. After about 150 psec electron collisional heating dominates. The low ion temperature approach to laser design which is discussed in this chapter should be reasonable even if the ions

are heated to several electron volts (which would be the case if some weak anomalous heating should occur), although the gain would be less due to the additional contribution of the increased ion thermal Doppler broadening.

In the preceding chapter we were concerned briefly with the issue of electron-ion instabilities on account of the increase in ion temperature such instabilities would produce were they to occur. Since the high gains achieved were in part due to the low ion temperatures which occurred, it was of interest to check to see that anomalous heating didn't spoil the gain. In the present case we again have a low ion temperature, which by itself is not so crucial as it is in the case of the Ne II laser. The line-pumping approach can be implemented in the case of a low ion temperature as well as in the case of a high one, and there are features to recommend both. In the low temperature case one can take advantage of large velocity gradients to aid in the penetration of resonance pump radiation, as is being examined in this section. The low temperature approach has the unfortunate property of leading to weak coupling of nearly degenerate states between which dipole-allowed collisional transitions may occur. For example, in the case of the 55 eV laser scheme, the coupling between the different levels of the N-shell is quite weak, and the gain on the $4^1F - 3^1D$ transition is severalfold lower than it would otherwise be if ion collisional processes were to play a role in the kinetics. A similar situation exists in the case of the

helium-like Ne IX 3-2 laser scheme. If the ion temperature were high (several tens of electron volts), then the design procedure would include the optimization of hydrogen concentration within the laser medium to provide protons for better ion collisional coupling of the $1s3p\ ^1P$ and $1s3d\ ^1D$ levels. We have performed calculations of this type, and from the results obtained, it is not at all obvious that the low temperature approach is the best. A high ion temperature could be produced in the laser medium simply by allowing the suprathermal electron driven shock wave across the entire laser medium before pumping an inversion. The high ion temperature approach to short wavelength laser design is a competing design philosophy with the one discussed here, and in the future will probably constitute a major area in the field of short wavelength laser design.

The reason for bringing up the point here has to do with our discussion of the status of electron-ion instabilities in the present design. In the present design, both electron and ion densities are very much higher than in the case of the Ne II laser discussed in the previous chapter, and for this reason the wave frequencies and potential instability growth rates are considerably higher. We have calculated growth rates in the collisionless limit with the result that growth rates on the order of 10^{12} /sec may be possible. The waves appear to be strongly damped by ion-ion collision, so much so that one hesitates to call them waves at all as the ion collisional damping rate exceeds the ion acoustic wave frequency. This result

is due to the moderately large Z involved (the ion-ion collisional coupling rate is proportional to Z^4) and the low ion temperature involved (the rate scales as $T^{-3/2}$, where T is the ion temperature). If the drift velocities are sufficiently large, it is possible to drive the plasma unstable, even with the heavy ion collisional damping which is present. The question of how large is sufficiently large requires some calculation, and our estimates suggest that the requisite drift velocities correspond to a situation in which the replacement of all thermal electrons within the laser medium within a very short time (less than 10 psec). In fact, the replacement of thermal electron occurs only as new electrons are produced, and the number produced in 10 psec is small compared to the total number of electrons present after the first couple of electrons have been removed on the average from the neon ions. Attaining the large drift velocities required to drive electron-ion instabilities appears difficult, and at present the laser plasma seems to be quite stable by virtue of collisional damping.

In Figures IV-14a-e are shown results pertaining to the 4-3 small signal gain. The pumping field on the $1s^2-1s4p\ ^1P$ transition is assumed to be constant at 0.015 photons/mode for this set of calculations. If there were no collisional or radiative losses from the $1s4p\ ^1P$ state other than radiative decay to the ground state, the upper state population would be very nearly 0.045 of the ground state population (the $1s4p\ ^1P$ state has threefold degeneracy). The ground state population peaks near 0.5 of the total population, and hence the $1s4p\ ^1P$ state could have no more than about 0.0225 of the total population (and would have very nearly that figure at low electron densities). In Figure IV-14a one observes a peak population of about 0.008 for the $1s4p\ ^1P$ level, down a factor of almost three from its loss-free value. The other N-shell states have lower fractional populations, and the populations of the N-shell states well exceed that of the M-shell states. The fractional inversion densities are shown in Figure IV-14b from which one observes that the lower state populations reduce the inversions by about 25% from the upper state populations. Figure IV-14c shows the small signal gain on some of the M-N transitions computed using the ion temperatures of Figure IV-13a. Although the inversion densities are strongest for the 4^1P-3^1S , 4^1D-3^1P and 4^1F-3^1D transitions respectively, the magnitude of the transition gains are in reversed order. The reason for this can be seen from the following plots--in Figure IV-14d are illustrated the stimulated emission cross sections at line center for the three lines, from which it is seen that the 4^1F-3^1D cross section is by far the

largest (due to the larger gf product of the transition and to the reduced homogeneous broadening as there are no dipole-allowed decays to the ground state). One observes the marked time dependence of the cross sections which is a measure of the importance of thermal Doppler broadening in the gain calculation. As the effects of proximity heating were not taken into account in these calculations, the corrected values will be lower by varying amounts depending on the line and time--the 4^1F-3^1D transition is likely to lose 40% of its peak gain to Doppler broadening induced by proximity heating. The Voigt parameters involved are shown in Figure IV-14e.

In Figure IV-15 are shown the absorption coefficients of the helium-like resonance lines which occur in the problem under consideration. As advertised earlier, the absorption coefficients are quite large. The absorption length on the $1s^2-1s4p\ ^1P$ line is near 25 μ .

So far in our discussion of laser kinetics we have been concerned with one operating point, among other factors we have chosen the initial gas density to be $2 \times 10^{18}/\text{cm}^3$. In Figure IV-16 we illustrate the results of a parameter study where the initial neon density is varied. The calculation is similar to the one described in the previous pages except that it includes twenty zones, each at a different density, and a coarser timestep. The small signal gains peak in the vicinity of $2-4 \times 10^{18}/\text{cm}^3$ and roll over slowly in density. This calculation does not include the effects of proximity heating; however, only the 4^1D-4^1F will be affected significantly by the increased ion temperatures. Figure IV-17 shows the time histories of the gains from which one observes the sensitivity of the inversion timing to initial neon density--at high densities (above $5 \times 10^{18}/\text{cm}^3$) the plasma passes through the helium-like sequence prematurely (the pumping pulse peaks at 135 ps and is at half maximum near 90 psec). Under the conditions of interest the highest gain occurs somewhat before the peak of the pumping pulse, hence, for this target and intensity one might prefer a thinner flashlamp (1000 $\text{\AA}$ of Zn was used here) so that the peak line emission would be better matched to the occurrence of helium-like ions (in the kinetics we have assumed line pumping which is constant in time whereas in reality it may be strongly peaked in time).

Figure IV-16. Gain on the Strong 3-4 Lines of Ne IX as a Function of Neon Density

In this calculation the flashlamp spectra as would be observed at the center of the mesh of Figure IV-9 were used to drive the ionization of neon at twenty different initial densities, from 10^{17} /cm³ to 10^{19} /cm³. A pump intensity of 0.015 photons/mode was used for the $1s^2-1s4p\ ^1P$ transition in all zones. At low densities one finds the 4^1P-3^1S gain to be highest as electron collisional transfer to the 4^1D and 4^1F is weak. As the density increases, the coupling does as well and the strongest gain occurs on the 4^1F-3^1D lines by virtue of its larger stimulated emission cross section. This curve was generated neglecting proximity heating which would affect the 4^1F-3^1D line most strongly. Qualitatively the curve would retain its overall shape but the magnitude would be lower with the inclusion of proximity heating. Countering this, the ion collisional coupling of the 4^1D-4^1F transition raises the gain of the 4^1F-3^1D line but preserves its general shape.

Figure IV-17. Small Signal Gain on the 4^1P-3^1S Transition as a Function of Time at Different Ion Densities

The peak gain occurs near $3 \times 10^{18} / \text{cm}^3$ on this line. The sensitivity of the timing to ion density is shown here, in particular the peak gain slightly precedes the maximum of the laser pump pulse. For initial neon densities between $10^{18} / \text{cm}^3$ and $4 \times 10^{18} / \text{cm}^3$ the time dependence and magnitude of the gain appear quite reasonable. At higher densities the electron collisional processes help bring the neon to the helium-like sequence prematurely and cause reduced upper laser state population (primarily due to N-O collisional transitions). At low initial neon densities, the gain peaks later in time and has lower values, simply because of a numerical decrease in the upper laser state population.

Ion temperatures for the parameter study are shown in Figure IV-18a and electron densities appear in Figure IV-18b. With regard to the former, one observes that electron-ion collisional coupling cannot bring the ion temperature to more than a few electron volts under condition of interest here. The situation including proximity heating is less clear at the moment; qualitatively the curves should rise faster early on before the ions have had a chance to arrange themselves in space (to move apart), but the end result is probably the same. Perusal of the electron densities reveals the relative importance of electron collisional processes under the different conditions of the parameter study.

In Figure IV-19 we show the results of another parameter study, namely the variation of the small signal gain as a function of $1s^2 - 1s4p \ ^1P$ pump strength. The initial neon density is $2 \times 10^{18} / \text{cm}^3$ as is the case in previous calculations so far reviewed. Proximity heating was not included. The point of the illustration is that there is a departure from linearity at high pump strength, and this departure is due to the fact that increased pumping of the N-shell leads to higher excited state populations, which in turn leads to a greater depletion rate of the helium-like isoelectronic sequence. Near a pump strength of 0.015 photons/mode, the gain is very nearly linear in pump intensity.

