**An Experimental Investigation of Stimulated Brillouin
Scattering in Laser-Produced Plasmas Relevant to Inertial
Confinement Fusion**

By

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To my wife, Alison, and my daughter, Chelsea.

In Memory of
Shauna Michelle Carver
1970—1988

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Preface

Research devoid of motivation is of little intellectual value. On the contrary, research *driven* by motivation can greatly enhance our understanding of nature. Were the work contained herein of interest only to other scientists, its value would be severely diminished. Often, studies appear to be devoid of value simply because the scholar fails to motivate the problem in a manner befitting the viewpoint or background of the audience. The difference between *good* and *great* lies in the ability to explain complex concepts in simple terms. It is the hope and belief of the author that any person who so chooses, regardless of his or her own scientific knowledge, may achieve a greater understanding of the problem that this research attempts to address. Toward that goal, the first chapter is presented. The remainder of the dissertation follows a more traditional format, at a level befitting those at least initiated in the sciences.

Finally, for purposes of historical perspective, it is important to note the world political situation in place during this scientific investigation. This research effort began shortly before the rather abrupt end of the Cold War and the subsequent collapse of the Soviet Union as a governing body. Though ironic, it is certainly true that the search for peaceful uses of fusion have at least in the United States been funded primarily by military expenditures that became available because of the Cold War. It is likely that other countries' campaigns in inertial fusion were similarly motivated. This irony is

only enhanced when one realizes that the end of the Cold War has left the future of peaceful fusion research uncertain. The unpredictable path to fusion energy will depend more on world politics than upon fundamental merit or potential promise, as this regrettably appears to be an axiom of human nature.

Keith S. Bradley

Abstract

Despite the apparent simplicity of controlled fusion, there are many phenomena which have prevented its achievement. One phenomenon is laser-plasma instabilities. An investigation of one such instability, stimulated Brillouin scattering (SBS), is reported here. SBS is a parametric process whereby an electromagnetic wave (the *parent wave*) decays into another electromagnetic wave and an ion acoustic wave (the *daughter waves*). SBS impedes controlled fusion since it can scatter much or all of the incident laser light, resulting in poor drive symmetry and inefficient laser-plasma coupling.

It is widely believed that SBS becomes convectively unstable—that is, it grows as it traverses the plasma. Though it has yet to be definitively tested, convective theory is often invoked to explain experimental observations, even when one or more of the theory's assumptions are violated.

In contrast, the experiments reported here not only obeyed the assumptions of the theory, but were also conducted in plasmas with peak densities well below quarter-critical density. This prevented other competing or coexisting phenomena from occurring, thereby providing clearly interpretable results. These are the first SBS experiments that were designed to be both a clear test of linear convective theory and pertinent to controlled fusion research.

A crucial part of this series of experiments was the development of a new instrument, the Multiple Angle Time Resolving Spectrometer (MATRS).

MATRS has the unique capability of both spectrally and temporally resolving absolute levels of scattered light at many angles simultaneously, and is the first of its kind used in laser-plasma experiments.

A detailed comparison of the theoretical predictions and the experimental observations is made. It is determined that there are significant qualitative and quantitative differences in the levels of scattered light, as well as the shapes of the spectra and the evolutions of the peak scattered power. Several possible mechanisms are considered to explain these discrepancies. It is believed that the theory fails to predict the observations because it does not incorporate many of the plasma details, as well as the early-time (non-steady-state) behavior of SBS.

Chapter 1

Our Insatiable Need for Energy

Most people have never considered the role that energy plays in their lives. This is ironic since plentiful energy resources have literally revolutionized the way modern society has evolved. Almost every technological innovation that man has developed depends on some source of energy. For example, without concentrated energy sources, our modes of transportation would be limited to animal-drawn carriages or ocean-worthy sailing ships that were common in the previous century. Even these modes depend on naturally occurring energy sources (animal feed, wind, etc.). Today our homes are filled with appliances that rely almost exclusively on electricity or the burning of natural gas and oil. Our dependency on plentiful, inexpensive energy has rapidly become an essential and irreplaceable part of our lives.

What is true about most of our energy resources to date is that they readily exist here on earth. Man's first productive use of energy occurred shortly after the discovery of fire. Man quickly learned that "burning" wood

created heat, and its use as a heat source for warmth and cooking for example, were immediately realized. The burning of wood was probably the first form of chemical energy that was liberated⁽¹⁾ for constructive use by man. Some would argue that mankind hasn't really progressed all that much, since even today the majority of the world's energy resources are derived from burning (converting chemical potential) of our natural resources⁽²⁾. Unfortunately, there are two basic problems with this form of energy production. First, the by-products that remain after burning wood and oil are environmentally toxic; an over-abundance of these products is widely thought to be unhealthy to our fragile eco-system. Second, and maybe even more crucial, these resources are far from limitless. Ours is the first generation that has had to face these two real threats from a dependency upon chemical energy. Our essential dependency on energy requires a search for new resources which are both cleaner and either renewable or at least extremely abundant.

Before continuing, it is important to address the overwhelming contribution of electricity. It is correct to note that electricity is in its own right, a very clean energy source. And with the possible exception of transportation, it is the dominant form of energy consumption in our lives. However, one must realize that the vast majority of electricity produced in the world today results from conversion of chemical energy, primarily through the burning of coal. Thus, although electricity itself doesn't threaten

(1)The term "liberation" is the most appropriate here because energy is not created; the energy stored in the chemical bonds formed in the molecules in wood or any other burnable substance are simply converted or "liberated" to a more readily usable form.

(2)Though we still burn wood, most chemical energy today is liberated by burning petroleum, natural gas and coal.

the eco-system, its methods of production often do, and it certainly isn't limitless.

It wasn't until the beginning of the 20th century that serious scientific research was conducted to seek new, alternative sources of electrical energy. Among those which have survived decades of scientific scrutiny are solar energy (collection and conversion of the light emitted by our Sun), hydroelectric energy (tapping the massive flow of the world's waterways to drive the turbines of electric generators), wind energy (harnessing the flow of wind to drive air turbines (windmills) to generate electricity) and geothermal energy (releasing the tremendous amount of heat that exists in the earth's core and converting it to electricity). All of these energy sources possess the virtues of being environmentally "clean" and virtually limitless. And, they have enjoyed varying degrees of commercial success. However, to this day none of these has come even close to providing us with enough energy to abandon or replace what little traditional resources remain. This is primarily due to the efficiency with which we have been able to convert these sources of energy into usable forms. In fact, research into new resources has had to focus almost exclusively on their efficient conversion. In time, it is possible that many of these together will be able to provide mankind with a small but significant fraction of his energy needs.

In 1945, man's ability to produce (convert) energy changed in a manner as dramatic as the original discovery of fire millennia earlier. The first artificially produced, large-scale nuclear reaction on earth literally shook the world when the first atomic device was exploded in the New Mexico desert in 1945. The physical process was nuclear fission—the combining of two heavy nuclei to produce a single nucleus whose mass was less than the sum of the original constituents. This reduction in total mass was released as energy in

the form of fast, energetic particles (neutrons) as is predicted by Einstein's famous mass-energy equivalence expression⁽³⁾, $E = mc^2$. It is a sad but recurring theme of humanity that the developments that led to this discovery were driven by the search for a weapon which could destroy people and structures. However, a detailed discussion of this concept would be merited in a dissertation in sociology, not physics. That is not to say that scientists are not at least partially responsible for the implementation of their discoveries — on the contrary, they are morally obligated to strive to influence how their acquired knowledge is used (or abused, depending upon one's personal point of view).

Though Szilard first formulated the concept of a neutron chain reaction in 1934 and fission was first experimentally confirmed by Dunning and others at Columbia University in 1939¹, it wasn't until the development of the fission bomb several years later that the useful potential for nuclear energy was widely realized, if for no other reason than nuclear processes liberate tremendous quantities of energy per unit reaction—they are almost incomprehensibly efficient—which is why it makes for a tremendously powerful weapon. Today in many countries of the world nuclear fission is the primary source of usable energy. However, as attractive as fission energy may seem, it is extremely important to note three major disadvantages. First, the fuel of fission, though some of which is naturally occurring on earth, is extremely rare. The technology to enrich the heavy isotopes of uranium necessary for today's fission reactors does exist, but because it is expensive and difficult, it cannot really be considered a limitless supply.

⁽³⁾It is historically interesting to note that, although these first experiments resulted in what later became the atomic bomb or atom bomb, the reaction by which the tremendous amount of energy is released is not atomic, but nuclear. The lay person's nomenclature has, over the years, been refined to account for this misnomer. They are now more commonly and more correctly referred to as nuclear weapons.

Second, the by-products of fission are both abundant and radioactive, which if uncontrolled makes them a serious threat to our eco-system, not to mention human health directly. Third, and probably perceived as the worst disadvantage it is undeniable that "run-away" or uncontrolled fission reactions can prove disastrous, as was evidenced in Chernobyl in 1986. In chemical reactions, adding more "fuel to the fire" will also result in the conversion of more energy. However, because of the ways in which we convert this energy, it is relatively trivial to control the reaction rate, if by no other means limiting the amount of fuel available. In fission reactors, on the other hand, it is far from trivial. Much technical effort has been devoted to the development of safe reactors, and the relative success or failure of this endeavor is now dominated by controversy and politics rather than technology.

Fusion is another nuclear process which liberates vast amounts of energy from very little fuel. In principle, fusion is far more attractive than fission because it does not produce the large quantities of toxic or radioactive waste that nuclear fission does⁽⁴⁾. The fuel proposed for fusion reactors, isotopes of hydrogen, are relatively abundant in the Earth's oceans (deuterium)²⁻⁷ or can be relatively easily produced from other elements⁸⁻¹⁵ (for example, tritium produced from lithium). Based on current trends in world population growth, it has been predicted that the Sun will self-extinguish before the Earth runs out of deuterium and tritium¹⁶. Thus, as

⁽⁴⁾Though the byproducts of the most probable laboratory fusion reaction (helium ions, neutrons) is neither toxic nor radioactive, the neutrons emitted during the reaction are captured preferentially by high-atomic-number (high-Z) materials that are in proximity to the reaction. For this reason, low-Z materials will most likely be used in the construction of containment vessels and other peripheral devices in future reactors. Even these materials will become radioactive from this process over time, and will need to be replaced. However, this very small amount of material is minuscule in comparison to the vast quantities of radioactive byproducts produced in nuclear fission reactors.

far as mankind is concerned, fusion is a potentially limitless source of energy. These two attributes alone have been the driving force behind the search for a cost-effective means for producing usable fusion energy.

Despite its great attraction, the efficient and safe production of fusion energy in the laboratory has eluded scientists since its inception. Numerous phenomena have complicated this seemingly simple goal. The work presented in this dissertation is the culmination of an effort to better understand one such phenomenon, namely parametric excitation of plasma density perturbations and the resulting loss of energy absorption through the scattering of laser light out of the plasma. This contribution attempts to contribute to the ultimate goal— to attain efficient, inexhaustible and safe production of fusion energy.

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Chapter 2

Plasma, Fusion, and Stimulated Brillouin Scattering

The research presented in this thesis is highly technical, and depends on at least a rudimentary understanding of several fundamental physics concepts. Since fusion is a thermonuclear process which requires tremendous energy densities to occur, the fuel is rapidly converted to plasma before the reaction begins. Consequently, the concept of a plasma is first discussed. Then, a brief discussion of thermonuclear fusion follows. Some of the plasma physics concepts necessary to understand the research discussed herein are then presented to the reader. Finally, a detailed technical discussion and derivation (the details of which appear in appendix A) of the coupled wave equations which form the basis of the theory of stimulated Brillouin scattering are given.

2.1 Plasma

Before one can understand the physics of fusion, one must understand the concept of plasma. When a gas is heated to sufficiently high temperatures, one or more of the electrons which were originally bound to the nuclei become free. This gas is now composed of a collection of negatively-charged electrons and positively-charged ions. Such a collection of electrons and ions is known as a plasma. Plasma and gases differ fundamentally in their behavior, as will be discussed further below. Plasma, the fourth state of matter (the other more familiar three being solid, liquid and gas) comprises the vast majority of matter in the universe, despite the fact that it rarely occurs naturally on earth (the earth's ionosphere and the plasma produced during a lightening strike are two exceptions). Because plasma is composed of charged particles, plasma can display collective behavior—it can behave as a fluid. However, since it is composed of particles, it can also behave by the laws of particle kinetics. It is this complexity that makes plasma such a rich topic of fundamental interest¹. However, scientific interest in plasmas has vastly increased since laboratory fusion was first hypothesized, and is one of the driving forces behind the research discussed in this dissertation.

Many fundamental plasma phenomena are important to fusion. Possibly the simplest phenomenon is that of plasma oscillations. The plasma frequency, ω_{pe} , is the natural oscillation frequency of a plasma, where the restoring force, $\vec{F} = -e\vec{E}$, originates from the electric fields generated by the spatial separation of positively-charged ions and negatively-charged electrons. Here, e and $\vec{E}$ are the charge of the electron and the electric field, respectively. The plasma frequency is given by

$$\omega_{pe} = \sqrt{\frac{4\pi n_e e^2}{m_e}} \quad (2.1)$$

where n_e and m_e are the electron density and mass, respectively. The concept of plasma frequency is critical to laser-plasma interactions. Light of frequency equal to the plasma frequency is specularly reflected, since the electrons can overcome their inertia and respond rapidly to shield out the electric fields of the electromagnetic radiation. Conversely, light of frequency greater than the plasma frequency is transmitted since the electrons cannot respond rapidly enough. The density in the plasma where the frequency of the light wave equals the plasma frequency is called the critical density, and similarly the location of the critical density is referred to as the critical surface. Light cannot penetrate beyond the critical surface in a plasma with an increasing density profile, as is often encountered in plasma produced from a solid target.

Another fundamental concept of interest is the Debye length. As inner-shell electrons of an atom can partially shield the positive charge of the nucleus and effectively reduce the electric field experienced by the outer-shell electrons, so too can plasma shield out electric fields. Since the electrons in a plasma move under the forces of attraction from the ions, they tend to shield out any net charge. The Debye length, λ_D , and the related concept of the number of particles in a Debye sphere, N_D (which is equal to the number of particles in a sphere of radius λ_D), are given by

$$\lambda_D = \sqrt{\frac{k_B T_e}{4\pi n_e e^2}} \quad (2.2)$$

and

$$N_D = \frac{4\pi n_e}{3} \lambda_D^3, \quad (2.3)$$

where k_B is the Boltzmann constant ($k_B = 8.63 \times 10^{-5} \text{ eV} \cdot ^\circ\text{K}^{-1}$, where $^\circ\text{K}$ is degrees Kelvin). The concept of Debye length is crucial to an understanding of the fundamental difference between a plasma and an ionized gas; the electric field a single Debye length away from an internally or externally imposed collection of charge in a plasma is reduced by $1/e$ due to the shielding by particles of opposite charge. Over many Debye lengths, the plasma appears to the observer to be neutral. Hence, unlike a simple ionized gas, a plasma is said to maintain quasi-neutrality—the net charge of a plasma on a macroscopic scale is neutral, despite the fact that it is composed of discrete charged particles. Similarly, the Debye sphere is a measure of the extent to which a plasma behaves collectively. Many particles in a Debye sphere ensures collective or fluid-like behavior, whereas only a few lead to longer-range interactions similar to those which dominate the behavior of an ideal gas.

Finally, there is one other fundamental plasma physics concept—Landau damping—which is crucial to the research reported in this dissertation. Landau damping is a non-collisional damping of waves in plasmas. A complete understanding of Landau damping requires a knowledge of kinetic theory, which is beyond the scope of this chapter. However, a qualitative understanding can be obtained from simple considerations.

Since plasma particles possess charge, they can respond to the electric fields of the waves, if the particles are resonant with the wave. In the limit of

a zero temperature plasma, all particles possess no random motion. Of course, no such limit can exist. However, in a plasma of finite temperature, the plasma particles possess a distribution of random velocities. A measurement of temperature is quantitatively related to the mean energy of the plasma particles. Strictly speaking, temperature is meaningless unless the distribution of particle speeds is Maxwellian. Nonetheless, the concept of temperature is often applied to non-Maxwellian plasmas in an attempt to quantify the kinetic energy contained within the random motions of the particles.