Figure IV-18a Ion Temperatures as a Function of Time for Different Initial Neon Densities

Using the same spectral data as in previous figures we have calculated the ion temperatures including electron-ion collisional coupling. The electron-ion coupling is observed to be relatively weak at the low density of the laser medium--at $10^{19}/\text{cm}^3$ initial neon density the ions are only a little above 10 eV by the end of the problem. The use of ion collisions to assist collisional transfer of the $4^1\text{P}-4^1\text{D}$ transition (0.2 eV) or the $3^1\text{P}-3^1\text{D}$ transition in the 151 eV scheme requires collisional coupling alone at these low densities and short pulses. Proximity heating (not included in the above figure) can help some, however, the bottom line is that attainment of ≈ 100 eV ion temperatures (where proton collisions could be of considerable use) does not appear likely. It was for this reason that much of the present research has focused on the low ion temperature approach.

Figure IV-18b Electron Densities as a Function of Time for Different Initial Neon Densities

Continuing the examination of the parameter study, here are presented electron densities from the same calculation as for the previous three figures. One observes collisions in the transient ionization process. For initial neon densities below about $2 \times 10^{18}/\text{cm}^3$, the curves roll over between 80 ps and 90 ps (the pumping pulse is at half maximum at 90 ps). At higher initial neon densities electron collisions play an increasingly more dominant role in the stripping process. The operating point which has been selected (2×10^{18}) is where both collisional and photo- (and Auger) processes play comparable roles in the transient ionization.

Figure IV-19. Variation of the small signal gain on the M-N lines as a function of resonant line pump strength.

In this plot we show the peak gain reached as the $1s^2 - 1s4p \ ^1P$ transition pump radiation field is varied. The X-ray spectra at the center of a symmetrically irradiated target (the same problem as has been considered earlier in this section, see Figure IV-9) was used to drive kinetics in 10 zones, each of which saw a different pump intensity. One observes from the results that the gain is very nearly linear in the pump photon field at low pump intensity and becomes increasingly more nonlinear at high pump intensities due to the additional depletion of the helium-like ions. This figure also indicates the need for a bright pump source as gains below 10 cm^{-1} are unlikely to be easily observable in a SHIVA experiment given the constraint that only 4 mm length can be pumped.

In the following sequence we examine some results from two dimensional studies of the target. The goal of this type of target design is to set up usefully large velocity gradients using suprathermal electrons to push the ions via the large macroscopic electrostatic fields set up as the suprathermal electrons transport throughout the laser medium. We assume that the hot electrons can wander throughout the target uniformly to the extent to which they are allowed by charge neutrality. We assume further that locally the distribution is isotropic and can be described by a Maxwellian at the suprathermal electron temperature. Given these assumptions one notes that pressure gradients can be set up if one region of plasma is more highly ionized than a neighboring region of plasma, as the pressure is dominated by the suprathermal electron contribution which in turn is proportional to the local electron density (and hence the degree of ionization). One way to produce ionization gradients (and therefore pressure gradients) is to shield parts of the laser medium from the flashlamp radiation, and insofar as the ionization process is dependent on the magnitude of the pumping flux, the shadowed regions will be ionized more slowly resulting in the desired ionization gradients. At $2 \times 10^{18}/\text{cm}^3$ initial neon density the photonic contribution to the plasma stripping is considerable and the scheme is self consistent.

In Figure IV-20 is illustrated a snap-shot of the electron density shortly before the peak of the pumping pulse for a symmetric two dimensional target. In this design flashlamp and filter systems of the specifications given in Figure IV-5 of Section A are placed a few microns to the left

Electron Density
($1/\text{cm}^3$)

A: $1.44\text{E}+19$

B: $1.46\text{E}+19$

C: $1.48\text{E}+19$

D: $1.49\text{E}+19$

E: $1.51\text{E}+19$

F: $1.53\text{E}+19$

G: $1.54\text{E}+19$

H: $1.56\text{E}+19$

I: $1.58\text{E}+19$

J: $1.59\text{E}+19$

Figure IV-20 Electron Density in Space for a Symmetric Target

In this figure and several following figures we show the results of a two dimensional calculation of the laser medium as it is ionized and as velocity flows are set up. The initial neon density is $2 \times 10^{18}/\text{cm}^3$ and the flashlamp-filter system is the same Zn/CH target as shown in Figure IV-5. The time of the snapshot is 130ps, or 5 psec before the pumping pulse maximum. The region of highest electron density is roughly the center of the target, which is furthest stripped as it sees the full radiation from flashlamps on both sides. Above and below the stripping proceeds more slowly as the plasma in the outer arms lie in the shadow of the central wedges and see only about half the radiation field seen at the target center.

and right of the laser medium and irradiated uniformly and coincidently (that is, the pumping pulses are identical and are incident at the same time on both sides) with 90 psec FWHM Gaussian pulses of peak intensity 6×10^{14} W/cm². The wedge-like regions which protrude into the laser medium are assumed to be infinitely absorbing for photons. The idea is that the wedge-like protrusions will shield the arms from radiation from the flashlamp on the opposite side. The central region will see the full radiation field from both sides and strip more rapidly. Velocity gradients should be strongest for y-directed velocities (up-down in the following figures), and indeed calculations confirm this. The only issues to be settled concern the magnitude of the gradients and whether they are sufficient to allow penetration of the pump radiation into the interior of the target.

Figure IV-20 shows isocontours of the electron density resulting from symmetric irradiation of the laser target. One observes nearly one electron per neon ion difference in average stripping between the central region and the wings, with the sharpest gradients occurring near the vertices of the central restriction. In Figure IV-21 are shown the resulting pressure contours and in the following figures one can view velocity contours. Velocities of several times 10^6 cm/sec are obtained corresponding to several thermal units. In Figures IV-22 and IV-23 we show the x-directed (right-left) and the y-directed components of the velocity vector. In the following figure (Figure IV-24) is shown the y-directed velocity contours in units of the neon ion thermal velocities, from which one observes in excess of ten thermal

Pressure₀
(ergs/cm³)

A: 2.44E+11

B: 2.47E+11

C: 2.50E+11

D: 2.53E+11

E: 2.56E+11

F: 2.59E+11

G: 2.61E+11

H: 2.64E+11

I: 2.67E+11

J: 2.70E+11

Figure IV-21 Pressure in space for a Symmetric Target

The pressure is dominated by the contribution from the suprathermal electrons. In this calculation we have modelled the suprathermal distribution as being isotropic and Maxwellian, and have assumed that the thermal electrons are displaced and heated, such that the electrons are characterized by one density and temperature. This may present some problems in the present target as the distribution peaks in the center of the target--in an electron transport calculation the hot electrons may have trouble getting there. In fact, from the pressure gradients which can be deduced from the above figure the electrons actually push outward from the center of the target. The asymmetric targets to be examined shortly produce spatial distributions which are more intuitive. The pressure in this calculation includes the effect of the macroscopic electrostatic field which is set up as the electrons are constrained by approximate charge neutrality.

Velocities
(cm/sec)

A: -2.00E+06
 B: -1.56E+06
 C: -1.11E+06
 D: -6.67E+05
 E: -2.22E+05
 F: 2.22E+05
 G: 6.67E+05
 H: 1.11E+06
 I: 1.56E+06
 J: 2.00E+06

Figure IV-22 X-directed Velocities for the Symmetric Target

Here are shown isocontours of the x-component (right-left) of the ion fluid velocity at 130 psec in the problem considered in the previous several figures. Although the spatial zoning may be crude the story is pretty clear; namely, that strong velocity gradients may be set up. The magnitude of the gradient should be compared against the thermal ion velocities which are near 3.5×10^5 cm/sec in the present calculation. The velocity gradients are set up in the x-y plane (shown here), while the short wavelength laser pulse would propagate in the z direction (out of the page) and be unaffected by the perpendicular velocity gradients. Inhomogeneities in the laser pumping pulse may set up z-directed velocity flows, and part of the design procedure must be to ensure that the ratio of parallel to perpendicular velocities is small.

Velocities
(cm/sec)

A: -2.00×10^6
 B: -1.56×10^6
 C: -1.11×10^6
 D: -6.67×10^5
 E: -2.22×10^5
 F: 2.22×10^5
 G: 6.67×10^5
 H: 1.11×10^6
 I: 1.56×10^6
 J: 2.00×10^6

Figure IV-23 Y-directed Velocities for the Symmetric Target

The y-directed velocity gradients are observed to be larger than the x-directed gradients of the previous figure. The up-down directed velocity flow is useful for allowing penetration of resonance pump radiation incident from the upper and lower regions of the flashlamp. Some of the jaggedness in the contours comes from the relatively coarse spatial zoning used, and some is due to the limited number of angles used in the integrations involved in determining the x-ray spectrum at each point.

Figure IV-24 Y-directed Velocities Normalized to the Local Ion Thermal Velocity

In this figure we have focused on the region of the plasma closest to the vertex of the upper restriction. The velocities have been normalized to the local ion thermal velocities where the ion temperatures were computed including contributions from electron-ion collisions, proximity heating (estimated using a crude analytical model corrected to be in near agreement with the Monte Carlo results), and from nearly isentropic $p dV$ work. One observes that the gradient is near one thermal velocity unit per 5μ . The comparison of the bulk velocity gradients with the local ion velocities is relevant since Doppler broadening dominates the linewidths of the helium-like resonance lines.

units worth of gradient. The ion temperature resulting from electron-ion collisional coupling alone is shown in Figure IV-25a, and in Figure IV-25b we show a calculation of the ion temperature where proximity heating has been taken into account. A simple analytic formula was used which included only the contribution of neutral atoms which are ionized in proximity to the charged ions. The Monte Carlo results naturally included all combinations; however, it is the ionization of the neutrals which dominates the proximity heating rate. A multiplicative factor near unity was included to bring the analytic result into near agreement with the result shown in Figure IV-13b.