If there is a sufficient quantity of particles whose velocity approaches the phase velocity of a wave in the plasma, then those particles experience a nearly DC electric field from the wave, and they can be accelerated or decelerated, depending upon whether the particle speeds are slightly faster or slightly slower than the phase velocity of the wave. In the case of a Maxwellian distribution, the number of particles decreases monotonically with increasing velocity. If the phase velocity of the wave is on the order of a few times the mean velocity of the particles in a Maxwellian plasma, then there are many more particles which travel slightly slower than the wave than there are which travel slightly faster. Since more particles will be accelerated than decelerated in this situation, the result is that a net amount of energy is transferred from the wave to the slower particles, thus modifying the distribution and damping the wave.

Quantitatively speaking, Landau damping scales with the quantity $\exp(-v_{ph}^2/v_{th}^2)$, where v_{ph} is the phase velocity of the wave and v_{th} is the thermal velocity of the particles. From this it is clear that if the phase velocity of the wave is much greater than the thermal velocity of the particles

($v_{ph} \gg v_{th}$), Landau damping is negligible. On the other hand if they are comparable ($v_{ph} \approx v_{th}$), then clearly the wave can be significantly damped.

A derivation of these and other pertinent, fundamental plasma quantities can be found in Chen¹.

2.2 Thermonuclear Fusion

Whereas nuclear fission is the separation of a single heavy nucleus into two lighter nuclei, fusion is the process of combining two light nuclei into one heavier nucleus. Fusion can occur between any two elements that are lighter than the naturally dominant isotope of elemental iron ($Z=26$, $A=55.85$)⁽¹⁾. Thermonuclear fusion is the process which powers our own Sun. Like fission products, the total mass of the fusion products is less than the original reactants. The difference in mass is converted to kinetic energy in the products. This mass-energy conversion or "equivalence" is governed by Einstein's famous relation.

$$E = mc^2 \quad (2.4)$$

where E , m and c are energy, rest mass and the speed of light, respectively. It is readily apparent from this equivalence why fusion is such a promising form of energy; the energy liberated by converting matter to energy is the mass multiplied by the speed of light *squared*. Thus, a very small mass can be converted to a tremendous amount of energy in a nuclear reaction.

⁽¹⁾Strictly speaking, the peak of the nuclear binding energy curve, which dictates the above-mentioned behavior, occurs for any element with 56 *nucleons*. In principle, any element lighter than iron can possess isotopes with 56 nucleons, due to the extra neutrons associated with these isotopes. Since iron's dominant isotope possess approximately this number, it is typically quoted as the element occurring at the peak of the nuclear binding energy curve.

There are many different nuclear reactions which can occur. However, most require fuel temperatures which are not readily obtainable in a laboratory setting. A few, however, occur at temperatures that are thought to be obtainable within the laboratory. They include

$$D + T \Rightarrow {}^4\text{He}(3.5\text{MeV}) + n(14.1\text{MeV}) \equiv (17.6\text{MeV}) \quad (2.5)$$

$$D + {}^3\text{He} \Rightarrow {}^4\text{He}(3.6\text{MeV}) + n(14.7\text{MeV}) \equiv (18.3\text{MeV}) \quad (2.6)$$

$$D + D \Rightarrow {}^3\text{He}(0.82\text{MeV}) + n(2.45\text{MeV}) \equiv (3.27\text{MeV}) \quad (2.7)$$

$$D + D \Rightarrow T(1.01\text{MeV}) + p(3.02\text{MeV}) \equiv (4.03\text{MeV}) \quad (2.8)$$

where D and T are deuterium and tritium, the first and second isotopes of Hydrogen, respectively, n and p are the neutron and proton, respectively, and the kinetic energy of each product is given in parentheses. The reaction given in (2.5) is currently considered the most promising path to a reasonable fusion reactor since its reaction cross section peaks at a much lower temperature than the others. However, it requires tritium, a radioactive and rare isotope of hydrogen, which the others do not. Expression (2.6) produces the same products at roughly the same energies as expression (2.5), and does not require tritium, but its reaction cross section is many orders of magnitude lower at similar temperatures. To achieve a similar reaction cross section requires much higher plasma temperatures which are currently difficult to obtain in the laboratory. Expressions (2.7) and (2.8) require similar plasma temperatures to obtain a reasonable reaction cross section. However, expression (2.8) is unique in this set and consequently attractive since it

produces a proton rather than a neutron. Unfortunately, one cannot drive reaction (2.8) without driving reaction (2.7) and thus avoid production of neutrons. Despite this, and although (2.7) is three times as likely to occur as (2.8), the production of protons affords the possibility of electromagnetic conversion of the kinetic energy of the proton into electrical energy, thus bypassing the inefficient heat cycle necessary to extract energy from the other reactions' products. Finally, it is important to note that there are many other reactions that can occur between other elements, which are not listed here because they are currently even less achievable in the laboratory.

2.3 The Lawson Criterion

In order for a fusion reaction to be sustained, not only must the fusing nuclei be sufficiently energetic (sufficiently high temperature), but they must be in close enough proximity (sufficiently high density) for a sufficient duration. These latter two requirements can be expressed mathematically by the Lawson criterion²

$$\eta\tau \geq 10^{14} (\text{sec} \cdot \text{cm}^{-3}) \quad (2.9)$$

where η is the density of fusible ions and τ is the energy confinement time. It is clear from this expression alone that to obtain a fusion reaction, one must either contain ions at moderate density for a long duration or contain ions at very high density for a shorter duration.

2.4 Magnetic vs. Inertial Confinement of Hot Plasmas

Our Sun confines its particles through gravitation. The mass of the Sun is sufficiently large that the fuel for fusion cannot escape its own gravitational attraction. In the laboratory we are far less fortunate. The gravitational attraction of the Earth is simply too small⁽²⁾, and the exceedingly high temperatures required to even sustain a plasma in a laboratory make mechanical containment virtually impossible. Schemes have been and are continuing to be devised to contain the extremely hot plasma for a sufficient duration so that the fusion reactions can take place in the laboratory. Two such schemes that are rigorously studied are magnetic containment and inertial confinement fusion.

Magnetic containment relies on the fundamental laws of electromagnetism to contain the charged particles in the plasma. Since the particles in a plasma are charged, their motion can be modified or controlled by the external application of magnetic fields. Magnetic containment attempts to fulfill the Lawson criterion by attempting to contain moderate plasma densities ($\sim 10^{14} \text{ cm}^{-3}$) for long periods of time ($\sim 1 \text{ sec}$). This is a very complex problem since plasmas are diamagnetic; they tend to generate magnetic fields which oppose the externally applied fields. This can eventually result in an unstable or collapsed containment "vessel". In an effort to perfect these techniques, various devices have been postulated and built over the past 30 years in an attempt to attain sufficient containment of thermonuclear plasma, including Tokamaks³, stellarators⁴, and mirror machines³. Briefly, a Tokamak is a torus-shaped vacuum chamber with externally-applied toroidal

⁽²⁾This of course is *highly fortunate* since a gravitational field sufficient to sustain fusion reactions on Earth would render the planet uninhabitable and consequently the search for energy sources would be moot.

and vertical magnetic fields which both contain and stabilize the plasma. It has, to date, been the most successful magnetic containment device built. Stellerators are similar to Tokamaks, but have externally applied poloidal magnetic fields to improve stability. Mirror machines are cylindrical vessels with magnetic "mirrors" that reflect particles at each end. Though the perfect magnetic containment system has not yet been developed, much progress has been made in perfecting this concept in search of a reliable fusion reactor.

Inertial confinement fusion (ICF) is based on a fundamentally different principle, and addresses the Lawson criterion from a completely different perspective than magnetic containment. In ICF⁵, very high fuel densities are obtained ($>10^{25} \text{ cm}^{-3}$) for extremely short periods of time ($<10^{-10} \text{ sec}$) by spherically imploding solid or gas-filled microballoon targets with lasers, particle beams, or intense x-rays. ICF targets consist of hollow spherical shells which contain an outer "ablator" layer and an inner layer of D-T fuel. Additionally, the shell is filled with a low density D-T gas. In direct-drive ICF, the outer shell material of the capsule is ablated by direct incidence of the driver beams upon the fuel pellet, and through conservation of momentum or the "rocket-effect", the remaining D-T shell is compressed radially inward, which also serves to compress the low density D-T gas. As stagnation occurs and the inner D-T gas is heated and thermonuclear ignition occurs. This burn front then propagates radially outward through the dense D-T shell and high gain is achieved⁵⁻⁷. Figure 2.1 illustrates the process of direct-drive ICF. In indirect-drive ICF, the particle or laser beams are incident upon a high-atomic-number (Z) host material, which is converted to a very hot plasma through collisional absorption of the driver energy. This hot, high-Z plasma re-radiates x-rays, and these x-rays "drive" (compress and

heat) the fuel capsule in a manner similar to direct drive ICF. Figure 2.2 illustrates the process of indirect-drive laser fusion.

2.5 The NOVA Laser Facility

There are several ICF research facilities which are currently in operation around the world. The largest facility is the NOVA laser facility at Lawrence Livermore National Laboratory⁸. NOVA is currently the world's most powerful laser, capable of producing in excess of 100 terawatts of power at the fundamental laser wavelength of 1.054 μm . NOVA is a neodymium-doped glass (Nd:glass) laser system comprised of ten arms seeded by a single oscillator. Each arm consists of 16 single-pass Nd:glass rod or disk amplifiers, which are radiatively pumped with Xenon flashlamps to obtain the necessary population inversion in the Nd:glass. There are 10,000 capacitors, capable of storing 60 megajoules of electrical energy and delivering 10 megaamperes of current and 120 gigawatts of power to the flashlamps⁹. The total gain of each amplifier chain is $\sim 10^7$. The optical path length of each arm from oscillator to target is approximately 280 meters and the final beam diameter before focusing on target is approximately 74 cm. Despite this large beam diameter, the final $f/4.3$ focusing optics can produce a $\sim 150 \mu\text{m}$ diameter spot on a target containing 80% of the delivered energy.

NOVA has two separate target "bays" or laboratories, each containing a spherical vacuum chamber. The larger chamber is approximately 172 inches in inside diameter and can receive light from all ten arms, whereas the smaller chamber, which is approximately 60 inches in inside diameter, can receive light from only two arms. There is a "switchyard" which handles routing of the beams, and it is possible to route two beams to the smaller

chamber and the remaining eight to the larger chamber to perform separate but simultaneous experiments in both target chambers. In the ten-beam target chamber, all beams are distributed every 72° along two opposing cones of 50° half-angle whose axes are horizontal and counter-incident (figure 2.3). In the two-beam chamber, the two beams are horizontally opposed and counter-propagating. Each chamber possesses a large battery of instruments for diagnosing the laser beams, the plasma that is produced, and the vacuum and other support systems. Target alignment, laser alignment, and routine target and laser diagnostics are maintained and controlled by highly skilled and specially trained technicians, and experimental diagnostics are controlled by physicists. All experiments discussed in this dissertation were performed in the ten-beam target chamber.

Each of the Nova beams passes through a pair of potassium dihydrogen phosphate (KDP) crystal arrays, giving each beam the independent capability of delivering light on target of wavelength $1.054\mu\text{m}$ (1ω , fundamental frequency), $0.527\mu\text{m}$ (2ω , second harmonic), or $0.351\mu\text{m}$ (3ω , third harmonic). The measured second-harmonic (2ω) conversion efficiency, which was the frequency used for all interaction experiments reported in this dissertation, is shown in figure 2.4. Third-harmonic (3ω) conversion efficiency is approximately 25% more efficient above $1\text{GW}/\text{cm}^2$. Energy calibration of each of the NOVA beams is conducted during periodic, dedicated calibration shots using separate whole-beam glass calorimeters that are located on the inside chamber wall opposite each beam entrance port. Fundamental and third harmonic energies are calibrated to within $\pm 10\%$ and second harmonic energies are calibrated to within $\pm 25\%$ ¹⁰. Prior to an experiment, the oscillator and rod amplifiers are discharged to ensure that the energies produced are within the accuracy requested by the

experimenters. During an experiment, energy at the fundamental frequency is measured for each beam, and the harmonic energy delivered to target is inferred from this measurement and the data collected from dedicated calibration experiments. Because each beam has a separate amplifier chain, each may have a different output energy and may be independently advanced or retarded in time with respect to other beams. This, coupled with the frequency conversion capability, gives the experimenter the ability to “preform” the plasma with some of the beams at one or more frequencies, energies, and spot diameters, then conduct laser-plasma interaction experiments with a separate beam of different energy, frequency, spot diameter, and arrival time. The experiments reported in this dissertation made extensive and unprecedented use of these capabilities.

The Nova oscillator has significant temporal pulse-shaping capabilities, which have proven very important in the study of ICF. For example, this capability allows the experimenter to preform a plasma with a lower intensity “foot” or prepulse and then heat the plasma with a high intensity main pulse. Although all beams are seeded by one oscillator, the final fundamental (1ω) pulse shape for each beam may differ since the gain of each amplifier is limited by how intensely it is optically pumped. In other words, if the amplifiers from different beam lines are operated to produce different amounts of amplification, different pulse shapes at the end of the amplifier chain will result. Further, since the conversion efficiencies of the KDP crystals are both output-frequency and input-intensity dependent, it is quite possible to produce two or more beams with different frequencies and pulse shapes during a single experiment, from the single oscillator output.

When frequency conversion to a particular harmonic is performed, relatively small fractions of the energy incident upon the KDP crystals are

transmitted as unconverted (1ω) or converted (alternate-harmonic) light. Because the focusing optics are refractive, their focal lengths are frequency-dependent. Thus, the residual fundamental and alternate-harmonic light which is transmitted to the target chamber will have different focal positions. In order to prevent light of unwanted frequency from reaching the target, there is a beam block or hard stop in the center of the beam path which results in an annular beam delivered to the target chamber. The resulting hole in the center of the beam beyond the focal plane, coupled with the frequency-dependent focal length, ensures unwanted light from reaching the target. This effect is demonstrated for example in figure 2.5, where green light ($2\omega, 0.527\mu m$) is focused on target, and unwanted frequencies ($1\omega, 1.054\mu m$ and $3\omega, 0.351\mu m$) are diverted away from the target.

2.6 Parametric Instabilities

Many plasma effects can reduce the coupling efficiency of the laser to the plasma fuel in laser fusion. For obvious reasons, it is essential to optimize this coupling. Plasmas can support parametric instabilities which can contribute to this inefficiency.

The concept of parametric instability is not unique to laser-plasma interactions. One of the interesting aspects of a parametric instability is that all of the existing resonances that are driven unstable must be present for any of them to be driven unstable. As an intuitive example of this, consider a mechanical system of two different masses connected by a spring (figure 2.6). If mass #2 is fixed in space and mass #1 is allowed to move freely in one dimension, the system has one resonance frequency given by $\omega_1 = \sqrt{k/m_1}$ where k is the spring constant and m_1 is the mass of object #1 (figure 2.6.a).

Reverse the situation and one similarly finds $\omega_2 = \sqrt{k/m_2}$. Since these are resonant frequencies, an initial impulse is all that is required to sustain these oscillations, as long as there is no damping present in the system. Now, suppose that there is an energy source within the system which can independently drive one of these disturbances (figure 2.6.b). In our mechanical analog, this could be simply a motor and cog connected to m_1 , pushing and pulling it at the resonance frequency ω_1 while m_2 is fixed in space. Under the absence of damping, the driven resonance can grow indefinitely and become unstable.

Now, suppose that both masses are free to move in one dimension (figure 2.6.c). In this case, the resonance frequency of the system has been changed to $\sqrt{\omega_1^2 + \omega_2^2}$. Unlike before, the system is no longer resonant at the frequency ω_1 , and if driven at that frequency by the pump, the system will no longer become unstable. If the pump now drives the system at the frequency $\omega_0 = \sqrt{\omega_1^2 + \omega_2^2}$ then the system once again has the potential to become unstable.

Though this illustrates the idea of resonance coupling quite well, it is subtly different from a parametric instability. This is because in a parametric instability, the two *original* resonances grow unstable at the expense of the pump. In the example above, a *new* resonant mode is formed and is driven unstable by the pump, while the original resonances do not grow unstable. For a more accurate, but less intuitive example analogous to parametric instability, see the example given in Chen¹.

In laser-plasma interactions, the coupling that is necessary to drive a parametric instability arises from Coulomb interaction and the associated electric fields. This can be easily understood by examining figure 2.7. Consider a plasma with an initial spatial disturbance in density, δn_e . It can

easily originate from thermal fluctuations in the plasma. Now, a light wave of frequency ω_0 (the laser) is incident upon this density disturbance. The charged particles in the density fluctuation experience transverse oscillations of velocity $\vec{v}_{osc} = e\vec{E}_0/m\omega_0$, under the influence of the electric field of the laser, $\vec{E}_0$. These oscillating charged particles, which constitute an oscillating current, re-radiate a scattered light wave with electric field, $\delta\vec{E}_s$, at frequency ω_s . The electric field from the pump beats with the scattered light wave resulting in a spatially varying electric field energy density, $\vec{E}_0 \cdot \vec{E}_s$. A spatially varying electric field pressure proportional to $-\vec{E}_0 \cdot \vec{E}_s$, then serves to further increase the density perturbation, δn_e , which in turn increases the amplitude of the scattered light wave, $\delta\vec{E}_s$, and the plasma and scattered light waves grow unstable.

There are two classes of parametric instabilities that can be observed in plasmas. The first are purely growing instabilities. In this class of instability, the plasma disturbance (one of the two daughter waves) is a purely growing disturbance with zero frequency and finite wave vector. Since this disturbance has zero frequency, it is stationary in the plasma. Examples of purely growing parametric instabilities in laser plasmas are the oscillating two stream¹¹ and filamentation¹²⁻²¹ instabilities. The second class of instability are so-called “decay” instabilities since the “parent” laser light wave decays into traveling “daughter” waves. Each daughter wave has a finite frequency and wave vector, as defined by the conditions for resonance, or “matching” conditions

$$\omega_0 = \sum_i \omega_i \quad (2.10)$$

and

$$\vec{k}_0 = \sum_i \vec{k}_i. \quad (2.11)$$

Here, the subscript 0 refers to the parent or pump wave and the subscript i refers to the i th daughter wave. These matching conditions, which are necessary conditions for a real solution to the coupled partial differential equations governing each process (and are derived later in this chapter, as well as appendix A), are equivalently expressions of conservation of total wave energy and total wave momentum, respectively. Examples of parametric decay instabilities in laser plasmas are stimulated Raman scattering (SRS)²²⁻³², stimulated Brillouin scattering (SBS)³³⁻⁴⁰, two plasmon decay (TPD)⁴¹⁻⁴⁹, and ion acoustic decay (IAD)⁵⁰⁻⁵³. The remainder of this dissertation will be devoted to the study of stimulated Brillouin scattering.

2.7 Stimulated Brillouin Scattering In Laser Plasmas

Stimulated Brillouin scattering is a parametric decay instability which occurs between a pump light wave (ω_0), an ion acoustic wave (ω_{ia}), and a scattered light wave (ω_s). [In the following formalism, for reasons of brevity and established tradition, the faster scattered light wave will be assigned the subscript 1 ($\omega_s = \omega_1$), the slower ion acoustic wave will be assigned the subscript 2 ($\omega_{ia} = \omega_2$), and the laser or pump wave will retain the subscript 0.] SBS is of particular concern to laser fusion since it can, in principle, scatter most of the laser light out of the plasma and consequently, drastically reduce coupling efficiency.

The fundamental theory of SBS in a uniform plasma can be derived from the two coupled wave equations describing the two driven daughter waves

$$\left(\frac{\partial^2}{\partial t^2} - c^2 \nabla^2 + \omega_{pe}^2 \right) \tilde{\tilde{A}} = - \frac{4\pi Z e^2}{m_e} \tilde{n}_i \tilde{A}_{laser} \quad (2.12)$$

$$\left(\frac{\partial^2}{\partial t^2} - c_s^2 \nabla^2 \right) \tilde{n}_i = \frac{n_{0i} Z e^2}{m_e M c^2} \nabla^2 (\tilde{A}_{laser} \cdot \tilde{\tilde{A}}) \quad (2.13)$$

where damping and plasma flow have been neglected for the purposes of illustration. Here, Z , e , m_e , M , n_{0i} , and $\tilde{n}_i$ are the atomic number, electronic charge, electron mass, ion mass, ion background density and ion fluctuation density, respectively. $\tilde{A}_{laser}$, $\tilde{\tilde{A}}$, c , and c_s are the background (laser) vector potential, the fluctuating (scattered light wave) vector potential, light speed, and sound speed, respectively. A detailed derivation of these two expressions can be found in appendix A. Note that, to arrive at expression (2.12) from expression (A.16) appearing in the appendix, the appropriate substitution $\tilde{n}_e \approx Z \tilde{n}_i$ has been made so that it is clear that the density fluctuations of interest occur in the ions.

The first expression describes a light fluctuation (scattered light wave), $\tilde{\tilde{A}}$, being driven by the coupling of the laser to an ion density fluctuation ($\tilde{n}_i \tilde{A}_{laser}$). The second expression describes an ion density fluctuation (ion acoustic wave), $\tilde{n}_i$, being driven by the coupling of the laser light to the light fluctuation ($\tilde{A}_{laser} \cdot \tilde{\tilde{A}}$). In other words, as described before, the pump and each daughter wave interacts to increase the amplitude of the other daughter wave, at the expense of the pump, resulting in unstable growth of both daughter waves.

A solution of these two coupled wave equations results in a dispersion relation which describes the process. By taking the laser to be a traveling wave with frequency ω_0 and wave vector $\vec{k}_0$, the vector potential then has the form

$$\vec{A}_{laser}(\vec{x}, t) = \vec{A}_{laser} e^{i(\vec{k}_0 \cdot \vec{x} - \omega_0 t)}. \quad (2.14)$$

where the real part of the vector potential is defined as

$$\text{Re}(\vec{A}_{laser}) \equiv \frac{1}{2}(\vec{A}_{laser} + \vec{A}_{laser}^*) \quad (2.15)$$

and (*) denotes complex conjugate. Substituting this into the two coupled equations and Fourier-analyzing the two equations are transformed into⁵⁴

$$D(k, \omega) \vec{\tilde{A}}(\vec{k}, \omega) = \frac{4\pi e^2}{m_e} \frac{\vec{A}_{laser}}{2} \left[\vec{n}_e(\vec{k} - \vec{k}_0, \omega - \omega_0) + \vec{n}_e(\vec{k} + \vec{k}_0, \omega + \omega_0) \right] \quad (2.16)$$

and

$$(\omega^2 - k^2 c_s^2) \vec{n}_e(\vec{k}, \omega) = \frac{Z n_0 e^2}{m M c^2} \frac{k^2 \vec{A}_{laser}}{2} \cdot \left[\vec{\tilde{A}}(\vec{k} - \vec{k}_0, \omega - \omega_0) + \vec{\tilde{A}}(\vec{k} + \vec{k}_0, \omega + \omega_0) \right] \quad (2.17)$$

where $D(k, \omega) = \omega^2 - \omega_{pe}^2 - k^2 c^2$ and where we have re-transformed $(n_{0i}, \vec{n}_i) \Rightarrow (n_{0e}, \vec{n}_e)$ to follow convention. By letting $(\vec{k}, \omega) \Rightarrow (\vec{k} \pm \vec{k}_0, \omega \pm \omega_0)$ successively into expressions (2.15) and (2.16), the two equations can be combined. Since we are interested in the low frequency ion waves, we take

$\omega \ll \omega_0$ in equation (2.16). We can neglect the terms evaluated at $(\vec{k} \pm 2\vec{k}_0, \omega \pm 2\omega_0)$ as being very off-resonant, resulting in the dispersion relation

$$\omega^2 - k^2 c_s^2 = \frac{1}{4} k^2 v_{os}^2 \omega_{pi}^2 \left[\frac{1}{D(\vec{k} - \vec{k}_0, \omega - \omega_0)} + \frac{1}{D(\vec{k} + \vec{k}_0, \omega + \omega_0)} \right] \quad (2.18)$$

where $\omega_{pi} = \sqrt{4\pi n_e Z e^2 / M}$ is the ion plasma frequency and M is the ion mass.

Growth rates and damping thresholds for SE in homogeneous plasmas can now be obtained from this dispersion relation. The most commonly considered case is direct backscatter, which is the direction of largest growth in a homogeneous plasma. Substituting $\omega = kc_s + i\gamma$ into the dispersion relation and retaining only the down-shifted (Stokes) component, we arrive at the ion acoustic wave number and growth rate

$$k = k_{ia} = 2k_0 - \frac{2\omega_0}{c} \frac{c_s}{c} \quad (2.19)$$

$$\gamma = \frac{1}{4} \frac{k_{ia} v_{os} \omega_{pi}}{\sqrt{k_{ia} \omega_s c_s}} \quad (2.20)$$

If we were to include terms in the wave equations which accounted for damping of the scattered light wave and ion density perturbations, we would arrive at a threshold condition for SBS to grow. This condition is given simply by $\gamma > \sqrt{\gamma_{ia} \gamma_s}$ where $\gamma_{ia}(\gamma_s)$ is the amplitude damping rate of the ion acoustic (scattered light) wave⁵⁴. Typically, the scattered light wave is

damped by collisions between the electrons and ions, whereas the ion acoustic waves are damped by Landau damping⁵⁴.

Damping of SBS can originate from other sources as well. The above derivations assumed that the plasma was homogeneous. However, in the laboratory, it is nearly impossible to produce a homogeneous plasma. Most plasmas have gradients in temperature, density, and flow velocity, and these can each potentially stabilize SBS. To understand this phenomenon, one must consider the dispersion relations for the light waves and ion acoustic waves in a plasma with finite density, temperature, and flow velocity. The dispersion relation for electromagnetic waves in an unmagnetized plasma can be readily determined from the undriven counterpart to the scattered light wave equation derived earlier, equation (2.12)

$$\omega_{0,s}^2 = \omega_{pe}^2 + c^2 k_{0,s}^2 \quad (2.21)$$

where the subscripts 0,s represent the laser and scattered light waves, respectively. Similarly, the dispersion relation for the ion acoustic wave can be arrived at from the undriven counterpart to equation (2.13) and is given by

$$\omega_{ia} = k_{ia} c_s \quad (2.22)$$

where

$$c_s = \sqrt{\frac{3T_i}{M_i} + \frac{ZT_e}{M_i(1 + k_{ia}^2 \lambda_{De}^2)}} \quad (2.23)$$

is the sound speed, and Z, T_e, T_i, M_i and λ_{De} are the ion charge state, electron temperature, ion temperature, ion mass and Debye length, respectively.

This dispersion relation, however, is incomplete. Since plasma flow often approaches or even exceeds the sound speed (the velocity at which ion acoustic waves travel in a plasma) significant Doppler shifting of the ion acoustic waves can occur. Thus, a more complete expression for the dispersion relation of an ion acoustic wave in an unmagnetized plasma is

$$\omega_{ia} = \bar{k}_{ia} \cdot \bar{u} + k_{ia} c_s. \quad (2.24)$$

where $\bar{u}$ is the plasma flow velocity.

These dispersion relations, along with the frequency and wave vector matching conditions, uniquely describe SBS for a given set of plasma conditions $(n_e, n_i, T_e, T_i, \bar{u})$ at a given angle of observation. Now, in an inhomogeneous plasma, the plasma properties and their gradients are functions of both position and time. Hence, for each position, time, and direction of scattering there is a unique resonant triad, as dictated by the matching conditions and plasma properties. But, because slightly off-resonant growth is possible, there is effectively a finite-sized “matching region” in the plasma for a given set of plasma and wave parameters. As any of the waves travels outside of this region, the plasma conditions change, and the frequency and wave vector matching conditions are no longer fulfilled. At this point, the process is no longer resonant. The growth rate of SBS must overcome this convective energy loss in order to prevail.

The convective nature of SBS was first addressed by Rosenbluth in 1972⁵⁵. To arrive at a mathematical description of convective gain of SBS, WKB theory is invoked to transform the second-order coupled wave equations to first-order coupled mode equations⁵⁶ (shown here in 1-D for simplicity)

$$\left[\frac{\partial}{\partial t} + v_{g1} \frac{\partial}{\partial x} + \Gamma_s \right] a_1 = \gamma_1 a_2^* e^{i\Phi} \quad (2.25)$$

and

$$\left[\frac{\partial}{\partial t} + v_{g2} \frac{\partial}{\partial x} + \Gamma_{ia} \right] a_2^* = \gamma_2 a_1 e^{-i\Phi} \quad (2.26)$$

where $a_j = \sqrt{k_j} A_j$ and v_{gj} are the complex action amplitude (and where the asterisk denotes complex conjugate), and group velocity of wave j , respectively. Here, $\Phi \equiv \phi_0 - \phi_1 - \phi_2$ is the phase mismatch of the three waves of the triad: at resonance $\Phi = 0$. Also, Γ_s (Γ_{ia}) is the amplitude damping rate of the scattered light (ion acoustic) wave, and $\gamma_{1,2}$ are the growth rates of the fast and slow daughter waves. A detailed and complete transformation from the wave equations to the mode equations is given by Batha⁵⁷.

One of the primary goals of this research was to test the validity of the convective theory of SBS in a series of well designed laser plasma experiments. Though the details of the convective theory which evolve from expressions (2.25) and (2.26) is crucial to this goal, they are beyond the scope of this chapter, and will be discussed in a later chapter.

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Figure 2.1 Illustration of a direct-drive fusion pellet. Lasers incident upon the outer layer of the pellet ablate the outer surface. By conservation of momentum, the pusher layer compresses the D-T gas at the center. Heating of the gas occurs from compression and stagnation of the pusher. Ignition occurs in the gas and burn then propagates radially outward through the pusher, resulting in high gain.

Figure 2.2 Illustration of an indirect-drive target. Lasers enter the radiation case through the entrances holes and form a plasma on the inside of the radiation case. This plasma efficiently converts the laser energy into x-rays, which then drive the fuel pellet.

Figure 2.3 Orientation of all ten NOVA beams. The light blue beams are the East beams (2, 4, 6, 8 and 10) and the dark blue beams are the West beams (1, 3, 5, 7 and 9). Beams are distributed along two opposing cones of 50° half angle and the beams on each cone are equally spaced at 72°. Even and odd beams are not opposing or interwoven—that is, there is a plane of symmetry between the even and odd beams.

Figure 2.4 Conversion efficiency of second harmonic (2ω) on NOVA. Third harmonic (3ω) conversion efficiency is approximately 25% more efficient over this range of input fluences. The diameter of the KDP crystal array is approximately 64 cm.

Figure 2.5 Annular focusing of different harmonics simultaneously on NOVA. Focal Length of final focusing lenses are wavelength dependent. Lens positions are adjusted so that the desired wavelength is focused at the appropriate location with respect to the target. Residual light of unwanted wavelength misses the target due to the annulus-shaped beam and different location of best-focus.

Figure 2.6 Mechanical analog illustrating a feature of a parametric instability. In (a) mass #1 is free and mass #2 is fixed. In (b) mass #1 is driven at resonance. In (c) mass #2 is now free to move as well, and the system is driven at the new resonance, $\omega_0 = \sqrt{\omega_1^2 + \omega_2^2}$ which would not be at resonance without both masses free to move.

Figure 2.7 Illustration of a laser-plasma instability. Starting from any point on the circle, this process can continue to grow unstable indefinitely until some non-linear process quenches it. The quantities at the top and bottom are the daughter waves and they are driven by the forces that appear at the left and right. The daughter wave energies grow at the expense of the energy in the laser.

Chapter 3

A Review of Past Experiments in Stimulated Brillouin Scattering

Because of its impact on inertial confinement fusion (ICF) as well as its intrinsic interest, stimulated Brillouin scattering (SBS) has been studied in laboratory plasmas for many years. There have been many experiments conducted in a variety of plasma environments which have attempted to understand the scaling of SBS with various plasma and driver parameters, as well as to detect and understand levels of instability threshold and saturation. The comprehensive discussion of all experiments is far too long to present here in detail. However, it is both instructive and insightful to review a few which had a significant impact on both the understanding of SBS and the direction that future experiments have taken.

3.1 Microwave Experiments

Though the vast majority of productive ICF experiments have taken place using lasers, there have been other testbeds for understanding SBS. Microwave experiments are the most notable example. Though several laser systems capable of reaching sufficient energy densities existed throughout the scientific community by the time the first microwave experiments were conducted, they were extremely large and expensive to build, operate and maintain. For the university laboratory, microwave sources were a relatively inexpensive and readily available alternative to lasers. Further, since the wavelength of the microwaves is much longer than the laser's, more traditional methods of plasma production could be employed to produce plasmas near the critical-density for microwave experiments.