The calculations described above do not include the effects of the suprathermal electron driven shock wave (which compresses the laser medium encountered as it travels inward from the inner walls of the flashlamp and filter system). We do not expect a similar shock to develop at the interface between the brass and laser medium regions, primarily for the reason that this boundary is not located within microns of the critical surface where suprathermal electrons are generated. It is possible that a large suprathermal electron flux will run along the inner edge of the brass next to the neon plasma, in which case a similar shock would develop, and the laser target would have to be redesigned.

The problem of resonance line radiation transfer was discussed in section C, and there it was discussed that there were difficulties

Ion Temperature (eV)

- A: 2.40E-01
- B: 2.51E-01
- C: 2.63E-01
- D: 2.74E-01
- E: 2.86E-01
- F: 2.98E-01
- G: 3.09E-01
- H: 3.21E-01
- I: 3.32E-01
- J: 3.44E-01

Figure IV-25a Ion Thermal Temperature in Space (Neglecting Proximity Heating)

In this figure we give the ion temperatures for the symmetric target at 130 psec. In this case the temperature at the target center is near a third of an electron volt. The ion heating rate due to electron ion collisions is large enough to bring the ion temperature over 1eV within the next 25 psec. The contribution due to proximity heating was not included in the calculation shown here, but is given in the following figure. The gradients illustrated here come about because the ionization of the center part of the target was most advanced, hence the ions were exposed to more electrons over the course of the problem. Of interest to note is that the variation is only 0.1 eV over the mesh, which is large compared to the 0.3 eV mesh average. When proximity heating is included the variation is not much more than 0.1 eV (excluding the noise at the vertices in Figure IV-25b). Proximity heating in the approximation used here is not so sensitive to the X-ray photoionizing flux.

Ion Temperature
(eV)

A: 9.00E-01
B: 9.22E-01
C: 9.44E-01
D: 9.67E-01
E: 9.89E-01
F: 1.01E+00
G: 1.03E+00
H: 1.06E+00
I: 1.08E+00
J: 1.10E+00

Figure IV-25b Ion Thermal Temperature in Space

Inclusion of proximity heating in the calculation has increased the ion temperature by about 0.75 eV from the values shown in the previous figure. This figure is noisier than the previous one due to a different boundary condition imposed. In the previous figure the velocities at the boundaries were set to zero, mostly to reduce the noise observed here. The basic problem has to do with the zoning near the vertex, which is not sufficiently fine to result in an accurate determination of the shock heating contribution to the ion temperature near the boundary next to the vertices. In reality, the brass will blow in some (which is not included here) and the numerical problems discussed above are outside the region of validity of the model.

involved in carrying out such a calculation in the case of the present laser design. Rather than present the results of a transfer calculation, we instead examine a much simpler calculation wherein the pump radiation is assumed to suffer no absorption at all while penetrating the laser medium. Although the velocity gradients produced are substantial, it is unlikely that penetration would be satisfactory were 120 μ of tranquil laser medium present. So far we have not included the effect of the suprathermal electron driven shock which crashes in on the region from both sides (from the data of Figure IV-10d one might estimate the width of the tranquil region to be near 40 μ at 130 psec). The value of the gain when no absorption of the pump radiation occurs models the limit where the shock pinches off the laser medium, and for this reason it is a relevant quantity. In Figure IV-26 we show the gain computed under this approximation. The peak gain calculated occurs at the target center, and has a magnitude slightly in excess of 50 cm^{-1} .

The symmetric target discussed so far does provide strong velocity gradients, although the largest gradients occur in the up-down component of the velocity vector. Penetration is not as good as might be the case if the gradients in the right-left component were larger. There is at least one way to improve upon the symmetric approach considered so far, and in the next sequence of figures we illustrate the point. We consider a target with a geometry similar to the one discussed previously, and irradiate one side before the other. Infrared laser pumping pulses incident on the two sides are the same (90 psec FWHM of

Gain (cm^{-1})

A: $1.89\text{E}+01$

B: $2.34\text{E}+01$

C: $2.80\text{E}+01$

D: $3.25\text{E}+01$

E: $3.71\text{E}+01$

F: $4.16\text{E}+01$

G: $4.62\text{E}+01$

H: $5.07\text{E}+01$

Figure IV-26 Gain on the $4^1\text{F}-3^1\text{D}$ Transition of Ne IX for the Symmetric Target

We assume isotropic resonance radiation emission over the length of the flashlamps to the left and to the right of the laser medium; constant in time at a value of 0.015 photons/mode. Radiation transfer was avoided as described in the text for reasons discussed in section C, rather the local pump field was computed assuming no absorption takes place (due to the laser medium--the shields are infinitely absorbing) along the path between a point and the flashlamps. At the target center 80% of the "sky" is filled with flashlamp and hence the pump strength is $80\% \times 0.015$. Part of the variation of gain over the target is due to differences in local pump strength, but some is also due to differences in local concentrations of the helium-like ion due to variations in ionization rates.

peak intensity $6 \times 10^{14} \text{ W/cm}^2$ (as was employed in the previous calculations), and a delay of 25 psec has been added to one. In principle different pulse shapes and intensities might be delivered, and the flashlamp and filter systems might differ. We have not explored the possibilities yet, and for present purposes we use the same flashlamp and filter system as before.

The idea behind the target is to set up velocity gradients directed more through the central restriction than parallel to it. Pump radiation should be able to penetrate more easily in this case. In Figure IV-27 we show isocontours of the electron density. The upper part of the target is irradiated first and the snap-shot is for 130 psec (as well as following figures) where the peak of the prompt pulse occurs at 135 psec (similar to the illustrations shown earlier in this section). The gradients are stronger than before as the lower half of the problem is largely shadowed for early times. The target is no longer four-fold symmetric. The upper half is narrower; the idea being that the velocity gradients in the tranquil region will be minimal (the object is to create velocity gradients in the lower half of the target) and so there is no reason to make it thicker than the distance that the suprathermal electron driven shock would travel by the time lasing is to occur. In the lower half of the target we have left room for a delayed shock to propagate and for the laser region to occur.

Electron Density
($1/\text{cm}^3$)

A: 0.97×10

B: 1.05×10

C: 1.11×10

D: 1.17×10

E: 1.22×10

F: 1.28×10

G: 1.34×10

H: 1.39×10

I: 1.45×10

J: 1.51×10

Figure IV-27 Electron Density for the Asymmetric Target

In a calculation similar to the one described for the symmetric case we have simulated the behavior of an asymmetric target. The mesh included 135 points in one plane and reflective symmetry was used. The ionization has not proceeded as far as in the symmetric case as is expected, although the maximum electron density calculated is quite close to that for the earlier calculation. In this calculation only the momentum equation was integrated; however, with the large velocities obtained the changes in local ion density would be noticed if the mass equation were integrated.

Figure IV-28 shows the velocity flow as a function of space, and in Figure IV-29 we show velocity isocontours where the velocity is measured in units of the local thermal ion velocity. As is noted from the latter figure the velocity gradients are stronger than in the symmetric case. We shall not press further into an examination of other properties of the target, as in many ways the results are very similar to those of the symmetric case.

Figure IV-28 Velocity Flow for the Asymmetric Target

From this figure one can see the general direction in which fluid motion occurs. An arrow is located at each meshpoint and is pointed in the direction in which a local fluid cell is flowing. The length of each arrow is proportional to the magnitude of the local fluid velocity. As has been discussed in section C, the zoning around the vertices sometimes causes some numerical problems, in this example is no exception. The target dimensions in this case are similar to those of the symmetric target. 420μ is the width of the target along the flashlamp side, and the vertices are again separated by 84μ . The width of the upper section is 40μ and that of the lower is 70μ .

Figure IV-29 Isocontours of Velocity Normalized to the Local Thermal Ion Velocity

On the left is the x-directed (right-left) component of the normalized velocity and on the right is the y-directed (up-down) component. The range of the graphs are restricted so as to focus on the lower part of the target where lasing would most likely occur. Although the contours range between -10 and $+10$, in both plots values clearly go higher and lower. Much of the raggedness of the plot is due to the presence of shadows when a limited number (40) of angles were used in the angular (θ) integration for the photon spectrum. The results shown above are some of the more important of this section as they illustrate that (at least theoretically) usefully large and properly directed velocity flows can be achieved through the present approach to X-ray laser design.

E. Laser Design in the Case of the Ne IX 3-2 Laser

Having devoted considerable discussion to the design of the Ne IX 4-3 laser, we now turn to the problem of laser design in the case of resonantly pumped 3-2 lasers. The design of the higher energy (the Ne IX 3^1D-2^1P line is at 151 eV and the Ne X $3d-2p$ lines are at 189 eV) lasers is more difficult than in the case of the 55 eV laser scheme, and the purpose of this section and the following one is to comment briefly on some of the issues involved which cause the difficulties.