The first microwave experiments that observed SBS in an underdense plasma were conducted by Huey *et al.*¹. In these experiments, plasma was preformed from various gases (hydrogen, helium, neon, and argon) using a direct current filament discharge. SBS was driven using 3-16 GHz microwaves, and the peak plasma density was approximately 10% of the critical density for the pump microwaves. The scattered microwave wavelength shifts were not measured, but peak backscatter reflectivity proved to be less than 5%. In another experiment conducted by Pawley *et al.*², the plasma was formed by the same microwaves that were used to drive SBS. The host gas in these experiments was limited to helium. The maximum saturated amplitude of ion density fluctuations ($\delta n/n$) was measured with collective Thomson scattering to be less than 6%. The dominant saturation mechanism at densities below $0.2 n_c$ was determined to be ion tail heating, where as at densities from 0.3 - $0.5 n_c$ the dominant

mechanism was believed to be pump depletion. In a third series of experiments conducted by DeGroot *et al.*³, intensity-scaling of SBS was studied in an argon plasma produced by an energetic electron beam. To ensure that a sufficient ion noise source was present, standing microwaves were introduced into the plasma which not only enhanced ion fluctuations, but also served as the pump wave for SBS. The instability was driven in plasma with a peak density less than half-critical, and with electron and ion temperatures less than or equal to $\sim 2\text{eV}$. In these experiments, the plasma was significantly perturbed by the microwave pump, with the ratio of electron quiver velocity to electron thermal velocity approaching unity ($v_o/v_e \cong 1$). Further, the level of ion fluctuations was indirectly measured to exceed $\delta n/n \geq 1$. Eventually, nonlinear saturation occurs, resulting in ion wave collapse and strong ion heating.

Though microwave experiments are useful for the indirect study of laser-plasma interactions, the results obtained must be put into context. First, microwaves would make a particularly poor choice as an ICF driver since the densities necessary to produce fusion are many orders of magnitude above critical density for microwaves. This would result in most of the energy being deposited in the thin outer surface layer of the target plasma. Second, because of microwaves' extremely long wavelengths compared to optical light, it is very difficult to devise experiments that make the quantity L/λ_0 very large, where L is the length of the interaction region in the plasma and λ_0 is the vacuum wavelength of the pump. In contrast, most ICF laser-plasma experiments tend to produce $L/\lambda_0 \gg 100$. As such, in microwave experiments, the instability can grow only over a few wavelengths at most. This is a particularly important difference if the instability grows as it travels through the plasma (convective growth).

3.2 Experiments involving CO₂ laser-driven SBS

One of the first laser systems developed that were both powerful enough and reliable enough to be considered as a viable candidate as a laser fusion driver was the CO₂ gas laser. Because this laser relies on a gas as the lasing medium, many of the energy density and damage issues associated with solid state lasers were obviously avoided. Because of their relative availability and their direct relevance, there were many experiments conducted to measure SBS driven by CO₂ lasers, especially in the early years of laser fusion research.

A large class of experiments were conducted using preformed plasmas—that is, the plasmas in which the SBS was driven were produced prior to the experiment by other means. Two particularly popular methods of plasma pre-production were z-pinch devices⁴⁻⁷ and pulsed arc- or diode-discharge devices⁸⁻¹¹. Both of these methods used gases as a plasma host material.

Turechek *et al.*¹¹ conducted a series of experiments in helium and argon underdense ($\leq 0.25 n_c$) plasmas that were preformed by diode discharge. The intensity of the laser driver was approximately 5×10^{10} W/cm². The presence of SBS was confirmed by measuring the frequency shift of the scattered light, and by a direct measurement of the low frequency of the ion acoustic waves via the time-dependent Stark effect, in which a forbidden transition in He-I is induced by the electric field of the ion acoustic wave. Levels of direct backscatter were not reported.

In a series of experiments conducted by Herbst *et al.*^{8, 9} backscatter reflectivities were found to saturate at about 5%. Sidescatter was also

measured, but was determined to result from dielectric reflection from the overdense ionization front caused by the pump laser rather than from SBS. The conclusion was drawn that many of the features of the scattered light were attributable to the ionization front and subsequent heating produced by the pump laser, as well as initial plasma conditions which varied during the sequence of experiments.

Because the plasmas were preformed using methods not involving the laser which was used to drive the SBS, it is probable (though it was rarely measured) that the conditions of the plasma were quite different than those of laser-ablated plasmas. In fact, in instances where there was a delay between the end of the preform pulse and the start of the laser driver, it is probable that the plasma was rather quiescent at the start of the driver laser. Since plasma turbulence can lead to enhanced ion density fluctuations which can in turn enhance SBS, quiescent plasmas are favorable for studying SBS since the levels of ion wave noise are more predictable. This is particularly evidenced by the fact that many of these experiments reported quite low levels of SBS backscatter, suggesting that the noise levels of ion waves in the plasmas were at thermal levels. Of course this does not apply in cases where the probe laser itself creates plasma turbulence, as was postulated to occur in the experiments conducted by Herbst *et al.*^{8, 9}.

In contrast, there have been many experiments conducted with CO₂ lasers that relied on the laser radiation itself to form the plasma as it entered a gas¹²⁻¹⁹. In one experiment conducted by Burnett *et al.*¹⁴, in which the plasma density was below $n_c/2$ and the laser intensity was greater than 1×10^{14} W/cm², instantaneous levels of backscatter reflectivity exceeded 50% in some instances. It was postulated that the highly perturbative effect of the laser on the plasma resulted in wave breaking and subsequent enhanced

levels of density perturbation that were not readily predictable by fluid theory. Furthermore, it is also probable that these enhanced density perturbations were responsible for the high levels of observed backscatter reflectivity. Similarly, a series of experiments were conducted by Ng *et al.*¹⁹ in overdense plasmas created by a laser incident on a supersonic, high density gas jet. Backscatter reflectivities on the order of 20% were observed, and many phenomena including ion heating, beam filamentation, and ion turbulence were postulated to be responsible for the observations. Observed red and blue shifts confirmed that the observed light was not due to specular reflection from the critical surface. In other experiments with similar laser and plasma parameters, backscatter reflectivities were observed to saturate at levels below 10%^{15, 17, 18}, albeit at much lower interaction laser intensities.

There have been a few CO₂ laser experiments reported in which the plasmas have not possessed critical or quarter-critical surfaces²⁰⁻²³. In theory, these should simplify interpretation of SBS data since one can rule out reflection at the critical surface and other phenomenon occurring at the quarter-critical surface which can enhance SBS noise sources. Counterintuitively, most of these experiments observed levels of SBS in excess of those observed in similarly produced plasmas possessing these densities. Levels of backscatter reflectivity were observed from 10%²³ to complete (~ 100%) back reflection²⁰. It is crucial to note however, that in all but one of these experiments (Walsh *et al.*²³), it was concluded that the scattered light was the result of stimulated Compton scattering (SCS) rather than SBS. A more detailed discussion of SCS will follow later in this chapter.

As with microwaves, it was eventually determined that CO₂ lasers made less-than-favorable ICF drivers, primarily for three reasons. First, like

microwaves, the penetration of 10.6 μm laser light from CO_2 lasers into the very dense core of the ICF fuel was severely limited by the relatively low critical density ($\sim 1 \times 10^{19} \text{ cm}^{-3}$), though clearly this problem was not as severe as it was for microwaves. Second, the amount of energy deposited into the plasma perturbations induced by the laser electric field scales as the square of the quiver velocity, and the quiver velocity, in turn scales linearly with laser wavelength. Thus, the relatively long emission wavelength of CO_2 lasers tends to produce rather turbulent plasmas. Finally, the coupling of the CO_2 laser to the plasma often resulted in the production of hot electrons, which were a threat to ICF since they can penetrate into the core of the plasma and preheat the fuel before compression is accomplished.

3.3 SBS experiments conducted with solid-state lasers

By the mid-1970's, several solid state laser systems had been built and were operating with sufficient reliability to conduct ICF experiments. The dominant laser medium used in these lasers was neodymium-doped glass (Nd:glass). The *SHIVA* laser facility at Lawrence Livermore Laboratory and the *OMEGA* laser facility at the University of Rochester were the two most energetic of that era. The 1.06 μm emission line of Nd:glass lasers was a large improvement over the 10.6 μm emission line of CO_2 lasers, at least in terms of attaining the goal of inertial confinement fusion. By the 1980's technology had improved sufficiently to develop solid crystals that possessed unique nonlinear optics properties which allowed harmonic frequency conversion at sufficiently high energy densities for practical experiments. Consequently, there were many experiments conducted at laser wavelengths of 0.527 μm ($2\omega_0$), 0.351 μm ($3\omega_0$) and even 0.264 μm ($4\omega_0$).

By the time solid-state lasers dominated ICF research, most experiments did not rely on preforming plasma by other means. And in fact, by this time, virtually all of the ICF-driven experiments started with solid material: thin foils, glass microballoons, solid disks and rods, or cavity-structured or hohlraum targets²³⁻⁴³. However, there is at least one notable exception. Mayer *et al.*⁴⁴ conducted backscatter experiments with a Nd:glass focused onto a puffed nitrogen gas jet. This experiment differed distinctly from most of the other solid-target experiments of its era since the scale lengths of various plasma parameters were much larger than with solid targets. Density scale lengths in excess of 50 laser wavelengths were easily attainable with this technique. It is speculated that this is what lead to high levels of observed direct backscatter at about 40%, where the scattering was found to saturate for incident intensities around 2×10^{15} W/cm². There were a few anomalous experiments within this campaign where backscatter reflectivity was found to exceed 60%. It is also important to note that during the time in which the laser was driving SBS, it was concluded that the plasma was nearly stationary. This could be crucial since the widely held and used convective theory predicts that gradients in plasma flow tend to stabilize SBS growth. Finally, note that spectral analysis of the backscattered light showed that a large fraction of the observed fluence was attributable to specular reflection from the critical surface in the plasma rather than SBS.

In fact, for the vast majority of experiments that have been conducted to study ICF and SBS in particular, the plasmas have possessed critical surfaces. Since plasmas produced in an ICF reactor would always possess such densities, these experiments most closely resemble reactor targets. However, as with the previously mentioned experiment, this can confuse an

SBS measurement since the scattered light from SBS is at nearly the same frequency as the laser, which makes it very difficult to distinguish between contributions from specularly reflected light and SBS. In fact, for many of the experiments^{31, 33, 34, 37, 43, 44}, after specular reflection at the critical surface was accounted for, it was concluded that SBS contributed to only a very small fraction of the backscattered light. In one such experiment conducted by Tanaka *et al.*³⁷, the foil targets were tilted with respect to the laser axis, with the goal of tilting the specularly reflecting critical surface. The result is that the direct backscatter from SBS is scattered at a significantly different angle than the specularly reflected laser light. It was concluded that SBS backscatter was a relatively small contribution ($< 5\%$) to the total amount of scattered light. Others drew similar conclusions based on spectral analysis of the scattered light.

Many experiments were conducted where the angular variation of scattered light was measured^{24, 31, 35, 40, 41, 43}. In a series of experiments conducted by Drake *et al.*²⁴ on thick gold planar targets with $0.527\ \mu\text{m}$ laser light, SBS sidescatter was measured to exceed 10% of the incident laser light. In a similar experiment conducted by Tanaka *et al.*³⁵ where glass microballoons were irradiated with $1.06\ \mu\text{m}$ laser light, sidescatter which was attributed to SBS exhibited levels of scatter in excess of 50% of the incident light. Whereas Drake's experiments possessed planar plasmas, Tanaka's experiments possessed spherical plasmas. In both experiments, the angular distribution and spectral analysis of the scattered light was used to confirm the origin of scattered light. These experiments suggest that under these circumstances, SBS may be a convective instability.

Because of the complications posed by the presence of a critical surface, there were a few experiments conducted where the peak density was below

the critical density for the pump wavelength^{27, 28, 42}. Labaune *et al.*²⁷ conducted thin foil experiments with two different interaction wavelengths (0.527 μm and 0.264 μm) in an attempt to understand the wavelength scaling of SBS in underdense laser plasmas. Despite the relatively long density scale lengths ($L_n/\lambda_0 \approx 100 - 200$) that existed in these plasmas, and the relatively high irradiation intensities ($\leq 2 \times 10^{15} \text{ W/cm}^2$), backscatter reflectivity did not exceed about 7% for 0.527 μm light and 0.1% for 0.264 μm light. These low reflectivities were attributed to collisional absorption, which is more efficient at shorter wavelengths.

There were a couple of experiments conducted which measured the coupling between SBS and the two plasmon decay (TPD) instability. TPD is a parametric instability in which a light wave (laser) decays into two electron plasma waves. Since the frequencies of the electron plasma waves are nearly ω_{pe} , the plasma frequency, conservation of energy dictates that TPD occur very near the quarter-critical density for the light wave. In experiments conducted by Tanaka *et al.*⁴² the SBS instability was driven with 0.351 μm light in plasmas whose peak density was $0.66 n_c$. Peak backscatter reflectivities of 8% were measured. Intensities as high as $1 \times 10^{15} \text{ W/cm}^2$ were incident on the underdense plasmas, but it was concluded that both thermal and ponderomotive self-focusing, which can significantly increase the beam intensity within the plasma, contributed significantly to the observed results. Further, observation of the $3\omega_0/2$ emission from the plasma led to the conclusion that there was direct coupling between SBS and two plasmon decay (TPD). Similarly, an experiment conducted by Young *et al.*²⁸ was designed specifically to determine the dependence of SBS on the presence and eventual disappearance of both a critical- and quarter-critical-surface. Measuring the temporal dependence of the scattered light as the peak plasma

density evolved first below critical density, then below quarter-critical density, several conclusions were drawn. First, reflected light from the critical-surface seeded SBS, which in turn enhanced the levels of observed backscatter. Second, the beating of electron plasma waves produced by TPD resulted in enhanced ion wave noise, which further seeded SBS. After the peak density dropped below quarter-critical, levels of SBS were similar to those predicted by convective theory.

3.4 Stimulated Compton Scattering

There were many experiments conducted, some of which were discussed above, where stimulated Compton scattering (SCS) was concluded to contribute significantly to the observed scattered light^{20, 21, 29}. As such, a brief discussion of SCS will now be given.

Whereas SBS is the resonant decay of a light wave into another light wave and an ion acoustic wave, there are certain circumstances that can occur in these plasmas where it is inappropriate to think of the resulting density perturbation as a wave. In the limit that the phase velocity of the ion acoustic wave, v_{ph} , is much greater than the thermal speed of the ions, v_{th} , the ion acoustic waves are weakly damped and can grow to significant levels—SBS can become significant. However, in the limit that the ion acoustic phase velocity approaches the ion thermal speed, the bulk of the ions can efficiently Landau damp the ion acoustic waves. Scattering in this limit is primarily a particle kinetics effect, with a fraction of the energy from each scattering event being directly transferred to the particles.

For ion acoustic waves, Landau damping scales exponentially with the quantity

$$-\frac{v_{ph}^2}{v_{th}^2} = -\frac{ZT_e}{T_i} - 3 \quad (3.1)$$

where $T_e(T_i)$ is the electron (ion) temperature. It is clear that in the limit $ZT_e \gg T_i$ Landau damping will be negligible. During laser plasma experiments, the relatively low inertia of the electrons relative to the ions often results in such a limit, since it takes a finite period of time for the electrons and ions to equilibrate (become isothermal). However, when the electrons and ions eventually do equilibrate in temperature, Landau damping can be significant. In all of the experiments discussed above, SCS was found to occur late in the experiment, after a period of time sufficient to attain the condition $ZT_e \approx T_i$.

3.5 Summary

In conclusion, despite the plethora of experiments that have been conducted to study SBS in laser-driven plasma experiments, none has possessed all of the attributes necessary to clearly test convective theory. Experiments involving high laser intensities violated the linear theory's requirement of a weakly perturbed plasma ($v_{os} \ll v_{th}$) and most likely resulted in enhanced levels of ion density fluctuations ($\delta n/n$). Further, experiments conducted with $10.6 \mu\text{m}$ light (CO_2 lasers) or microwaves usually violated the linear theory's requirement of a weakly driven instability ($\gamma_0 \ll \omega_{ia}$). Experiments conducted with plasmas produced by ablation of solid material and driven with solid-state lasers typically possessed critical or quarter-critical density surfaces, which tend to confuse SBS measurements by introducing other methods of scattering light near the laser wavelength

(e.g. dielectric reflection at the critical surface) or again, by enhancing ion density perturbations (e.g. TPD-driven ion acoustic waves at the quarter-critical surface). Finally, many experiments in which the plasma was produced from a gas discharge or gas puff, the plasmas were stationary which prevents a test of the convectively stabilizing effect of inhomogeneous flow. In light of this, we have conducted weakly-driven solid-state laser experiments in plasmas with no critical or quarter-critical surfaces which also possess gradients in plasma flow. Further, titanium plasmas were chosen as a testbed to exploit the relatively high atomic number ($\bar{Z} \approx 20.5$) to quench Landau damping, consequently allowing SBS to grow significantly above the thermal noise level. By using a $0.527 \mu\text{m}$ interaction beam at relatively low intensities, these experiments fulfilled the limiting criteria of weakly-perturbed plasma and weakly-driven instability. Subsequently, these are the first experiments which clearly test the convective theory of SBS in laser-produced plasmas.

3.6 References

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Chapter 4

Titanium Thin Foil Experiments on NOVA

In the previous chapter it was shown that there have been many experiments over the past 15 years that were conducted to gain a better understanding of stimulated Brillouin scattering in laser-driven plasmas. However, each of these experiments were in regimes where it was difficult to clearly test convective theory. In most cases, especially in CO₂ laser-driven experiments, the laser parameters did not conform to the assumptions made in the development of convective theory, namely weakly-driven SBS in a weakly-perturbed plasma. Many others were plagued by characteristics which allowed other phenomena to occur, which may have either seeded or suppressed the growth and production of SBS. Finally, some made insufficiently quantitative study of the laser and plasma parameters they produced for the results to be sufficiently significant.

Consequently, it was the primary goal of this experimental campaign to devise a set of experiments where all of the necessary laser, plasma, and

diagnostic instrumentation requirements were combined to clearly test convective theory of SBS. Specifically, these experiments were conducted in plasmas which a) were relatively homogeneous over many laser wavelengths in size, ($L_n/\lambda_0 \sim 10^3$, where L_n is the plasma density scale length) b) did not possess critical- or quarter-critical surfaces, c) were weakly-perturbed by all lasers and d) possessed finite uniform flow gradients. Further, the SBS was weakly-driven in the low intensity experiments that were conducted in this campaign. Finally, in an effort to determine the plasma density and laser intensity scaling of SBS, experiments were conducted at two different peak plasma densities and two different laser intensity regimes that differed by an order of magnitude. Each experiment's role in studying the scaling of SBS can be found in table 4.1. This unique combination of experimental characteristics not only served as a test bed for convective theory, but also served the secondary goal of studying SBS in plasmas similar in nature to what will be produced by next-generation laser fusion facilities. Hence, this experimental campaign simultaneously addressed both fundamental plasma physics and practical issues concerning laser-fusion.

4.1 The Targets

In order to produce plasmas with the necessary parameters to address the goals of the experimental campaign, unique targets were developed. This would not have been practical without the previous target development experience gained in previous experiments on NOVA. In addition, much of the design of these targets was driven initially by extensive modeling with LASNEX¹, a large and complex hydrodynamics code developed and run at Lawrence Livermore National Laboratory.

LASNEX is a two-dimensional, cylindrically symmetric, hydrodynamics code developed and extensively used at Lawrence Livermore National Laboratory primarily for modeling the behavior of laser-produced plasmas¹⁻³. LASNEX uses a movable grid (Lagrangian) to track the evolution of the plasma at consecutive time steps. LASNEX uses a flux-limiter to account for inhibited free-streaming energy flow. LASNEX has the ability to calculate plasma conditions without assuming local thermodynamic equilibrium. LASNEX provides the scientist with two-dimensional "snapshots" of the plasma density, electron and ion temperature, and flow field, as well as various other physical parameters. LASNEX was used extensively in this experimental campaign, both in the initial design of the targets and in subsequent predictions of SBS behavior.

Titanium thin-foil targets were chosen for these experiments. There was initial concern that Landau damping could damp a sufficient fraction of the ion acoustic waves so that SBS would never reach detectable levels. Titanium was chosen to minimize ion Landau damping (a physical picture of Landau damping is given in the next section) by maintaining the quantity $|-ZT_e/T_i|$ large, since Landau damping scales exponentially with this quantity. A schematic of the target is shown in figure 4.1. The titanium foils were 1.5 mm in diameter, and were either $0.5\mu\text{m}$ or $1.5\mu\text{m}$ thick, depending on which peak density was desired for a particular experiment. These thin foils were suspended on formvar substrates which were 1000^{+000}_{-200} Å thick. LASNEX predicted that the extremely mass-limited formvar substrates would have no impact on the hydrodynamics, plasma or atomic physics. A special "butterfly mount" was devised to support the titanium-formvar assembly, and was so-named because of its resemblance to the wings of a butterfly. The butterfly mount consisted of a shaped, nominally 1mm

diameter wire frame and a vertical section that was mounted to the target positioner assembly. The "wings" of the butterfly were angled forward to form a three-dimensional assembly that resulted in a flat, two-dimensional foil.

The butterfly mount served two purposes. First, in order to measure the expansion of the plasma in a direction along the original target normal, it was necessary to allow for a line-of-sight of the target that was "edge-on" with respect to the thin foil. Since the wire used to form the frame was necessarily much thicker than the titanium foil itself, the wire would have completely occluded any edge-on viewing diagnostic had the wings of the butterfly mount not been swept forward. Second, target alignment on NOVA is accomplished with an orthogonal three-camera system, two of which are located in the target chamber equatorial plane. For these experiments, the necessary orientation of the target with respect to the laser beams resulted in the plane of the target coincident with the target chamber equatorial plane. Thus, without the forward-swept wings, translational alignment tolerance in the equatorial plane would have been limited to the thickness of the wire frame, in this case $\pm 500\mu\text{m}$. At first, this may seem to be sufficiently accurate for the purposes of the campaign. However, results from previous experiments suggest that the spectral shape of scattered light is significantly dependent on the location of the foil with respect to target chamber center, since deviation from dead-center can result in a radiation imbalance on either side of the foil and a subsequent global drift of the resulting plasma. The extremely narrow bandwidth (10's of Å) inherent in light scattered by SBS makes these Doppler shifts significant in the observations.

4.2 The Lasers

The plasma was formed using eight (8) of the ten arms of the NOVA laser facility. The NOVA laser beams are all incident upon target chamber center at a 50° polar angle. Hence the beams all lie on the surface of two opposing cones which possess a 50° half-angle, five on each cone with the West beams numbered odd and the East beams numbered even. The beams are equally spaced in angle along each cone, resulting in beams that are separated in azimuth by 72° . Beams on opposite sides are not "interwoven"; that is, beams 5 and 6 for example, are contained in the 270° azimuthal plane. For these experiments, four (4) frequency-tripled ($3\omega_0$, $0.351\ \mu\text{m}$) beams were incident simultaneously on each side of the titanium foil. This particular wavelength was chosen for enhanced plasma penetration and absorption, which was necessary to achieve the required large, low density, high temperature plasma. Each beam was defocused to form an overlapping spot at the center of the target, resulting in a quite uniform illumination circle that was approximately $950\ \mu\text{m}$ in diameter. LASNEX predictions and previous experiments suggested that this $950\ \mu\text{m}$ diameter illumination on the 1.5 mm diameter foil would result in a sufficiently planar plasma⁴. The beams were 2.1 ns full width at tenth-maximum (FWTM) "flat-top" pulses that had rise and fall times of approximately 100 ps. Total energy contained within the beams was nominally 22 kJ, with a resulting intensity at the surface of the foil of approximately $1.5 \times 10^{15}\ \text{W/cm}^2$. A typical set of preform laser pulses is shown in figure 4.2. Energy balance on each side of the target was maintained to within $\pm 15\%$ throughout the irradiation.

4.3 The Plasma

The eight "preform" beams quickly exploded the foil and heated the resulting plasma. A series of experiments was dedicated to diagnosing the plasma conditions. In these experiments, a ninth beam (beam line 6) was defocused to approximately $400\mu\text{m}$ in diameter at the center of the plasma and illuminated the plasma approximately 200 ps after the start of the preform beams. This diagnostic beam was operated at $0.527\mu\text{m}$, to distinguish resulting signals from those originating from the preform beams. Further, since it was the same beam and wavelength that was used to drive SBS during interaction experiments, it tracked the density profile relative to critical density in the same manner as during the interaction experiments. Since one of the main goals was to produce a plasma without a quarter-critical density surface for the interaction laser, it was important to be able to measure the density evolution with respect to $0.527\mu\text{m}$ during the characterization experiments. Using the same wavelength for the diagnostic beam during the characterization experiments and the interaction beam during the interaction experiments afforded this ability.

The evolution of the peak density was measured by observing the stimulated Raman backscatter produced from the $0.527\mu\text{m}$ diagnostic beam. This technique has been successfully employed in past experiments and has typically agreed with code predictions to within a factor of two in density³. Stimulated Raman scattering (SRS) is a parametric instability similar to SBS, in which an incident electromagnetic wave interacts with an electron plasma wave to create a scattered electromagnetic wave. Like SBS, it too must obey frequency and wave vector matching conditions (conservation of energy and momentum, respectively) for resonance, as well as dispersion

relations in order for the waves to propagate in the plasma. From the frequency matching condition, $\omega_{sc} = \omega_0 - \omega_{epw}$, and the electron plasma wave dispersion relation, $\omega_{epw}^2 = \omega_{pe}^2 (1 + 3k_{epw}^2 \lambda_{De}^2)$, it can be seen that as SRS occurs at higher densities, the wavelength of the scattered light wave increases. Thus, the longest wavelength scattered from the plasma originates from the highest density below quarter-critical density. Assuming one knows the electron temperature to sufficient accuracy, the long wavelength cutoff of the Raman emission is a direct measurement of the peak density in the well-underdense plasma.

As an example, figure 4.3 shows a comparison of the evolution of peak electron density calculated and measured via observation of the Raman long-wavelength cutoff. To model the plasma produced by the thin (thick) target, LASNEX assumed an initial $0.5\mu\text{m}$ ($1.5\mu\text{m}$) thick titanium foil, 1.5 mm in diameter, deposited upon a $1500\text{ \AA}$ thick formvar substrate. The foil was symmetrically irradiated for a 2 ns duration with $0.351\mu\text{m}$ wavelength laser light, focused to a $872\mu\text{m}$ diameter (FWHM) spot on each side of the foil, with a total energy deposition of 23 kJ (21.5 kJ). Then at $t=1.8$ ns, a 2 ns duration, $0.532\mu\text{m}$ wavelength interaction laser was incident upon the plasma and was focused to a $68\mu\text{m}$ radius (FWHM) spot at the center of the plasma. The energy contained in the interaction pulse was 1.25 kJ (1.5 kJ). A flux limiter of $f=0.03$ was employed to inhibit heat transport, and was chosen in light of favorable results from previous runs whose accuracy had been confirmed by experiment³. From figure 4.3 it is clear that there is a discrepancy between predicted and measured peak density that differs roughly by a factor of two. This discrepancy is consistent with many previous comparisons between experiment and LASNEX predictions^{3, 5, 6}. It is noteworthy that the evolution of the predictions and the experiments track one another; that is,

the shape of the temporal profiles are similar, despite the differences in amplitude. This is important for two reasons. First, a simple scaling factor suffices to resolve the discrepancy, and this is convenient when using LASNEX results to model certain effects analytically. Second, it leads to the possibility that the discrepancy may be due to an anomaly resulting from certain choices made in characteristics of the grid. Procassini and Estabrook have shown that modifications to the way in which energy is accounted for in the code can partially or completely diminish this discrepancy⁷. Since the LASNEX results represent a complete set of plasma predictions, they were used in unmodified form to produce the convective theory predictions that are presented in chapter 7, and the impact that this decision has on the results will be discussed in detail at that point in the dissertation.

The electron temperature can also be measured by observing SRS. Before one can understand this, one must understand Landau damping. Briefly, Landau damping is a collisionless, resonant damping that occurs between electrostatic waves and particles in a plasma. A simple physical explanation is given by Kruer⁸, and is given here in slightly modified form for the purposes of this particular discussion. When plasma particles experience an oscillating electric field, they tend not to absorb any net energy unless collisions occur, allowing the oscillation energy to be converted to random scattered or thermal energy. However, in the regime where the phase velocity of an electrostatic wave nearly matches the thermal speed of the plasma particles, a large fraction of the plasma particles are resonant with the wave—that is, these particles experience nearly constant electric fields which can either accelerate or decelerate the resonant particles, depending on the relative phase. Particles with speeds slightly less than the phase velocity are accelerated and gain kinetic energy, whereas particles with speeds

slightly faster than the phase velocity tend to be decelerated and lose kinetic energy. If there are more particles whose speeds are slower versus faster than the phase velocity, a net damping of the wave can occur. Such is always the case in a Maxwellian plasma since this distribution of particle speeds monotonically decreases with increasing speed.

Landau damping of electron plasma waves is considered significant for waves whose wavenumber, k_{epw} , satisfies the relation $k_{epw}\lambda_{De} \geq 0.3$ ⁹, where $\lambda_{De} = \sqrt{k_B T_e / 4\pi n_e e^2}$ is the electron Debye length. Since these short wavelength electron plasma waves will be significantly damped, the intensity of the Raman scattered light associated with these particular waves will be subsequently very low. If one assumes that the shortest wavelength (taken here to be 10% of the peak spectral intensity) occurs for $k_{epw}\lambda_{De} = 0.3$ then an observation of the shortest wavelength of SRS emission as defined above is a direct measurement of the peak electron temperature. A comparison of the calculated temperature and that inferred from the short wavelength cutoff for a typical experiment is given in figure 4.4.

An 8× magnification soft x-ray microscope which employs grazing-incidence cylindrical mirrors for imaging was employed to observe the symmetry of the resulting plasma density, by measuring the bremsstrahlung (free-free) emission from the plasma. There are four separate channels, each of which are filtered differently to be sensitive to different energy bands. The channel used to collect the data presented here (8XE2) had a 76μm thick PVC bandpass filter and the focusing mirror was coated with titanium. In addition to the bandpass filtering, attenuation of the signals was performed with a 50μm thick beryllium filter, which also serves as a blast shield to protect the instrument from plasma and solid debris. This combination of filters results in an on-axis transmission of approximately 12% at an x-ray

energy of approximately 2.3 keV. The shape of the resulting transmission curve is asymmetric, with -3dB points (the wavelengths where transmission has decreased by half) at approximately 1.8 and 2.5 keV. The view of this telescope is at an angle of 25° from the target normal. Figure 4.5 shows a typical result from the $8\times$ telescope, where the image has been corrected for the offset in polar angle. In other words, the image presented is estimated to be that which would be observed if the $8\times$ telescope were looking along the original target normal. From this diagnostic, it is clear that axial symmetry is quite accurate for this type of experiment.

Diagnostics were used to measure peak density, temperature, and symmetry. However, other, equally crucial plasma characteristics were not measurable. The typical temperatures and scales of laser-plasma experiments typically precludes the use of traditional plasma diagnostic techniques. Specialized techniques such as Thomson scattering of laser light have been developed elsewhere, and other schemes such as imaging of bremsstrahlung emission from the plasma have been suggested, but these have not been employed on Nova primarily due to cost and complexity. For this reason, LASNEX predictions were extensively employed to predict the detailed behavior of the plasma.

4.4 The Interaction Laser

The same laser (beam line 6) and wavelength ($\lambda_0 = 0.527\mu m$) that was used as a diagnostic probe during the characterization experiments was later used as the interaction laser during the SBS experiments. This section will focus on its use as an interaction beam.

The interaction beam was operated at the frequency-doubled ($2\omega_0$) wavelength of $0.527\mu\text{m}$ so that signals originating from it could be separated in frequency from those originating from the $0.351\mu\text{m}$ preform beams. Operating in the visible region of the spectrum also made it possible to make extensive use of fiber optics in the diagnostics used to measure SBS. During the SBS experiments, the laser was focused with $f/4.3$ optics to a $175 \pm 30\mu\text{m}$ diameter spot at the center of the plasma, where the beam achieved its maximum intensity. Because the NOVA beams are not diffraction limited, the length of the beam waist along the axis of the laser is rather long—approximately $500\mu\text{m}$ on either side of “best focus”. Here, the length of the beam waist is defined as the length over which the diameter of the spot remains within $175 \pm 30\mu\text{m}$. The spot diameter is defined as that diameter within which 80% of the delivered power is contained, and intensities are calculated based upon this definition. It is important to note that there is significant transverse structure within this 80% power circle, again attributable to the less-than-perfect quality of the NOVA laser optics. Figure 4.6, shown here compliments of the author of reference¹⁰ shows a photograph of the intensity profile at best focus of a similar but different NOVA beam, operated at $0.527\mu\text{m}$. This structure may play a significant role in the seeding of instabilities such as SBS, and will be discussed further in the final chapter.