The 3-2 scheme is illustrated in the case of helium-like Ne IX in Figure IV-1. Resonant line radiation drives population from the $1s^2\ ^1S$ ground state to the $1s3p\ ^1P$ level, and the transition energy is 1074 eV. Collisional de-excitation to the $1s3s\ ^1S$ and $1s3d\ ^1D$ states occurs, resulting in gain on the $1s3s\ ^1S-1s2p\ ^1P$ and $1s3d\ ^1D-1s2p\ ^1P$ transitions. Inversion of the $1s3p\ ^1P-1s2s\ ^1S$ transition does not occur, and the reason for this is that although $1s2p\ ^1P$ radiative decay is rapid, the $1s2s\ ^1S$ state is coupled only weakly to the $1s2p\ ^1P$ state (this collisional coupling is the dominant decay mode of the $1s2s\ ^1S$ state). The gain of the $1s3d\ ^1D-1s2p\ ^1P$ transition is usually larger than that of the $1s3s\ ^1S-1s2p\ ^1P$ transition under optically thin conditions, due to the larger gf value involved (where g is the statistical weight of the lower state, and f is the absorption oscillator strength); however, radiation trapping of the $1s^2-1s2p\ ^1P$ line destroys the gain on the $1s3d\ ^1D-1s2p\ ^1P$ transition

before it does on the $1s3s\ ^1S-1s2p\ ^1P$ transition. Present interest will be focused primarily on the $1s3d\ ^1D-1s2p\ ^1P$ line at 151 eV.

In the case of the 55 eV laser, the initial neon density was chosen through optimization of the small signal gain; namely, that there was a local maximum in the gain near $3 \times 10^{18}/\text{cm}^3$ of initial neon gas. Further analysis of the target design showed that a choice of $2 \times 10^{18}/\text{cm}^3$ as an operating point was quite reasonable, as the absorption length for the resonant pump radiation was comparable to the effective width of the laser medium at the peak of the infrared laser pump pulse. The electron density of the laser medium was less than the suprathermal electron density likely to be present, which allowed coupling between the momentum density of the hot electrons and that of the ions, resulting in the development of very strong ion fluid velocity gradients. These gradients allowed improved penetration of the resonance pump radiation, and the target design resulting looks quite promising given that 0.015 photons/mode can be delivered at the $1s^2-1s4p\ ^1P$ transition energy.

In the present case of the 151 eV laser, the maximum gain does not occur below an initial neon density of $10^{19}/\text{cm}^3$, as is shown in Figure IV-30a. Selection of an initial neon density is no longer determined by simple gain optimization over one parameter, rather other physical constraints must be imposed. One constraint might be to require that the suprathermal electron density exceed the electron density of the laser medium, which limits the initial neon density

Figure IV-30a Small Signal Gain as a Function of Initial Neon Density for the $1s3d\ ^1D-1s2p\ ^1P$ Line at 151 eV

We have calculated the transient ionization and inversion kinetics in neon at several densities, and have plotted the results for peak gain achieved at each density as a function of the initial neon density. The spectrum was that of the center of the symmetric target described in the last section, and the $1s^2-1s3p\ ^1P$ line was pumped with 0.015 photons/mode radiation intensity that was constant in time. Proximity heating was included in this and all other calculations in this section and section F. At low density, the gain rises more rapidly than linearly as a function of neon density, due to the fact that there is more neon (and hence more inversion density if the fractional value were constant), and that the $1s3p\ ^1P$ to $1s3d\ ^1D$ coupling is improved by the increased number of electrons present (hence raising the fractional inversion density). The maximum in the gain curve occurs above $10^{19}/\text{cm}^3$, at which point electron collisional destruction and increased ion heating counters any benefit derived from the addition of more neon.

to below about $3 \times 10^{18}/\text{cm}^3$ when the infrared laser pump intensity is $6 \times 10^{14} \text{ W}/\text{cm}^2$. This constraint could be exceeded at the expense of losing the strong velocity gradients described in the previous section. That this would be ill-advised can be seen from Figure IV-30b, wherein the $1s^2-1s2p \ ^1P$ absorption coefficients are shown. The absorption length of this line is near 3μ at the time of peak gain for a neon density of $2 \times 10^{18}/\text{cm}^3$, and decreases roughly linearly with the inverse neon density. A few optical depths can be tolerated on the trapped $1s^2-1s2p \ ^1P$ transition, but past five optical depths severe problems occur. At neon densities above about $3 \times 10^{18}/\text{cm}^2$, where velocity gradients are much weaker, the width of the laser medium would be constrained to be below 10μ wide lest severe trapping spoil the laser gain. Such a severe requirement is unreasonable for many reasons, including the facts that the tolerance is too tight to allow target construction given available resources, and that with the factor of five spatial inhomogeneities of the infrared laser pump beams present, there can be no assurance that a large fraction of the laser would have the correct laser medium width at the same time given that 90 psec pulses are used.

An initial neon density between $2 \times 10^{18}/\text{cm}^3$ and $3 \times 10^{18}/\text{cm}^3$ is probably a reasonable operating point given the assumptions made about the suprathreshold electron densities present at a pump intensity of $6 \times 10^{14} \text{ W}/\text{cm}^2$. We shall choose $2 \times 10^{18}/\text{cm}^3$ as an operating point for present purposes, and note with delight that much of the analysis

Figure IV-30b Time-dependent Absorption Coefficients of the $1s^2-1s2p\ ^1P$ Transition at Different Neon Densities

From the same calculation as for the results shown in the previous figure, we show here the absorption coefficients at line center for the $1s^2-1s2p\ ^1P$ helium-like Ne IX resonance line. The absorption coefficients peak earlier in time at the higher densities, corresponding to the more rapid transient ionization on earlier production of the helium-like $1s^2\ ^1S$ state. Of interest in this figure are the magnitudes of the coefficients. At $8 \times 10^{18}/\text{cm}^3$, the absorption coefficient peaks above $10^4\ \text{cm}^{-1}$, corresponding to an absorption length under 1μ . There is no way that this author can envision to design a laser using the present approach which such a high absorption coefficient as the trapping problems are so severe. We shall choose to work at $2 \times 10^{18}/\text{cm}^3$ where the length is 3μ , and try to establish sufficiently large velocity gradients to relieve the trapping problem.

(especially the two-dimensional hydrodynamics studies) of the previous section can be applied to the problem of target design implementing the present scheme. We assume in the following that a target similar to the asymmetric one discussed at the end of the last section will be used for the design in this case, and in what follows, focus on the kinetics issues which are important in the determination of the gain. Before leaving the topic; however, we note that this scheme can benefit from the presence of more suprathermal electrons (i.e.; if there were more hot electrons present, one could work at a higher neon density, in which case the small signal gain for the corresponding optically thin problem is higher), if the radiation trapping of the $1s^2 - 1s2p \ ^1P$ line could be successfully reduced. One approach to this problem is to increase the infrared laser intensity incident on the flashlamp. We offer as speculation an alternate approach. Dense matter can scatter hot electrons, and therefore one might arrange to have electron reflectors outside of the target to help direct outgoing suprathermal electrons back into the laser medium. We have not calculated an example implementing this approach, and so we cannot say yet whether it works or not.

We now turn to a discussion of some of the population kinetics issues. The transient ionization proceeds much the same as in the analogous example considered in section D (see Figure IV-11b), which might be expected since the conditions are identical with the exception that the $1s^2 - 1s3p \ ^1P$ line is being pumped rather than the $1s^2 -$

$1s4p\ ^1P$ line (in both cases pumping occurs at 0.015 photons/mode). The helium-like sequence is not depleted so rapidly here due to the slower ionization and excitation (followed by ionization) rates of the M-shell levels as compared with the N-shell levels. In Figure IV-31a we show the fractional level populations of the L-shell and M-shell states of Ne IX. The helium-like $1s^2\ ^1S$ population peaks at about 60% of the total neon population, hence pumping the $1s^2 - 1s3p\ ^1P$ line with 0.015 photons/mode would result in a $1s3p\ ^1P$ fractional population of 0.027 if the only processes involved were photoexcitation and radiative decay. The calculated fractional population peaks near 0.020, which indicates the relative importance of the other processes involved (in the next section when we consider the analog of the scheme in the hydrogen-like sequence, we shall find that the relative losses are much greater). Collisional transfer of population to the $1s3d\ ^1D$ upper laser state is slow, leading to a $1s3d\ ^1D$ fractional population that peaks near 0.005. The relative values of the populations of the two states is an indication of the weakness of electron collisional coupling at high electron temperatures (if the states were in relative equilibrium, their populations would be proportional to their statistical weights, which gives a ratio of 5/3 --- instead, the ratio comes out to be about 1/4). The gain resulting is shown in Figure IV-31b, and one notes that the peak value of 34 cm^{-1} is lower than in the 55 eV laser calculations, and is sufficiently low that when other effects are included (primarily

Figure IV-31a Fractional singlet L-shell and M-shell levels in the case of $1s^2 - 1s3p\ ^1P$ pumping

A calculation of the transient ionization and the inversion kinetics of neon with $1s^2 - 1s3p\ ^1P$ resonant line pumping was carried out under a set of assumptions very similar to those made in the kinetics calculations reviewed in section D. In this figure are the fractional level populations of the important excited singlet states, and in the following figures are other quantities computed in the same calculation. The initial neon density is $2 \times 10^{18} / \text{cm}^3$, and the X-ray spectrum from the center of the symmetric target described in section C and section D was used to drive the transient stripping of the neon. The $1s^2 - 1s3p\ ^1P$ resonant pump radiation was kept constant at 0.015 photons/mode, that value being chosen to allow comparison with the 55 eV laser results of the last section and those of the 189 eV laser discussed in the next section.

Figure IV-31b Small signal gain for the Ne IX 3-2 laser scheme.

In this figure we show the gain on the $1s3d\ ^1D - 1s2p\ ^1P$ line at 151.3 eV and on the $1s3s\ ^1S - 1s2p\ ^1P$ line at 149.8 eV. From the previous figures one observes that the $1s3s\ ^1S$ and $1s3d\ ^1D$ levels have nearly the same population (the 1D level has the higher population), while there is a large disparity in the small signal laser gains occurring on the two lines. This result is due to the much stronger oscillator strength of the $1s2p\ ^1P - 1s3d\ ^1D$ transition. The gain persists for roughly 30 psec, and is limited in time by the depletion of the helium-like $1s^2\ ^1S$ population. The gain should be above 50 cm^{-1} for a SHIVA laser target (because of the 4 mm length constraint), and a pump intensity of 0.022 photons/mode would be required to achieve this in the optically thin limit. A higher value would be needed when trapping is included, although the factors determining it are a function of other target conditions, and have not been analysed yet.

the reduced penetration of the $1s^2 - 1s3p \ ^1P$ pump radiation when absorption is included, and the strong trapping of the $1s^2 - 1s2p \ ^1P$ radiation which may occur) then a laser experiment corresponding to this model would probably fail if the length were constrained to 4 mm. Raising the resonant pump radiation intensity can help (at least in theory, where such is possible through simple modification of a generator deck), and a value of 0.025 photons/mode may prove to be sufficient if the trapping problem is handled well, as the gain is linear in the pump intensity here. The stimulated emission cross section is shown in Figure IV-31c, and the absorption coefficients of the K-shell resonance lines are shown in Figure IV-31d.

The weakness of the $1s3p \ ^1P$ to $1s3d \ ^1D$ electron collisional coupling is responsible for the low gain achieved at 151 eV. A pump intensity of 0.015 photons/mode is considered to be relatively strong, at least by this author, and the basic problem is that the radiation is not being used efficiently. Ion collisional transfer occurs at higher ion temperatures; for example, proton collisions are effective above 50 eV, which is considerably higher than the ion temperatures calculated for this problem. If a strong infrared laser were tuned to the transition energy (0.46 eV) and its beam directed axially down the length of the laser target within the laser medium, then the slow radiative decay of the transition could be stimulated to high enough rates to be noticed in the kinetics. Were the transition stimulated such that decay occurred in 0.5 psec, then the peak calculated gain

Figure IV-31c Stimulated Emission Cross Sections for the Ne IX 149 eV and 151 eV Lines

From the same calculation as for the previous two figures, we have plotted here the stimulated emission cross sections at line center for the two singlet 3-2 inverted transitions. The points to note here are that the cross sections are not strongly dependent on the ion temperature (the increased line broadening due to the rising ion temperature is primarily responsible for decrease in cross sections here), and that the $1s3d\ ^1D-1s2p\ ^1P$ cross section is about an order of magnitude larger than the $1s3s\ ^1S-1s2p\ ^1P$ cross section (due to the difference in oscillator strengths). From Figure IV-31a one observes that the $1s3s\ ^1S$ and $1s3d\ ^1D$ levels have comparable populations, and it is the difference in cross sections which results in the disparity in gains calculated for the two lines.

Figure IV-31d Resonant absorption coefficients of the helium-like Ne IX K-shell resonance lines.

In the design of the 55 eV laser, the penetration of the $1s^2 - 1s4p \ ^1P$ pump radiation could not occur were the laser medium to remain stationary, due to the small absorption length involved (25μ). For this reason we pursued the use of large velocity gradients to help matters. The same problem occurs again here, except that the problem is aggravated as the absorption length of the $1s^2 - 1s3p \ ^1P$ line is 15μ . Trapping of the $1s^2 - 1s2p \ ^1P$ transition is a serious problem due to the 3μ absorption length which occurs, and it is not yet known whether the strong velocity gradients produced within the asymmetric target described in the last section are sufficient to relieve this trapping.

exceeds 100 cm^{-1} . The time-dependent gain and fractional level populations are shown under these conditions in Figure IV-32. The conclusion which can be drawn from this is that a scheme involving the use of an additional laser to stimulate one-photon decay of the $1s3p \ ^1P-1s3d \ ^1D$ transition would be an excellent approach were it not to impose unreasonable demands of the experimental group involved. A factor of three in gain could easily make the difference between a successful SHIVA experiment and a long wait for more powerful infrared driven lasers. The situation regarding two-photon Raman gain has not been examined yet.

The asymmetric target discussed in section C may be used for a target design (aside from differences in the flashlamp material, which in this case is likely to be either nickel or cobalt, depending on which has the preferred pump spectrum). The velocity gradients are considerable; however, the calculation of the gain including line radiation transfer has not been carried out for this target as discussed in section C. It is not known whether the velocity gradients are sufficiently large to relieve the $1s^2-1s2p \ ^1P$ radiation, which is the key problem in the successful operation of this type of laser. The absorption length is 3μ at the time of peak gain, and the velocity gradient is large enough to result in a thermal Doppler unit shift every $3-5\mu$ in the two transverse velocity component. The situation looks quite hopeful, but it is not yet known whether the conditions are favorable enough for high gain to develop.

Figure IV-32 Gain and Fractional Level Populations for the 151 eV Laser in the Case of $1s3p\ ^1P-1s3d\ ^1D$ Pumping

The weak coupling of the $1s3p\ ^1P$ and $1s3d\ ^1D$ levels is responsible for the relatively low gain of the 151 eV line. We have redone the calculation of the last sequence with additional rates included between these levels to model the effects of optical pumping of the 0.46 eV transition with an infrared laser. The coupling rate was taken to be $2 \times 10^{12} \text{ sec}^{-1}$ in the $1s3p\ ^1P-1s3d\ ^1D$ direction, and 0.6 times this value for the reverse process. The infrared laser radiation intensity should be about 750 MW/cm^2 for this purpose. In the absence of this radiation, the $1s3p\ ^1P-1s3d\ ^1D$ line develops a gain of about 10 cm^{-1} (the absorption due to inverse bremsstrahlung is much less if the electron temperature exceeds 100 eV), and in principle could supply the requisite pump photons if the local reflectivities could be made high enough.

F. Laser Design in the Case of the Ne X 3-2 Laser

We conclude this chapter with a brief examination of the issues involved in an implementation of the resonant line pumping scheme in the hydrogen-like isoelectronic sequence. The 3-2 transitions occur at 189 eV (the scheme is shown in Figure IV-33). Line radiation resonant with the 1s-3p transitions causes photopumping of the 3p states, and electron and ion collisional processes transfer population to the 3s and 3d states. Due to the rapid decay of the 2p states, population inversions can develop on the 3s-2p and 3d-2p transitions at 189 eV, although the gain is much weaker on the 3s-2p lines. The $1s_{1/2}$ - $3p_{1/2}$ line is at $10.23964\text{\AA}$ and the $1s_{1/2}$ - $3p_{3/2}$ line is at $10.23850\text{\AA}$, and as the difference in energy between the lines is 0.13 eV, the two transitions may be pumped with different intensities. The details are dependent on the characteristics of the resonant flashlamp line involved, and in Table IV-1b candidate lines are listed in gallium and zinc. It is likely that copper and nickel will have better pump lines as discussed in section A; however, the requisite spectroscopic analysis remains to be published and so one cannot say at this time if it is so. In the calculations reviewed here we have assumed that both lines see the same pump field, and have kept the pump intensity constant at 0.015 photons/mode (even though the gain resulting is too low to be useful for a SHiVA design where the laser length is limited to 4mm) to allow comparison with results presented in the last two sections.

Figure IV-33 The Hydrogen-like Ne X 3-2 Resonantly Pumped Laser Scheme

Photoexcitation of the 1s-3p transitions populates the $3p_{1/2}$ and $3p_{3/2}$ levels, and electron and ion collisional processes causes transfer of population to the 3s and 3d states. Due to the rapid radiative decay of the 2p levels, inversions can develop on the 3s-2p and 3d-2p transitions. In the analog of the scheme in the helium-like sequence (discussed in the last section), six singlet levels were involved. In the present case we are dealing primarily with nine doublet levels, and as a result the scheme has quite different properties, especially under conditions described here. The ratios of the statistical weights are different and there is now only one electron in the K-shell. These differences lead to lower gain (keeping the pump intensity constant) and less radiation trapping problems.

Many of the issues involved in the design of this laser are similar to those discussed previously. As was the case in the design of the 151 eV laser, the maximum gain as a function of initial neon density occurs well above $3 \times 10^{18}/\text{cm}^3$ (the curve of interest is shown in Figure IV-34a, and one notes that the gain curve for the $2p_{1/2}-3d_{3/2}$ transition is approaching a broad maximum towards the high density end of the plot, the peak of which occurs above $10^{19}/\text{cm}^3$). The density of $3 \times 10^{18}/\text{cm}^3$ is of interest, because above this neon density, the electron density of the laser medium is likely to exceed the suprathermal electron density, in which case no large velocity gradients would be produced. The penetration of pump radiation resonant with the $1s_{1/2}-3p_{3/2}$ transition and the escape of $1s_{1/2}-2p_{1/2}$ radiation is less of a problem than in the case of the analogous lines in the helium-like scheme (by a factor close to 6, as shown in Figure IV-34b), and for this reason one might attempt a laser design wherein the velocity gradients are not maximized (i.e.; one might work in the high neon density regime where the laser medium density exceeds the hot electron density, and the large suprathermal electron momentum density is not coupled into the ion fluid momentum density). Such an approach might be implemented by using an initial neon density near $10^{19}/\text{cm}^3$; however, from examination of the results shown in Figure IV-34a, one notes that the increase in peak gain achievable is less than a factor of two

Figure 34a Parameter Study of the Gain at 189 eV with Initial Neon Density as a Parameter

We have calculated the small signal laser gain on the 3-2p lines at different values of the initial neon density, and on the left is shown the peak values of gain at each density plotted as a function of neon density. On the right are time dependent results for the high gain $3d_{3/2}-2p_{1/2}$ line. The gain on the two lines having the $3d_{3/2}$ level as an upper state is a relatively slow function of the density, whereas the $3d_{5/2}-2p_{3/2}$ line varies more rapidly. This is due to the tight ion collisional coupling of the $3p_{3/2}$ and $3d_{3/2}$ states, while the $3d_{5/2}$ level is populated by electron collisional excitation. The latter process is relatively weak and hence shows a dependence on the density. From the time-dependent results, one notes that the gain peaks earlier at the higher densities, which is due to the faster electron collisional stripping of the neon because of the increased number of electrons present.