Estimates of the time-resolved laser beam energies after an experiment are inferred from a measurement of the temporal profile of the laser pulse using a relatively-calibrated optical streak camera, and a separate time-integrated measurement of the laser energy early in the amplifier chain using an absolutely-calibrated photodiode and previously-acquired KDP frequency conversion efficiencies. The calibration of the photodiodes and the frequency conversion efficiencies are obtained by periodic, dedicated

calibration experiments which employ whole-beam calorimeters within the vacuum chamber. Laser intensities are then obtained from these time-resolved energy measurements and previously performed transverse beam profile measurements. Intensities typically quoted are based on 80% of the total power delivered, and the radius of the 80%-inclusive power contour of the transverse profile of the beam. Typical uncertainties in the inferred intensities are $\pm 10\%$ for $0.351\mu\text{m}$ laser light and $\pm 25\%$ for $0.527\mu\text{m}$ laser light. Though the uncertainty in $0.527\mu\text{m}$ laser intensity may seem large, it is well within acceptable limits for the experimental campaign reported here, particularly since the campaign spans over two orders of magnitude in laser intensity.

For these experiments, interaction beam intensities varied from $4 \times 10^{13} \text{ W/cm}^2$ to over $5 \times 10^{15} \text{ W/cm}^2$, in an effort to obtain information on the scaling of SBS with laser intensity. The interaction beam was turned on approximately 1.8 ns after the start of the preform beams, resulting in a 200 ps overlap between the preform and interaction beams. Since the effective atomic number of the plasma was predicted to be $Z \approx 20.5$, the bremsstrahlung emission rapidly and efficiently cooled the plasma. By overlapping the preform and interaction beams, the electron temperature was ensured to remain high at least during the onset of SBS, again to significantly suppress Landau damping.

The interaction pulse shapes are a consequence of the design of NOVA. Since NOVA is seeded by a single oscillator, all ten beams are fed from the oscillator with an identical pulse shape. In order to deliver high intensity, 2 ns square preform pulses to the target at $0.351\mu\text{m}$ as shown in figure 4.2, the oscillator had to emit a pulse which ramped up in intensity during the pulse. This was necessary to compensate for saturation of the laser amplifiers.

Consequently, the interaction beam was also seeded with this same pulse shape. For the high intensity SBS experiments, the interaction pulse was also 2.0 ns square. However, for the low intensity experiments, the interaction pulse delivered to the target ramped up in intensity similar to the 1ω oscillator seed. Figures 4.7.a and 4.7.b show the temporal evolution of typical high and low intensity interaction beam experiments, respectively.

4.5 Summary

The above techniques were employed successfully to produce a plasma environment which was conducive to the study of SBS. Experiments previously conducted by others were often successful in measuring significant levels of SBS. However, for various reasons, the results did not offer a conclusive test of convective theory. In contrast, the results obtained from these experiments are not clouded by complications such as specular reflection from a critical surface, seeding at the quarter-critical surface due to two plasmon decay (TPD), or from plasma turbulence. Further, SBS is weakly driven in the low intensity experiments, a requirement for a true test of convective theory. Finally, complex instrumentation, discussed in the next chapter, allowed a detailed analysis of the temporal and spectral behavior of the scattered light originating from SBS at several angles simultaneously. The results of these experiments will be discussed in the following chapters.

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	Low Intensity Laser ($\sim 4 \times 10^{14}$ W/cm ²)	High Intensity Laser ($\sim 4 \times 10^{15}$ W/cm ²)
Low Density Plasma ($n_{\text{peak}} \leq 0.05 n_c$)	21070916 21111511 21111516	21070914 21071911
High Density Plasma ($n_{\text{peak}} \leq 0.15 n_c$)	21062616 21111514 21071915	21062614 21062618 21071913 21111508

Table 4.1 Nova shot numbers corresponding to all of the Titanium Foil experiments conducted to study SBS. Numbers in bold face indicate acquisition of streaked-spectra with MATRS. Those in normal face produced no streaked-spectra, but did produce time- and wavelength-integrated SBS scattering fluences.

Figure 4.1 Titanium foil target. The 1.5 mm diameter (0.5-1.5 μm thick) foil was mounted on formvar as a support. The butterfly mount allows diagnostics to have a clear edge-on line-of-sight of the plasma as it explodes and expands.

Figure 4.2 Temporal profiles of eight separate "preform" laser pulses operated at 0.351 μm .

Figure 4.3 A comparison of LASNEX and measured plasma density predictions for: (a) a thick target and (b) a thin target. The solid lines represent the LASNEX predictions. The measurements are based on observations of the long-wavelength cutoff of the stimulated Raman scattering (SRS) emission both in forward and rearward scattering, taken during several experiments dedicated for this particular purpose. It is clear that LASNEX predicts a consistently lower peak plasma density. However, the discrepancy is within a factor of two, making LASNEX predictions a valuable tool for analysis of SRS observations.

Figure 4.4 A comparison of measured and LASNEX-predicted electron temperature evolution. LASNEX predictions for thick target (#211106a) and thin target (#211207a) are shown. Measurements, based on the short-wavelength cutoff of SRS at two different wavelengths, were made only for a thick target.

Figure 4.5 Iso-intensity contours of soft x-ray "snapshots" of preformed plasmas obtained with the 8X microscope. The contours roughly represent curves of equal density integrated along the line-of-sight of the instrument. These plots appear here to indicate the qualitative symmetry obtained along the original target normal. Examples are given for (a) a thick target (#21111514) and (b) a thin target (#21111511).

Figure 4.6 Iso-intensity contour plot of the transverse profile at best focus of a Nova beam operated at 2ω . Clearly there is transverse structure in the beam, and depending on the design of the experiment in question, this structure may have impact on the observed results. This image was obtained at the 2-Beam facility at NOVA. Data courtesy of Paul Wegner.

Figure 4.7 Second-harmonic (2ω) temporal pulse shapes for a typical (a) low intensity and (b) high intensity experiment. Note that in order to obtain an order-of-magnitude difference in the intensity, the output pulse shape must change to account for saturation effects in the amplifier stages.

Chapter 5

MATRS: A New Instrument on NOVA for Measuring Absolute Levels of Time-Resolved Optical Spectra at Multiple Angles Simultaneously[†]

5.1 Introduction

The study of inertial confinement fusion has been one of the driving forces in the development of high temperature plasma diagnostics. To achieve fusion, it is necessary to understand the coupling of the driver to the plasma. In laser plasma experiments, many processes occur which can scatter part of the incident laser light, and much can be learned by measuring this light. Two such processes are stimulated Raman scattering (SRS) and stimulated Brillouin scattering (SBS), three-wave

[†] The material in this chapter will be submitted, in significantly modified form, to The Review of Scientific Instruments.

parametric instabilities in which an incident photon (the pump) decays into a scattered photon and an electron plasma- or ion acoustic phonon, respectively. This chapter describes a new instrument, the Multiple-Angle Time-Resolving Spectrometer (*MATRS*), developed to study these processes in plasmas pertinent to laser fusion. The discussion will describe the instrument in its most recent configuration, which was optimized for measurements of SBS driven by laser light at $0.527\text{ }\mu\text{m}$.

In recent years, there have been many new instruments developed, especially in the field of astronomy, which have been based upon Multiple Object Spectroscopy (MOS)¹. MOS is a technique which employs a single spectrograph and detector to simultaneously record many light sources by locating and distributing the signals at the object plane of the spectrograph. MOS techniques grew out of a need for the astronomy community to optimize the usefulness of large and expensive telescopes and associated detectors^{2, 3}. The development of fiber optics has made this technique both simple and cost effective, and there are many fiber-optic-based MOS's in use today³⁻¹¹.

Three phenomena must be considered when using optical fibers in instrumentation: attenuation, dispersion and the numerical aperture (NA) of the fiber being used¹². Attenuation and dispersion in fiber optics are usually measured in dB/km and ns/km, respectively. Since the lengths typically employed in laboratory instrumentation are short ($\ll 1\text{ km}$) these issues are rarely of concern. The NA of a fiber is a measure of its acceptance or emittance angle, is inversely proportional to f -number, and must be accounted for both in the design of the input and output optics. Light outside the acceptance angle will not be transmitted, and detector optics with higher f -numbers than the output f -number will be overfilled,

resulting in a loss of signal. A well studied effect associated with optical fiber NA is focal ratio degradation (FRD), whereby microbends in the core/clad interface result in a ratio of output to input angles greater than unity ($NA_{out} > NA_{in}$). FRD has been observed to increase with increasing input f -number¹³⁻¹⁵. Consequently, to avoid over-filling the output optics, a small f -number should be employed for fiber input. Conversely, an input f -number approaching the NA of the fiber results in loss of higher modes and subsequent attenuation of the signal. To minimize these competing effects, an optimum input f -number of $f/3$ to $f/4$ is typically employed¹.

For MATRS to be a useful diagnostic on NOVA, it needed to fulfill the contradictory virtues of being inexpensive to produce, simple to align and operate, require little diagnostic port space on the vacuum vessel, and transmit the detected signals to a remote location for detection by a single instrument. As such, a fiber-optic based MOS was an ideal candidate for the design of MATRS.

5.2 Diagnostic Design and Development

5.2.1 Light Collection

A series of six light collectors were assembled and distributed throughout the inside of the vacuum vessel to collect the scattered light at many different angles with respect to the laser beam. Figure 5.1 shows a schematic of the instrument as it is installed on NOVA. The location of each collector was chosen to be adjacent to a pre-existing photodiode which had been previously absolutely calibrated at $0.527\text{ }\mu\text{m}$, the interaction beam wavelength, which is also the nominal wavelength of the measured SBS.

Fluences measured by these photodiodes were used to absolutely calibrate MATRS for each experiment. Each collector assembly consists of a 100 mm diameter, 350 mm or 400 mm focal length quartz lens, two turning mirrors, an optical diffuser, and a 30 cm "pigtail" of 50 μm diameter step-index fiber. Each lens subtends a solid angle of approximately $\pi \times 10^{-3} \text{ sr}$ with respect to the center of the plasma. The lens was positioned to create a fiber input f -number of roughly $f/4$. For compactness, the image path is "folded" by two aluminum front-surface mirrors. The fiber input was placed at the focal plane of the lens.

Many experiments on NOVA produce fluences in excess of 200 J/sr, resulting in an energy density at the entrance of the 50 μm diameter fiber exceeding $3 \times 10^4 \text{ J/cm}^2$. Typical damage thresholds for optical fiber are several of orders-of-magnitude less than this, so filtering had to be incorporated in the design of the light collectors. To date, MATRS has been used in experiments employing two different laser wavelengths (0.351 μm and 0.527 μm). To study scattered light near 0.527 μm a low pass filter (CorningTM GG400) with a cutoff of 400 nm was used to eject the scattered 0.351 μm light, preventing it from reaching the fiber input. Further attenuation of the input signal above 400 nm was then accomplished with an optical diffuser placed in front of the fiber input, whose primary purpose is discussed below.

Without the optical diffuser, the field of view of the instrument was calculated to be approximately 200 μm at the center of the target, for a nominal demagnification of 4. It is believed that for most experiments, this is roughly the interaction length in the plasma. Such a narrow field of view required extremely accurate alignment of the instrument to ensure that all scattered light from the appropriate region of interest reached the fiber

input. To increase the field of view of the instrument, several techniques were tested. One possibility was to increase the size of the image plane, necessitating a larger diameter fiber. Large diameter fibers (2 mm diameter) and fiber tapers (2 mm to 100 μm) were tested in the laboratory. These proved to be extremely sensitive to input f -number and input angle, were relatively expensive, and were susceptible to optical and target debris damage. In contrast, an optical diffuser is inexpensive, relatively immune to damage and insensitive to input signal characteristics. Of course, the effective attenuation of an optical diffuser is much larger than either of the other two alternatives, but this proved to be of no consequence in the experiments conducted thus far.

A 25 mm diameter fused silica diffuser was produced by bead blasting the front surface of the substrate, and was placed approximately 3 cm in front of the fiber input. The effective attenuation of the diffuser at this distance was measured in the laboratory to be approximately ND 2.5 and the scattering was estimated to be lambertian. The addition of the diffuser increased the effective field of view to greater than several mm, far larger than the length of the resonant region of the plasma, as was confirmed by detection of laser light specularly reflected from the target stalk which is located at least 2.5 mm away from the center of the target.

5.2.2 Signal Transport

Light from each collector was transmitted over a 9 meter length of 50 μm diameter, step index, unjacketed fiber to the vacuum chamber wall. Unjacketed fiber was chosen for use inside the vacuum chamber to prevent out-gassing and ensure cleanliness. To protect the unjacketed fibers from

optical crosstalk and from ballistic target debris, a sheath of tinned copper braid was passed over each fiber. Signal were transmitted out of the vacuum chamber using a multi-fiber vacuum feedthrough, and were then finally transmitted to a remote diagnostics platform over a 16 meter length of optically-identical, jacketed fiber.

The fiber used was Corning 100/140 SDF™ (Short Distance Fiber) graded index multimode optical fiber, with a 100 μ m (140 μ m) diameter core (cladding). The manufacturer's measurements of attenuation (3.59 dB/km) are made at 850 nm. In contrast, MATRS operates in the 527 nm regime. However, one can extrapolate reasonable and conservative values of attenuation at the wavelength of interest for design purposes. Assuming an attenuation coefficient at 527nm of ≤ 10 dB/km, then the 30 meter length of fiber through which each signal passed experienced a maximum attenuation of 0.3 dB. Since it is estimated that each fiber junction (of which there are at least two) accounts for 3 dB loss, the attenuation due to the fiber itself is insignificant. A simple estimate for the chromatic dispersion can be made as well. Though the core index of refraction is measured and stated by the manufacturer at two separate wavelengths outside the operating range of MATRS [$n_{850} = 1.4805$ and $n_{1300} = 1.4748$, where n_{λ} indicates the index of refraction at test wavelength λ (given in nm)], one can assume that since over this range, the index of refraction changes less than 6 parts in 1000, then over the 10 nm range that MATRS operates in, the index of refraction should change even less. In fact, it is safe to assume that in a 10 nm wavelength range centered at 527 nm, the index of refraction changes far less than 1 part in 1000. In general, the difference in arrival times for two simultaneously launched signals of different wavelengths can be estimated by noting that

$$\Delta t = \frac{x}{c_0} \Delta n \quad (5.1)$$

where x, c_0 and Δn are the length of fiber, the vacuum speed of light, and the difference in index of refraction at the two wavelengths in question. For the values $x \sim 30m$, $c_0 = 3 \times 10^8 \text{ m/sec}$, $\Delta n \sim 10^{-3}$, one finds that the difference in arrival times for two signals with a difference in wavelength of 10 nm is $\Delta t \sim 1ps$, which is at least two orders of magnitude less than the temporal resolution of the instrument. So, neither attenuation nor chromatic dispersion from the fibers themselves proved to be a significant (or even detectable) issue.

Since variations in signal amplitude from separate collectors were expected to vary by many orders of magnitude, it was necessary to be able to filter the signal from each collector separately. For this reason, a small optical assembly consisting of two 25 mm diameter quartz lenses and a filter holder was produced for each separate signal path. Light exiting each fiber at the diagnostics table was collimated with the first lens, transmitted through the appropriate (aluminum-coated-glass) ND filter, and then refocused into another short length (0.5 - 1.5 m) of identical fiber which led to the spectrometer. Each filter holder had an effective attenuation of approximately ND2.5 without the use of a filter. Critical measurement of each attenuation source in the system was unnecessary since the photodiode adjacent to each collector inside the chamber allowed calibration of each measurement independently, for each experiment.

In addition to the signals from the six MATRS collectors, there is the existing capability on NOVA of collecting a portion of the light that is scattered directly back along the original path of the laser. Light collected

in this manner is transported to the diagnostics platform through a light pipe consisting of discrete geometrical optics¹⁶. In order to time-align the MATRS signals with the direct backscatter signal, approximately 15 meters of fiber was added to the direct backscatter channel of the spectrometer input.