Figure IV-34b Absorption Coefficients of the $1s_{1/2}-2p_{1/2}$ Transition at Different Neon Densities

The absorption coefficients for the $1s_{1/2}-2p_{1/2}$ K-shell resonance line were taken from the same calculation as the results shown in the previous figure. The time dependences are roughly the same in the two cases, which is as expected since both are determined from the population of the ground state of the sequence. Of major interest is the magnitude of the coefficients (the peak values range from 200 cm^{-1} to 650 cm^{-1}) as they are quite small, especially when compared against the analogous transition in the helium-like sequence. From the previous figure we have established that the scheme in the hydrogen-like sequence has low gain, and here we show that the trapping is no where as bad as was the case before. These factors are the primary differences between the two schemes from a designers point of view.

higher than what occurs at $2 \times 10^{18}/\text{cm}^3$, and the problems associated with radiation transfer effects in a relatively velocity-free laser medium probably offsets the increase in gain won by moving into this regime. We have not explored the high density regime for this laser scheme, and in what follows we shall be concerned with an operating point at $2 \times 10^{18}/\text{cm}^3$, in agreement with conditions selected for the 55 eV and 151 eV lasers.

In the next sequence we review a calculation of the population kinetics in the case of $2 \times 10^{18}/\text{cm}^3$ neon density. The pumping spectra is taken at the center of the symmetric target discussed in section D, and in many ways the calculation is similar to those presented in the last two sections. The primary difference is that the hydrogenic 1s-3p lines were resonantly pumped (at 0.015 photons/mode) rather than the helium-like 1-4 or 1-3 singlet lines as was the case in the previous two sections. The transient ionization proceeds much the same as in the earlier cases until the neon hits the final sequences. Where, in the case of the 55 eV laser (see Figure IV-12), the fractional population of the helium-like sequence peaked near 55% of the total neon population, here it reaches 70% (see Figure IV-35). Before, the hydrogenic sequence peaked above 70%; here it doesn't even reach 50%.

In the next figure (Figure IV-36), we present the fractional populations of the L-shell and M-shell states of the hydrogenic sequence. This figure requires some comment, as herein is contained

Figure IV-35 Fractional Populations of the Final Isoelectronic Sequences for the 189eV Laser

In this figure, and in the next few figures, we examine the population kinetics at an initial neon density of $2 \times 10^{18}/\text{cm}^3$. The calculation is similar to analogous calculations presented for the 55 eV and 189 eV laser schemes, including the use of the X-ray spectrum at the center of the symmetric target and the use of a local electron temperature which follows the suprathermal electron temperature. The dramatic effect which resonant photoexcitation with a strong resonant radiation field has on the kinetics can be appreciated through a comparison of this figure and its analog in the 55 eV laser section (see Figure IV-12). The peak hydrogenic population drops from 70% in the case of the 55 eV laser to 45% in the present case.

Figure IV-36 Fractional Populations of the L-shell and M-shell States of Ne X

The levels involved in the pathway leading to highest gain are the $3p_{3/2}$, $3d_{3/2}$ and $2p_{1/2}$ states. We have assumed in the calculation that strong ion collisional coupling occurs between the $3p_{3/2}$ and $3d_{3/2}$ states, and have used a rate coefficient of $5 \times 10^{12} \text{ sec}^{-1}$ for the process. The level populations are somewhat dependent on the value chosen, while the gain on the $3d_{3/2}$ - $2p_{1/2}$ transition is less so. When the rate is fast, better coupling occurs; however, the added homogeneous broadening of the laser transition offsets the additional inversion density resulting from the improved coupling. Here the $3p_{3/2}$ population is higher than that of the $3d_{3/2}$ state, and since the statistical weights are the same for the two levels, they are observed to be in approximate relative equilibrium. The $3d_{5/2}$ population is much less than that of the $3d_{3/2}$ state, hence the $3d_{5/2}$ - $2p_{3/2}$ gain is much lower.

the explanation of why the gain for this scheme is as low as it is (for example, the gain for the 3-2 scheme in the helium-like ion is 50% higher without laser-assisted pumping, and a factor of five higher with it). In section B we discussed the problem of ion collisional coupling between the $3p_{3/2}$ and $3d_{3/2}$ levels which are separated by only 2.28×10^{-4} eV. The ions behave as a Coulomb fluid, and we proposed that ion collisions took place through phonons present in the fluid. It appears that ion collisional transfer occurs when a neighboring ion is close (as occurs in the compression phase of a local phonon wave as seen by a target ion), and doesn't occur when the neighboring ions are sufficiently near equilibrium such that no large dipole component of the local electric field is produced. The frequency of the phonon waves is near 1.5×10^{12} radians/sec. We have examined the inversion kinetics with the coupling rate being 1.5×10^{12} /sec and 5×10^{12} /sec, and find that the gain is nearly identical in the two cases, although some difference occurs in the kinetics. The additional broadening which occurs when the rate is 5×10^{12} /sec offsets the gain in upperstate population such an increased rate produces, relative to the case at 1.5×10^{12} /sec. In the calculations discussed in this section, we have used 5×10^{12} /sec as a coupling rate. The result is that the $3p_{3/2}$ and $3d_{3/2}$ levels are tightly coupled in our calculations, while the $3p_{3/2}$ and $3d_{5/2}$ levels are only weakly coupled by electron collisional excitation. The $3p_{3/2}$ - $3d_{5/2}$ transition is

analogous to the $1s3p\ ^1P-1s3d\ ^1D$ transition in the helium-like sequence in that they represent bottlenecks in the inversion kinetics (the $3p_{3/2}-3d_{3/2}$ transition is analogous to the $1s3p\ ^1P-1s3d\ ^1D$ transition when optical pumping occurs, as described at the end of section E). Similarly, coupling between the $3p_{1/2}$ state and the $3d$ states is weak. The pathway for the production of highest gain is then photoexcitation of the $1s_{1/2}-3p_{3/2}$ line, followed by ion collisional coupling of the $3p_{3/2}$ and $3d_{3/2}$ states, leading to high gain on the $3d_{3/2}-2p_{3/2}$ transition. The small signal gains are shown in Figure IV-37, and one observes that the peak gain on the $3d_{3/2}-2p_{3/2}$ transition exceeds 20 cm^{-1} , well above the gains achieved on the other lines. The $1s_{1/2}$ population peaks near 45% of the total neon population, and with 0.015 photons/mode pumping the $1s_{1/2}-3p_{3/2}$ transition, the $3p_{3/2}$ fractional population would reach 1.35% if only photoexcitation and radiative decay on the $1s_{1/2}-3p_{3/2}$ line occurred. The calculated maxima shown in Figure IV-36 is about half that, indicating that coupling to other levels (primarily the $3d_{3/2}$

Figure IV-37 Small Signal Gain on the 2-3 Lines of Ne X at 189 eV

The small signal gains for the 3s-2p and 3d-2p transitions for the calculation being discussed are shown above. As mentioned previously, the gain on the 3s-2p lines is quite low (too low to be of use in the present approach to target design), due to the small 2p-3s oscillator strength. The inversion density is within a factor of two of that for the 2p-3d transitions; however, the stimulated emission cross section is smaller by several fold. The $2s_{1/2}-3p_{3/2}$ line is quite close to the $2p_{1/2}-3d_{3/2}$ line, the mismatch being about 0.020 eV. The peak absorption coefficient of the $2s_{1/2}-3p_{3/2}$ line peaks near 50 cm^{-1} , and because of the proximity in energy of the two lines, one is interested in the possible overlap between the lines. Some overlap occurs; however, the absorption coefficient of the $2s_{1/2}-3p_{3/2}$ line at the $2p_{1/2}-3d_{3/2}$ energy in this calculation peaks near 1.2 cm^{-1} , hence the lines are nearly distinct.

state) places quite a drain on the $3p_{3/2}$ population. The $3p_{3/2}$ and $3d_{3/2}$ levels have comparable populations, while the $3d_{5/2}$ level is relatively underpopulated. The $2s_{1/2}$ fractional population exceeds 0.011, well above that of the other levels, and therefore the 2s-3p lines do not invert.

The reasons why the gain is lower for the hydrogenic 3-2 scheme than that of the helium-like 3-2 scheme can now be outlined. If the 3p to 3d transition were fast enough to cause near statistical population between the two states (in the hydrogenic case, this is applied only to coupling between the $3p_{3/2}$ and $3d_{3/2}$ states), and if further all processes other than photoexcitation and radiative decay of the 1-3 transitions were included, the $1s3d\ ^1D$ state would be populated 2.5 times as much as the $3d_{3/2}$ state, given that the respective ground states were equally populated (i.e., the ratio of the statistical weights of the $1s3d\ ^1D$ state and the $1s^2\ ^1S$ state is five, whereas the ratio for the $3d_{3/2}$ and $1s_{1/2}$ states is two, and under the simplifying conditions described above, the upper laser state populations would be proportional to these ratios). The pump rate up to the $1s3p\ ^1P$ state is stronger relative to the losses than the pump rate to the $3p_{3/2}$ state relative to its losses, hence the fractional population of the upper laser state is decreased further. The peak fractional inversion densities are 1.4% and 0.4% for the helium-like and hydrogenic cases respectively, and so the natural statistical weight favoring of 2.5 is boosted to a factor of

four when losses are included. The helium-like $1s3d\ ^1D-1s2p\ ^1P$ stimulated emission cross section is larger by 20% than the $3d_{3/2}-2p_{1/2}$ cross section, which explains the difference in small signal gain (in the last section we had calculated slightly over 100 cm^{-1} for the gain when the $1s3p\ ^1P-1s3d\ ^1D$ was tightly coupled, and here the gain is 22 cm^{-1}).

The lower gain of the hydrogen-like 3-2 laser scheme under these conditions is a real effect, and reflects primarily the differences in the statistical weights involved brought about by the difference between LS and j-j coupling. This result implies a requirement for a higher modal photon density of the resonance pump radiation to achieve the same gain under optically thin conditions as in the case of the analogous scheme in the helium-like sequence. Along with the bad comes some good in this case, as the severe problems associated with radiation trapping which occur in the helium-like sequence are very much relaxed for the hydrogen-like system (see Figure IV-38). One should not immediately conclude that one scheme is better than another (although one might be tempted because the hydrogen-like scheme achieves considerably lower gain in the optically thin limit); rather, that they are different. If the laser length constraint on SHIVA were not so stringent, the radiation transfer imposed constraints would make the hydrogenic system the more pleasant of the two to work with. In order to obtain a small signal gain of 50 cm^{-1} , however, a pump intensity of nearly 0.035 will be required.

Figure IV-38 Absorption Coefficients for the $1s-2p$ and $1s-3p$ Lines of Ne X

From the calculation of the previous few figures we have plotted the absorption coefficients of the hydrogenic K-shell resonance lines. As noted earlier, radiation trapping on the $1s_{1/2}-2p_{1/2}$ line is less of a problem than in the analogous case in the helium-like sequence, since here the relevant absorption length is 25μ (admittedly not as many optical depths can be tolerated; even so, the difference is quite marked). Trapping of the $1s_{1/2}-2p_{3/2}$ line is not of great concern for the scheme, since the gain on the $3d_{5/2}-2p_{3/2}$ line is so low. Penetration of $1s-3p_{3/2}$ pump radiation is not a problem as the absorption length when absorption is at its worst is 80μ .

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V. Conclusions and Suggestions for Further Work

We have developed a general purpose computer code (XRASER) to aid in the design and simulation of EUV and X-ray laser systems. We have designed a laser at 27 eV implementing a photoionization-pumped scheme involving preferential ionization of 2s electrons from Ne I leading to a 2s-2p self-terminating inversion in Ne II. The design has been analyzed using XRASER and the laser fusion code LASNEX. The small signal gain in the absence of lasing is calculated to exceed 50 cm^{-1} (hence noise produced at one end of the laser experiences in excess of 20 gain lengths amplification while transversing the laser axially as the maximum irradiation length on SHIVA is 4 mm) when the laser flashlamps are irradiated uniformly with 90 psec FWHM pulses of peak intensity 10^{14} W/cm^2 . This approach appears to be promising and may be suitable for an early experiment in a laser-pumped short wavelength laser program.

We have also considered designs implementing resonantly-pumped photoexcitation schemes in helium-like and hydrogenic neon where the candidate laser transition energies are 55 eV, 151 eV and 189 eV. Some analysis of the designs has been carried out using XRASER and LASNEX, and the results show that gains of 50 cm^{-1} may be attainable if the resonant line radiation field exceeds 0.015 photons/mode. Previous authors have suggested using K-shell lines to pump K-shell transitions, or else using 1-2 transitions to pump 2-4 lines. We have extended the search to L-M transitions

in 3 to 10 electron systems and have found many candidate resonances to pump 1-3 and 1-4 transitions in H-like and He-like low Z ions.

Whether or not the resonant pumping radiation exceeds 0.015 photons/mode under the conditions of interest is not yet known as the calculations are difficult and experimental results have not yet become available. Future work on this approach should most immediately be directed towards the accurate determination of the strength of the resonant radiation fields and their relative displacement in energy.

In the case of the 55 eV laser, difficulties were encountered in transporting the pump radiation into the laser medium as the constraints on target width imposed by target fabrication considerations lead to a target whose width exceeds the penetration depth of the pump radiation. One solution to this problem is to set up strong velocity gradients in the two lateral fluid velocity components as a function of space, the idea being that associated with the velocity flows are corresponding bulk Doppler shifts which effectively increase the penetration depth of the pump radiation. The use of large velocity gradients to relieve radiation trapping has been considered in the case of recombination schemes; however, the use of such has not been explored in the case of resonantly-pumped schemes. The use of suprathermal electrons in conjunction with ionization gradients set up between shadowed

and unshadowed regions of the laser medium to drive large velocity gradients is proposed, and calculations show that the effect is a strong one. Ion velocities well in excess of 10^6 cm/sec are calculated for reasonable design parameters leading to a severalfold increase of the penetration depth.

It follows from calculations presented here that the resonantly-pumped approach to short wavelength lasers appears promising. We have examined a simple implementation of the scheme, one which will hopefully lead to experiments which are straightforward and require a minimum of effort from the laser program involved.

We now consider some of the issues of the different chapters.

In the first chapter we have given a brief review of the field of short wavelength lasers. The laser schemes were examined by pumping mechanisms including collisional pumping, recombination pumping, photoionization and photoexcitation pumping and charge exchange. In the work discussed here we have focused on a photoionization scheme and three photoexcitation schemes. There are other possibilities in both cases; for example, photoionization of 2p electrons in sodium leads to a 3s-2p self-terminating transition as suggested by Duguay in 1967. This experiment would involve more technical problems due to the use of sodium vapor than the one discussed in Chapter III. Ionization of 2s electrons from fluorine or oxygen are possibilities, leading to inversions at lower energies. The problems here involve obtaining monatomic species or else

working the scheme on diatomic molecules (working with photoionization of molecules is an area relatively unexplored in the context of EUV lasers). Photoionization of K-shell electrons from neon can lead to inversions in the oxygen-like sequence due to a favorable splitting in the Auger decay of the fluorine-like ion with a K-shell hole. There appear to be many systems in which vigorous pumping by a short pulse of X-rays can lead to EUV self-terminating inversions. The power and risetime requirements will vary for the different systems and it is likely that lower energy inversions can be pumped at considerably less than 10^{14} W/cm².

Extension of the scheme to 2s-2p inversions in fluorine-like silicon and argon have been examined in detail (although not discussed in this work) and the results do not appear at all promising. The problems center around two areas; that the radiative decay rates are much faster at higher Z which is unfortunate for a self-terminating scheme, and that there are many more electrons around so that collisional quenching of the inversion is much more severe. We have examined 2s-2p inversions in higher isoelectronic sequences of neon, silicon and argon and again the scenario appears bleak. The only reasonable approach to make these schemes work in silicon and argon is to provide a mechanism to actively depopulate the lower laser level; for example, by resonantly pumping it uses a coincident line in the flashlamp. In the case of neon, 2s-2p transitions in the more highly stripped

sequences could be lased if the target and pulse characteristics would allow for more broadband pumping X-rays near 500 eV.

In the case of line pumping schemes, we have examined hydrogenic and helium-like ions due to their simplicity and relative ease of production, and treated neon here for several reasons, including favorable line coincidences and continuity from early chapters. During the course of the research we have examined carbon, nitrogen, oxygen and fluorine as well (details may be published elsewhere in the future). Estimates suggest that neon may be a reasonable goal given the capabilities of the SHIVA laser systems; that if one discards Na, Mg and Al as not being easily obtainable in gas form, then silicon (in the gaseous form of silane) would be the next candidate when increasing Z , and the transition energies involved are sufficiently high to guarantee that providing the requisite resonant pump radiation will be difficult. It seems to the present author that one might profitably pursue helium-like and hydrogenic inversions in elements from $Z=6$ (carbon) to $Z=10$ (neon) using target designs similar to those discussed in Chapter IV; however, given the presently available pump characteristics lasing silicon will be difficult. Inversions in other sequences are possible, including 4-3 inversions in neon-like systems and 5-4 inversions in nickel-like systems. These sequences in particular may be useful in that they represent natural resting places for the ionization balances; that is, they

are relatively easy to produce. Line-pumping schemes in the more complex ions have yet to be explored in any detail; however, they should prove to be interesting in terms of EUV and soft X-ray lasers.

The use of a low energy laser pulse to stimulate Raman gain or to selectively channel population appears to be an excellent idea given adequate resources in a short wavelength laser program. We have not pursued the problem here as the experimental problems which ultimately would be involved are much too formidable to seriously contemplate a laser-assisted approach at LLL.