5.2.3 Instrumentation

A total of seven simultaneously collected signals were multiplexed into a single 0.5 m Czerny-Turner spectrometer. The spectrometer has a magnification of approximately unity. A streak camera was located at the image plane of the spectrometer to temporally resolve the spectra¹⁷⁻¹⁹. The streak camera, in its current configuration, has a 25 mm x 100 μ m input slit and a sweep rate of approximately 0.35 ns/mm. In second order for 0.527 μ m light, the spectrometer has a spectral dispersion of approximately 0.8 nm/mm over a spectral range of approximately 24 nm, which is limited in the current configuration by the length of the spectrometer exit slit. The microchannel plate (MCP) intensifier located at the output of the streak camera can resolve two events separated by approximately 100 μ m. However, focusing limitations of the electron optics in the streak tube (which are optimized for temporal resolution) and the finite dispersion of the electron currents emitted from single channels in the MCP are what limit the temporal and spectral resolutions to approximately 125 ps and 0.4 nm, respectively. The full duration of the sweep is approximately 10 ns. Typical experiments produce signals with bandwidths less than 10 nm and durations less than 4 ns, allowing the distribution of many signals across one streak record, as will be discussed in greater detail shortly.

The ends of the fibers were distributed and equally spaced at approximately 0.1 inch intervals in an aluminum multi-fiber holder which formed the input slit of the spectrometer. Figure 5.2 shows a photograph of the multiplexed fiber array. The multi-fiber holder was mounted to a positioning stage with three independent degrees of freedom. This stage was used to align the fibers as a single assembly, and the translation stage was mounted as a unit to the spectrometer housing for stability.

Distribution of the signals over the object plane of the spectrometer resulted in a similar distribution of spectrally resolved signals at the image plane of the spectrometer. A spatially infinitesimal or "point" light source with finite bandwidth located at the object plane of a spectrometer results in a spectrally resolved image (spatially extended in the transverse dimension) at the image plane. Of course, finite-sized sources are spectrally resolved to even larger images. Without the proper precautions, two closely-spaced light sources at the object plane can result in overlapping spectra at the image plane. The size of each object in MATRS is $50\mu\text{m}$ as defined by the diameter of the input fiber. To avoid spectral overlap of adjacent images, different fiber lengths were used for each signal to advance and retard the signals' arrival at the spectrometer input by up to ± 5 ns.

A narrow multilayer bandpass filter (approximately 10 nm FWHM, center wavelength at 525 nm) was placed inside the spectrometer directly in front of the fiber ends to prevent thermal and other stray broadband light sources that entered the light collectors from confusing interpretation of the measurement. Since the signals in these experiments proved to be very large, any loss due to overfilling of the spectrometer optics from FRD was

insignificant. Finally, the streaked spectra were recorded on black and white film for purposes of archival and analysis outside the laboratory.

5.2.4 Data Manipulation

Data that is recorded on film is digitally converted with a scanning photodensitometer for computer-assisted visualization and quantitative analysis. In order to prepare the data for this process, a calibrated neutral density wedge and light source are used to record a "transfer" wedge on the side of each page of film before an experiment. This results in a calibration signal that incorporates the response of the film to the spectrum produced by the image intensifier. After digitization of the data, a lineout of the wedge is obtained to convert film density to relative exposure. Finally, relative exposure for each data signal is converted to absolute spectral intensity by equating the total exposure (in "counts") to the total time- and wavelength-integrated fluence measured by the calibrated photodiodes located at the MATRS light collectors. The result is an absolutely calibrated, streaked image giving spectral intensity (power per unit solid angle per unit wavelength) of the SBS detected vs. time.

Algorithms for visualization and quantitative analysis of the data were written in a proprietary computer programming language, Interactive Data Language (IDL), and were executed on a VAX mainframe computer and various DEC workstations and graphics terminals. Appendix B gives an example of one of the algorithms used to process MATRS data. A computer-generated false-color-representation of a typical selection of MATRS data is shown in figure 5.3. The result of temporally and spatially multiplexing the signals into the streaked spectrometer is

clearly indicated. The signal located in each polygon represents the time-resolved SBS spectrum recorded at a different angle, here denoted with respect to direct forward of the interaction beam. Note that although streak camera time increases monotonically upward and spectrometer wavelength increases monotonically to the right, the signal at each separate angle possesses its own separate wavelength and time coordinate system, by virtue of both the imaging properties of the spectrometer and the configuration of the MATRS fibers at the entrance slit of the spectrometer.

Figure 5.4 illustrates many of the features of a typical piece of MATRS data. Scattered $0.527\text{ }\mu\text{m}$ laser light from the target stalk at the start of each experiment is recorded by each MATRS channel and serves as both an absolute wavelength and time fiducial for each individually recorded signal. This early $0.527\text{ }\mu\text{m}$ light originates from residual $1.054\text{ }\mu\text{m}$ light in the $0.351\text{ }\mu\text{m}$ plasma “preform” beams that was converted to $2\omega_0$ rather than $3\omega_0$. As explained in chapter 4, this off-harmonic light does not reach the plasma and therefore does not confuse the experiment.

5.2.5 Alignment

Alignment of each collector is performed in the laboratory before it is installed inside the vacuum vessel. A helium-neon (He-Ne) tabletop laser tuned to emit in the green ($\lambda_0 = 543.5\text{ nm}$, proprietarily named a “Gre-Ne”) is employed with a gradient index optic at the output aperture to efficiently couple the laser light into a $100\text{ }\mu\text{m}$ diameter fiber. The output of the fiber is rigidly mounted to an optical bench and each collector is placed at a distance identical to the distance between the target and the collector during an experiment. After appropriately adjusting the optics, each

collector is then installed inside the vacuum chamber, and is manually precision aligned *in situ* before the chamber is evacuated.

The light source used for precision alignment of the collectors *in situ* is the same used in the laboratory. A series of fixtures was developed to hold the fiber output, which is connected to the target positioner and is aligned to target-chamber-center (TCC) in a fashion identical to that used to align a target prior to an experiment. Because each MATRS collector is located at a different elevation on the inside chamber wall, and the target positioner has only one rotational degree of freedom, each fixture was produced with a different tilt angle so that the light emanating from the fiber at TCC would fill the optics of each collector. Once the fixture with the appropriate tilt is rotated to the appropriate angle for alignment of a particular collector, final precision alignment is then accomplished by adjusting the turning mirrors in the collector assembly, and is verified by temporarily removing the optical diffuser and visually inspecting throughput at the end of the pigtail, as well as remotely at the input to the spectrometer. After confirmation, the fibers are reconnected to the instrumentation and the diffuser is replaced.

5.3 Experiments

A series of experiments to study SBS in large underdense laser plasmas have recently been completed in which MATRS was employed as a primary diagnostic. Underdense plasmas were created by symmetrically exploding thin (0.5 - 1.5 μm) titanium foils (1.5 mm diameter) with eight beams of the NOVA laser facility. These preform beams had a wavelength of 0.351 μm , contained a total of approximately 22 kJ, and formed a 950 μm

diameter spot on target. Approximately 1.8 ns after the start of the preform beams, a 2 ns duration interaction beam was focused at the center of the plasma. The interaction beam had a wavelength of $0.527\ \mu\text{m}$, had a spot diameter of $<200\ \mu\text{m}$, and was varied in energy (intensity) from approximately 200 to 2500 Joules (3×10^{14} to $4 \times 10^{15}\ \text{W/cm}^2$). At the start of the interaction beam, the peak density was measured to be $0.05 - 0.15n_c$ and had an electron temperature of approximately $T_e \approx 3\text{keV}$. Significant levels of SBS were observed in many of the experiments. The data collected with MATRS has permitted a clear test of the convective theory of SBS in plasmas pertinent to laser fusion. The details of these experiments are reported in the remainder of this dissertation.

5.4 Conclusions

A new diagnostic, MATRS, has been successfully fielded at the NOVA laser facility. MATRS gives experimenters the new capability of recording absolute levels of temporally- and spectrally-resolved scattered optical light at up to seven angles simultaneously during an experiment. Initial experiments have proven MATRS to be a useful and robust instrument for the study of laser-plasma interactions in a parameter regime significant to laser fusion.

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Figure 5.1 MATRS on NOVA. Red cones represent individual light collectors distributed throughout Nova target chamber. Each assembly collects and transports light emitted from the plasma over a single optical fiber to a diagnostics platform containing a single spectrometer and streak camera. Absolutely calibrated photodiodes adjacent to each MATRS collector allow calibration of signal collected by each assembly. Additionally, geometrical optics collect light scattered directly back along the laser path.

Figure 5.2 MATRS multiple-fiber spectrometer entrance slit. Fibers from all MATRS collectors are distributed rectilinearly at the original location of the spectrometer entrance slit, forming a new effective entrance slit. Photograph (a) is the view of the new entrance slit assembly as would be seen from the inside of the spectrometer. The ends of the fibers have been outlined in white for emphasis. Photograph (b) shows the bundling of the single fibers that form the new entrance slit assembly. This assembly can be translationally adjusted in three orthogonal directions for precision alignment of the assembly at the entrance port of the spectrometer. Differing fiber lengths retard or advance arrival of various signals to prevent spatial overlap of spectrally resolved images at the exit slit of the spectrometer.

Figure 5.3 A typical MATRS streak page. The signal inside each rectangle represents the time-resolved spectrum observed at the angle that is indicated just below each signal. Global streak camera time increases upward and global spectrometer wavelength increases to the right. However, because of the method of layout of the MATRS signals on the streak page, each signal has its own separate wavelength and time coordinate system.

Figure 5.4 Details of an SBS streaked-spectrum recorded with MATRS. At $t=0$, residual 2ω light from the 3ω preform beams scatters off of the target stalk and into the MATRS collectors, denoting both $t=0$ and $\lambda=0.527 \mu\text{m}$. For a complete description of this effect, see chapter 2, as well as figure 2.5. At $t=1.8 \text{ ns}$ the interaction beam is turned on and the abrupt start of SBS is easily noted. At $t=3.8 \text{ ns}$ the interaction beam ceases and the SBS turns off, as expected.

Chapter 6

Observation of Stimulated Brillouin Backscattering at Far Above Convective Levels in Very Underdense, Inhomogeneous Laser Plasmas[†]

6.1 Introduction

Stimulated Brillouin scattering (SBS) is a three-wave parametric instability whereby an electromagnetic (em) wave decays into an electromagnetic wave and an ion-acoustic (ia) wave, and can occur in any medium which supports acoustic waves. SBS is of fundamental interest since it can suppress propagation of electromagnetic waves, as well as compete with absorption, filamentation and other instabilities¹. In

[†] The material appearing in this chapter has been submitted, in slightly modified form, to Physical Review Letters for publication.

plasmas² SBS can degrade implosion efficiency of laser-fusion targets, and broaden emission lines of photopumped x-ray lasers due to ion heating. SBS can be quite complex in inhomogeneous plasmas, since amplification can occur anywhere within the corona, and since various mechanisms can increase amplification through feedback or enhanced noise.

Past studies of SBS have been reviewed by Kruer³ and by Baldi *et al.*,⁴ and have included experiments using non-laser-produced plasmas, plasmas produced by ionization of gas jets, ionization of foams or thin foils, or by ablation of solid matter. In most previous experiments using inhomogeneous plasmas, the observed spectrum and amplitude of SBS may have been affected by reflection⁵ of laser light from the critical density (n_c), enhanced ion-wave noise produced at the quarter-critical density ($n_c/4$) as a consequence of two-plasmon decay^{6, 7}, and ion-wave noise resulting from turbulence driven by heat transport. Other experiments were in strongly-driven regimes⁸⁻¹¹ in which the growth rate of the instability in a homogeneous plasma, γ_0 , is much larger than the ion-acoustic frequency, ω_{ia} , or in strongly-perturbed regimes in which the electron oscillatory speed, $|\vec{v}_{os}| = |e\vec{E}_0/m\omega_0|$, is much greater than the electron thermal speed, $v_{th} = \sqrt{T_e/m}$. Here, e, m and T_e are the electron charge, mass and temperature (in energy units), and $\vec{E}_0$ and ω_0 are the laser electric field and frequency, respectively. These complications permit at best a tenuous comparison to the basic theory of SBS.

In contrast, the experiments discussed here unambiguously involve weakly driven ($\gamma_0/\omega_{ia} \sim 0.1$) SBS in a weakly-perturbed ($v_{os}/v_{th} \sim 0.05$) plasma. Further, these experiments include the effect of finite inhomogeneous plasma flow, which is thought to stabilize SBS. Finally, a

conclusive test of convective theory requires spectral measurements at more than two angles simultaneously. We report for the first time, experiments which possess the simultaneous attributes of a weakly-perturbed, very underdense plasma ($n_e/4$) with inhomogeneous flow, in which SBS is weakly-driven and in which multiple-angle, time-resolved spectra are obtained—all of which are necessary for a clear test of the convective theory of SBS.

6.2 Experiments

We produced plasmas for these experiments by symmetrically illuminating thin (0.5 ± 0.025 μm thick), circular (1.65 ± 0.1 mm diameter) titanium foils, mounted on 1000_{-200}^{+000} Å formvar supports, using eight laser beams of the NOVA laser facility ¹² to preform the plasma. A schematic of the experiment is shown in figure 6.1. Titanium was chosen to minimize the ion Landau damping by maintaining ZT_e/T_i large. All beams were incident upon the foil at 50° from the target normal and were focused with $f/4.3$ optics. These “preform” beams, of wavelength 0.351 μm , were defocused to produce a roughly uniform spot, 850 μm diameter full-width at half-maximum (FWHM), at the target plane. The total energy in the preform beams was nominally 22 kJ, with energy balance between both sides of the foil within 5%. The laser light was delivered in a 2 ns flat-topped pulse, with 100 ps rise and fall times.

Approximately 200 ps prior to the end of the preform beams, a 0.527 μm , 2 ns FWHM “interaction” beam, which ramped up in power during the pulse, was focused at the center of the underdense plasma to a 175 ± 30 μm

diameter spot containing 80% of the energy (based on measurements of other NOVA beams). The energy of the interaction beam was 0.16 ± 0.05 kJ for the data discussed here. At the start of the interaction beam, the plasma parameters were inferred from simulations using the LASNEX¹³ computer code to be $n_{peak} = 0.05n_c$, $T_e = 3$ keV, $\bar{v}_u = 7.7 \times 10^8$ sec⁻¹, $Z = 20.5$, and the FWHM of the density profile was 1.2 mm. These predictions were consistent with past experimental observations. The details of the experiment, including results of simulations and comparison to measurements that characterize the evolution of the plasma, are reported separately^{14, 15}.

The scattered light was measured at wavelengths near $0.527 \mu\text{m}$ using a streaked-spectroscopic system developed expressly for these experiments. The Multiple-Angle, Time-Resolving Spectrometer (MATRS) consists of six optical light collectors distributed in angle throughout the inside of the NOVA vacuum chamber¹⁶. There are three collectors in the forward ($\leq 90^\circ$) and three in the rearward ($\geq 90^\circ$) direction with respect to the interaction beam. Scattered light from the plasma is collected with refractive optics and directed out of the vacuum vessel via optical fibers to a 0.5 m Czerny-Turner spectrometer and an optical streak camera, where the ends of the fibers are equally distributed to form the entrance aperture of the spectrometer. Each collector is adjacent to a filtered, absolutely-calibrated photodiode for calibration. In addition to the six off-axis collectors, a fraction of the light scattered directly back along the original path of the interaction beam is also recorded¹⁷. In contrast to the off-axis signals, the direct backscatter signal is calibrated with a 2 inch diameter calorimeter and a fast photodiode, which permits isolation of the measured contribution due to plasma scattering from (early) beam reflections from

optics. Differing lengths of fibers are employed to retard or advance the signals, preventing spatial overlap of adjacent, spectrally-resolved images at the exit aperture of the spectrometer. In this manner, seven simultaneous, calibrated, time-resolved spectra were recorded on a single laser irradiation, with typical resolutions of 100 ps and 4 Å, although shifts in location of the peak as small as 1 Å may be detected.

6.3 Comparison of Data and Calculations

The data are compared to calculations based on linear, steady-state theory for plane waves interacting with a planar, inhomogeneous plasma^{2, 18}. In such theory, SBS is presumed to be convectively unstable and grow exponentially from thermal levels in the plasma as the disturbance propagates through the resonant region. Both bremsstrahlung and Thomson scattering of the probe laser contribute to the initial electromagnetic noise source¹⁹. Attenuation of the pump and daughter waves by the plasma within the resonant region reduce the gain exponent²⁰. All gradients are assumed to be along the original target normal and this leads, at most, to a factor-of-two error in the predicted scattered intensity. Calculation of the predicted spectral light flux follows from the previous work of Berger *et al.* using this reduced gain exponent¹⁹.