Other approaches are possible. Recombination and charge exchange schemes are being analyzed elsewhere in the context of laser-pumped short wavelength laser research. The approach so far appears to result in insufficient gain to be of use on SHIVA given the 4 mm length constraint. If the constraint were relaxed, then these approaches would become far more attractive for local use. Schemes pumped through electron collisional pumping (notably the 3s-3p laser schemes) appear to be quite promising at low energy. We have examined the kinetics of neon-like Krypton for possible gain under Z-pinch conditions, and one idea which arose from this exercise was the use of suprathermal electrons to do collisional pumping; for example, in neon-like Krypton, resulting in gain at 83 eV. The electron collisional rate coefficient which dominates the inversion kinetics (the 2p-3p monopole process) peaks near

3 keV, which can effectively be pumped by hot electrons generated at the flashlamp when irradiated with 1.06μ light near 10^{14} W/cm². One can imagine using a target similar to the ones discussed in the present work with krypton gas flowing through, and upon pumping from above and below, the krypton strips down to the neon-like sequence (and more highly stripped sequences as well) at which point lasing might occur. We have carried out some calculations on this approach and the results appear promising. Indeed, these calculations in the previous draft constituted the main body of an extra chapter of this thesis. This work may be presented elsewhere in the future. This approach may be adaptable to target design for CO₂ laser pump sources as the longer wavelength radiation is better suited for hot electron production.

Suprathermal electrons can in principle be used in other ways as well; for example, hot electrons would cause strong excitation on a resonance line in a low density plasma, and if the resonance line were strongly trapped the one could build up a very bright radiation field within a few linewidths of the resonance line center. Effective temperatures of the radiation fields on the trapped lines could approach the hot electron temperature, given that the hot electrons numerically dominated the local unbound electron density, such as is thought to occur in the short wavelength laser designs discussed in previous chapters. Such bright X-radiation could be used to drive a resonantly pumped short wavelength laser

scheme. Chapline has independently proposed this approach as a way to achieve high gain between 500 eV and 2 keV in a SHIVA-pumped target design (Chapline, unpublished (1979)). Chapline restricted his attention to H-like and He-like resonance lines; however, our calculations have shown that the 2p-3s transition of neon-like ions have the potential to develop very bright radiation fields with fewer optical depths in a related study. The use of this approach to lower transition energies has been examined in the course of the present research in laser-pumped targets and one of the major problems involved include the production of low density pump material as the fourth row transition metals are hard to get as gases or high density vapors under experimental conditions. It was decided on the basis of interactions with the laser program that only the simplest of targets were ever likely to be fabricated; hence, the class of targets described in Chapters III and IV are the ones having received the most attention.

The second chapter was devoted to a discussion of atomic physics parameters for neon and some of its L-shell ions. Oscillator strengths were collected from results given in the literature, as were some photoionization cross sections. The distorted wave method and empirical methods were used to fill in data not found in the literature. Electron collisional rate coefficients were computed from cross sectional data assembled from several sources. Auger processes were examined. Future work along these lines might be

devoted to the inclusion of the M-shell states in the first few sequences. The electron collisional matrix of cross sections coupling the L and M shells has yet to be given in full although the capability exists in several places to compute it. One reason for studying the model described in Chapter II was to gain experience in working with larger models, and a simple L-shell, dipole collisionally coupled model represents a low effort level step in that direction. One goal of some of the more recent work which the author is engaged in is directed toward the construction of a general set of L- and M-shell models for the 3-10 electron ions. Much work is being devoted toward this goal in the atomic physics and plasma communities, and work is even beginning on analysis of kinetics models in M-shell (11-28 electron) systems.

In the third chapter we have examined the issues of target design in the case of the 27 eV laser. Molybdenum was chosen as flashlamp material and LiF for filter material. The flashlamp material choice was somewhat arbitrary and it would be profitable to optimize the Z in the vicinity of 42. LiF was a reasonable choice for filter material, however Li or LiH would be superior from a theoretical point of view (as can be explained only by one who has never spent time in a target fabrication laboratory trying to work with Li and LiH). If not too impractical, some thought might be devoted to devising filters composed of these materials, perhaps in a sandwich configuration surrounded by parylene. The

simulation of the flashlamp and filter system suffers from errors associated with the use of the XSNQ model for non-LTE in LASNEX. In particular, the oscillator strengths differ from hydrogenic values by factors approaching 2 in the 3-4 transitions (the real values are lower), and the electron collisional rate coefficients are underestimated by varying factors (of order 4) in the XSNQ model. Correction of these defects would probably increase the total 3-4 radiative output; however, as to whether the spectral calculation is in better agreement with reality is not yet known. A non-LTE model for molybdenum should be constructed and considerable effort was devoted toward the project in the course of the research. The basic problem involved, assuming that all atomic physics parameters were easily computable (which in fact is close to the case at present), is in selecting a finite and small number of states adequate to describe the kinetics unless one is prepared to do calculations on a several thousand state model. We have examined models where several hundred levels total were used, and where 115 of them were included in the kinetics at a given time. In principle the extension to 1000+ level models is straightforward although we do not yet have experience with these kinds of calculations.

Perhaps a more straightforward approach to the problem of flashlamp and filter behavior would be to measure the time dependent spectra experimentally, although such a project is at present beyond

the state of the art. Much effort was put into obtaining time independent spectral information through a disk irradiation experiment on ARGUS; however, to date the experiment has not been successfully completed.

The suprathermal electrons were assumed to displace the thermal electrons completely, although detailed calculations of this process have not been carried out. The inclusion of kinetics models into LASNEX may allow the electron transport package of Kershaw to handle the problem. The two dimensional suprathermal transport problem has not been carried out in any approximation yet, and if the experiments yield negative results, this may be one place to look.

The third chapter concludes with a discussion of laser kinetics and gain, and gains in excess of 50 cm^{-1} were calculated. The accuracy of the calculation is limited at present by uncertainties in the spectra, in the suprathermal electron distribution, and in the model. The spectral errors are probably the most severe. Future work may be devoted toward re-examination of these issues. The designs work better with faster infrared laser risetimes, and some thought might well be spent on finding ways to get a sharp risetime on the standard 90 ps SHIVA pulse (perhaps through use of a saturable absorber at the back end). Experimental results on the 27 eV laser may be available within the coming years.

In Chapter IV we have examined line-pumping schemes in helium-like and hydrogenic neon. The first issue examined included the

search for L-shell resonance lines. We have examined most of the available literature pertaining to line positions suitable for resonant line pumping of one- and two-electron systems, and such a search has yielded many lines. The suitability of a line for pumping is determined by its oscillator strength and statistical weights of the states involved, and by the isoelectronic sequence of the transition. So far most of the resonances come from Ne-like, F-like and O-like sequences from the one side and Li-like and Be-like sequences from the other side. The middle sequences are the least well known. Spectroscopic work devoted to the analysis and classification of term systems is currently facing the very complex nitrogen and carbon sequences, from which it is likely that several important line resonances will be found.

The determination of the radiative output of the L-shell lines which have been found to be resonant or near resonant requires further work. Unfortunately it is the more complex systems which appear to be endowed with the best resonant line candidates which makes theoretical analysis difficult. Construction of detailed non-LTE models is underway and should definitely be a major area of future work. Experimental determination of absolute spectral output of the L-shell lines, preferably with 10 psec time resolution, would be a very difficult but also very valuable area of future endeavor. Spectroscopic work in the area of finding line positions relative to the line to be pumped would be very helpful.

A kinetics model of neon and its ions was constructed. Work can always be devoted to improving kinetics models, extending kinetics models and adding more processes. On the one hand the model should be as detailed as possible for reasons of accuracy and range of applicability, and on the other hand, it is not reasonable to run too detailed of a kinetics model in a large many-zoned two dimensional simulation. Work should not only be devoted towards development of more complete models, but also to the development of simple models which have sufficient physics to be interesting but are economical. We have done some work on both types, and the calculations described in Chapter IV were all performed using a model of the former class. This type of modeling should be pursued for other systems wherein potential short wavelength laser schemes reside, and for a laboratory dealing as extensively in plasma creation and utilization as LLL, it is hard to understand why studies of non-LTE kinetics have for so long not been included as an integral part of the laboratory's research effort.

The two dimensional hydrodynamics here are not state of the art algorithms, but on the other hand the amount of computer time spent on them relative to other parts of the code is so small as to make that issue unimportant. General angle dependent two dimensional line radiation transfer is an area which has not been extensively explored and the approach taken here may be of interest in other areas. Unfortunately, in the present class of designs where such

large velocity gradients are of interest the line transfer of XRASER is not yet able to provide an adequate description of the trapped radiation fields. Partial redistribution in energy can be included in a simple manner (namely, assume that the radiation field on the last timestep is similar spectrally at a given point to the one at the current timestep, and use the shape for an explicit determination of the redistributed emission function) and some future work should be devoted to that project as it is important in laser design.

The design work and simulations which have been done so far are promising, and should be continued. Once the pump radiation fields are more accurately determined then it should be possible to design a laser and be confident that it will produce measurable signals. Targets irradiated by an early pulse on one side, thus setting up a strong velocity flow (such as described in Chapter IV) represents the most promising approach to target design in the case of resonantly pumped schemes as viewed by this designer. Such design work should continue, in conjunction with better design tools.

The implementation of non-LTE into LASNEX would certainly aid in many of the calculations, although much remains to be done before it would serve the same function as XRASER. I would recommend maintaining XRASER and continuing development on it as well as pursuing LASNEX non-LTE. There is a place for both codes at present in a short wavelength laser design effort.

Experimental work should have gone in conjunction with the design work for the past several years. The reasons that it has not in the past are obscure. Future work should definitely include the development of an ongoing and supported experimental program.

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