MATRS was used to measure the time-resolved spectra at four angles simultaneously in an experiment where the interaction beam in a single irradiation ramped up in intensity from 4×10^{13} W/cm² after the initial 100 ps onset to 4×10^{14} W/cm² at the end of the pulse. Figure 6.2 shows the results and a comparison to convective theory. For direct

backscatter (180°) the intensity levels shown are lower limits; values are obtained by taking the reflectivity of the optics to be 1%, which was measured during a dedicated calibration shot and is also consistent with consideration of the optical design. However, attempts to measure the total scattered energy with the calorimeter occasionally led to much larger inferred optics and plasma reflectivities. To date no flaws have been found in the measurement that would account for this.

Figure 6.2.a presents intensity contours of the time-resolved spectra at four rearward angles. Figure 6.2.b presents a comparison of the evolution of intensity measured at the wavelength of maximum observed emission to that predicted using convective theory. Calculations predict that SBS should not significantly exceed thermal levels. However, measurements show levels six orders of magnitude higher than thermal noise. Further, theory dictates that once thresholds are exceeded, SBS should reach full amplification in the time required for an ion acoustic wave to traverse the matching region (~ 50 ps), which is incidentally shorter than the minimum amount of time that the instrument can resolve. In contrast, in the experiments the levels of scattered light take several hundred ps to rise to their peak amplitude—far longer than both the rise-time of the laser pulse and the time-resolution of the detecting instrument. Figure 6.2.c presents a comparison of the measured spectra to predicted widths based on the length of the interaction-beam waist and $\bar{V}_u$. Predicted lineshapes will be reported in a future publication. The measured spectral widths are typically twice as broad as those predicted. Other, similar experiments produced similar spectra, time dependencies and amplitudes. Similar conflicts with convective theory from experiments

driving SBS with a 10 ps pump laser will be reported by Baldis *et al.* in a future communication.

6.4 Discussion

Though a distribution of intensities in the laser beam can lead to increased scattering, uniform beam intensities far greater than the average intensity are necessary to explain these data. As an example, for the case of direct backscatter at 2.3 ns for the experiment reported above, the uniform beam intensity would have to be enhanced by a factor of approximately 18. If one considers the consequence resulting from structure in the beam, then as an additional example, 10% of the beam would have to be enhanced by a factor of over 21 to account for the observations. An assessment of the measured intensity profile of a similar beam, coupled with modeling by a filamentation code²¹ and analytical predictions²², suggests that a significant fraction of the beam area could reach 10 times the nominal intensity. It is conceivable that the limited accuracy of the assumptions used in this assessment could nearly double this number. If so, it is then possible that a combination of initial beam structure (hot spots) and significant filamentation could explain the observations. An equally plausible alternative may be that somehow SBS backscatter becomes absolutely unstable, perhaps with the assistance of modest filamentation.

6.5 Conclusion

In conclusion, long-scale-length, very underdense plasmas have been produced for studying SBS in a regime that is of particular relevance to the laser-fusion community and is of particular fundamental interest. Though convective theory is often invoked to predict SBS amplitudes for future applications, these results suggest that this procedure is questionable; convective theory is unable to explain the large amplitudes observed near direct backscatter. Clearly, further work is necessary to develop an ability to predict levels of SBS in inhomogeneous plasmas.

6.6 References

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Figure 6.1 Schematic of exploding titanium foil experiments on NOVA. Foil is exploded with 8 NOVA beams (four on each side of the foil) at $0.351\mu\text{m}$, only two of which are shown here for simplicity. The plasma expands, then the experiment is conducted with the interaction beam, operated at $0.527\mu\text{m}$. Light scattered from the plasma is collected at 6 oblique angles and direct backscatter simultaneously during each experiment using MATRS.

Figure 6.2 Backscattered light at 140°, 150°, 160° and 180° with respect to direct forward of the interaction beam: (a) Spectral-intensity contours of time-resolved spectra, spaced by factors-of-two, with the peak contour at each angle at 90% of the maximum; (b) Measured and predicted evolution of maximum spectral intensity. The measured intensities (solid lines) are averaged over approximately 3Å around the peak and have an approximate temporal resolution of 100 ps. Signals detected before the start of the interaction beam (≈ 1.8 ns) are not shown since they originate from preforming beam reflections rather than SBS. The predicted intensities are shown as broken lines. (c) Measured spectra (averaged over ~ 120 ps) and predicted spectral widths at 2.25 ns (dark curves and bars, respectively) and 3.0 ns (light curves and bars, respectively). The broken vertical line represents the laser wavelength. The ends of the bars show the wavelengths predicted to originate from the locations in the plasma corresponding to the ends of the approximately 1mm long beam waist, assuming perfect focusing. This is experiment #21070916.

Chapter 7

A Comparison of Observed and Predicted SBS Lineshapes

7.1 Introduction

As was previously mentioned, the primary goal of this experimental campaign was to study SBS in a plasma environment which was both of interest to fusion experiments and goals and which provided results which were easily and readily interpretable. Since future fusion experimental facilities will produce very large plasmas (many mm) it was important to create a plasma which was relatively homogeneous over as large a distance as possible. In order for the results to be a clearly interpretable test of the convective theory of SBS, various plasma parameters such as density, temperature, and flow, and various laser parameters such as intensity, energy, and temporal pulse shape had to be carefully controlled. Detailed explanations of the roles of these various parameters appear in the previous chapters.

This experimental campaign spanned a large range of peak plasma densities ($0.05 - 0.15 n_c$) and vacuum laser intensities ($4 \times 10^{13} - 5 \times 10^{15} \text{ W/cm}^2$) in an effort to record and understand the density and intensity scaling of SBS. Table 7.1 shows the nominal density and intensity that each experiment addressed, and each experiment is listed by its Nova-assigned shot number within this table. Those shot numbers appearing in bold-faced type yielded observable streaked-spectra at three or more angles simultaneously with MATRS. Those in plain type were experiments which did not yield sufficient streaked-spectra; it is speculated that MATRS was not operating optimally during these experiments, most likely due to alignment error. It is important to note, however, that the absolutely-calibrated photodiodes used to calibrate the MATRS streaked-spectra were operable during those experiments where MATRS achieved insufficient performance. Results from the photodiodes during all experiments indicated that there was considerable shot-to-shot repeatability in the measured flux amplitudes (better than a factor of two for identical experiments), despite the fact that no streaked-spectra were recorded. Such repeatability adds credence to comparisons of theory and predictions in cases where streaked-spectra were recorded during only one experiment (e.g. shots #21070916 and 21062616).

A simple measure of the time- and wavelength-integrated light scattered at different angles with respect to the interaction beam would yield insufficient information to determine the validity of convective theory. That is, other mechanisms such as spectral reflection can produce scattered light at wavelengths near those produced by SBS. Further, spectral widths versus scattering angle can appear different for convective and absolute SBS theory. Thus, in order to truly confirm the nature of the scattered light, spectral measurements were necessary—hence the development of MATRS. A

measurement of the time-evolution of the scattered light at several angles can further illuminate interesting behavior—such information was reported for a typical low intensity, low density shot in chapter 6. This chapter will focus on the spectral nature of the scattered light at many angles for each of the four quadrants of plasma density and laser intensity achieved in the experimental campaign as shown in table 7.1. Spectra measured at two separate times during each of these four chosen experiments is studied, and these measurements are compared to spectra predicted by convective theory. Such a comparison indicates the ability of convective theory to predict the observed results.

7.2 Convective Theory Predictions

Calculations were performed to determine the spectral shapes and intensities predicted by convective theory. In order to perform such calculations, many plasma parameters had to be known. As with any experiments, it was nearly impossible to measure all necessary plasma characteristics to sufficient accuracy to rely solely upon them in a theoretical prediction. As explained in chapter 4, the evolution of the peak density and peak electron temperature were estimated from the Raman rearward and forward scatter. However, a complete calculation requires knowledge of density, electron and ion temperature, and flow *profiles* and these quantities' gradients along the phase gradient at each point on the laser axis, as well as their evolution in time. Such a complete set of plasma properties are currently immeasurable in such plasmas, and there is much research currently taking place within the laser-plasma community to improve our ability to make such measurements. Since the current "state of affairs" is

thus, a decision was made to obtain these quantities through modeling of the plasma using the complex hydrodynamics code *LASNEX*¹. A description of *LASNEX* and the results it provides the experimental plasma physicist have been discussed in chapter 4. It was shown in chapter 4 that there is a significant discrepancy between *LASNEX*-predicted and Raman-inferred peak densities and temperatures. To analyze how this discrepancy affects lineshape predictions requires knowledge of the linear convective theory, so this analysis will be given after the theory is presented later in this chapter.

Figures 7.1 and 7.2 show the density, electron temperature, ion temperature, and plasma flow along the original target normal for a low density (0.05 μm thick target) simulation at 2.3 ns and 3.0 ns, respectively. Figures 7.3 and 7.4 show these same quantities at the same times for a high density (0.15 μm thick target) simulation. These quantities are taken along the original target normal for several reasons. First, these values are more readily available to the experimentalist in search of such results. Second, convective theory for a 1-D linear plasma depend on components of the plasma properties in the direction of the phase gradient as defined in chapter 4, assumed in these experiments to be along the original target normal. Since the interaction laser is not incident along the original target normal, the ideal calculation would involve plasma quantities along the laser axis. To account for this discrepancy, the calculations assume that all plasma quantities are "1-D" in nature along the original target normal, and scale lengths are adjusted accordingly for the laser's angle of incidence with respect to the density gradient (see appendix C). Based on these corrections and on a 2-D analysis of the *LASNEX* predictions, it is estimated that the subsequent error in predicted plasma properties along the laser path are within a factor-of-two².

The observant reader will notice that there is axial structure in the density and temperature profiles predicted by LASNEX. It is questionable whether this structure is real or accurate. It is unlikely that they are an artifact of the numerical method, since the size of the structure far exceeds the size of a typical numerical "cell". If this structure does arise from some phenomenon that is real and not numerical in nature, it is still questionable whether LASNEX accurately predicts their amplitude or location. The experimental campaign was explicitly designed to minimize such structure in the plasma (e.g. plasma turbulence). However, it was not possible to experimentally measure the success of this goal.

It is therefore appropriate to consider the impact that this or similar structure may have on the theoretical predictions reported below. One would expect that any perceptible effect would originate from vanishing gradients or "flat spots" in plasma properties, especially plasma flow, since gradients in the flow are considered to be the dominant stabilizing factor for SBS. Note, however, that LASNEX predicts very uniform flow gradients (figures 7.1 through 7.4) for these plasmas. Thus, one would predict that structure in plasma density and temperature would have a diminishing impact on the predictions of convective theory. It is interesting to note that a recent theory of Rose *et al.*³ suggests a mechanism whereby flat spots in density may be produced from hot spots in the laser. And, it is possible that flat spots in flow could result indirectly from these flat spots in density, as is discussed in greater detail below. So, despite their questionable origin or existence, the LASNEX structures are retained in the calculation as a means of studying the impact that such structure may have in the theoretical predictions.

In addition to extensive use of LASNEX predictions of plasma quantities and their gradients, the SBS flux calculations also relied on upstream⁽¹⁾ temporal measurements of the absolute intensity of the fundamental frequency laser light during each experiment and previous measurements of the second harmonic axial beam profile obtained from experiments which were dedicated for this purpose⁴. Figure 7.5 shows the temporal evolution of the inferred second harmonic power for the interaction beam in each of the four shots being analyzed in this chapter. Figure 7.6 shows the previously measured axial profile of the 2ω interaction beam and the “best fit” used in the calculations (this figure appears courtesy of Wegner⁴).

The SBS flux calculations are based on the convective theory briefly presented in chapter 2 (and supporting work in appendix A) and are similar to those explained in the previous chapter. The calculations were performed on a custom-developed series of Microsoft ExcelTM spreadsheets that were executed on various MacintoshTM platforms. The spreadsheets appear in both numerical and expression forms in appendix C. The calculations were performed at each of 80 separate points along the laser axis to determine the predicted scattered light wavelength and observable intensity outside the plasma resulting from each point. Various other quantities were calculated at each point as well, such as various instability thresholds and other secondary parameters necessary to complete the flux predictions. The wavelength and predicted intensity at each point in the plasma was then combined to form the predicted lineshapes which appear in the comparisons of predictions and observations found in section 7.4 of this chapter.

The expressions used to calculate the predictions are derived from the coupled mode equations, equations (2.25) and (2.26), where it has been

assumed that the instability is weakly driven ($\gamma_0/\omega_{ia} \ll 1$) in a weakly perturbed ($\tilde{v}_{osc}/\tilde{v}_e \ll 1$) plasma. These assumptions are key to appreciating the unique nature of the work presented herein, as is explained many times in previous chapters of this thesis. Despite the fact that most experiments to date have violated one or both of these assumptions, convective theory is often used to explain the results. These are the first experiments which do not violate these assumptions.

After significant manipulation of the coupled mode equations⁵, one finally arrives at the now-famous Rosenbluth expression^{6, 7}, shown here in slightly modified form⁸ to account for distinct sources of electromagnetic noise

$$I_{SBS} = I_{brems} e^G + I_{Thomson}(e^G - 1) + I_{other}. \quad (7.1)$$

The quantities I_{brems} and $I_{Thomson}$ are contributions to the initial noise source from bremsstrahlung and Thomson scattering of the pump from ion fluctuations in the plasma, respectively, I_{other} is simply the thermal (bremsstrahlung) emission beyond the resonant region that does not contribute to SBS, and G is the so-called gain factor, given by

$$G = \frac{2\pi\gamma_0^2}{K'v_{1x}v_{2x}} F(\sqrt{v_1v_2}/\gamma_0) \quad (7.2)$$

where K' is the first derivative of the wave vector mismatch, $K \equiv |\vec{k}_0 - \vec{k}_1 - \vec{k}_2|$, with the derivative taken along the direction of the mismatch, $\hat{x} \equiv \vec{K}/|\vec{K}|$, and

$$\gamma_0 = \frac{1}{4} \frac{k_2 v_{osc} \omega_{pi}}{\sqrt{k_2 \omega_1 c_s}} \quad (7.3)$$

is the growth rate for SBS in a homogeneous plasma, with $k_2, \omega_1, \omega_{pi}, v_{osc}$ and c_s being the ion acoustic wave number, scattered light frequency, ion plasma frequency, electron quiver velocity, and sound speed, respectively. Further, $\bar{v}_{1x}(\bar{v}_{2x})$ is the component of the group velocity of the scattered light wave (ion acoustic wave) along $\hat{x}$. For a plasma where the stabilizing plasma parameter changes linearly with position (in this case, the plasma flow) one can write the derivative of the wave vector mismatch as⁹

$$K' = \frac{\omega_{pe}^2}{2L_n c^2} \left(\frac{k_{2x}}{k_{0x} k_{1x}} \right) + \frac{c_s}{v_{2x}} \left[\frac{1}{2} \frac{|k_2|}{(1 + k_2^2 \lambda_D^2)^{3/2}} \left(\frac{1}{L_T} + \frac{k_2^2 \lambda_D^2}{L_n} \right) - \frac{k_{2x}}{L_u} \right] \quad (7.4)$$

where the scale lengths of density, temperature and flow are given by

$$L_n = \left[\frac{1}{n_e} \frac{\partial n_e}{\partial x} \right]^{-1}, \quad L_T = \left[\frac{1}{T} \frac{\partial T}{\partial x} \right]^{-1}, \quad L_u = \left[\frac{1}{c_s} \frac{\partial u}{\partial x} \right]^{-1} \quad (7.5)$$

respectively. Finally, $F(\sqrt{v_1 v_2}/\gamma_0)$ is an effective damping coefficient given by Williams¹⁰ which accounts for damping of the pump and daughter waves inside the resonant region, and is given by

$$F(y) = \frac{2}{\pi} \left[\cos^{-1}(y) - y \sqrt{1 - y^2} \right] \quad (7.6)$$

and $y = \sqrt{v_1 v_2}/\gamma_0$, with $v_1(v_2)$ being the damping rate of the scattered light wave (ion acoustic wave). Work by Berger *et al.*⁸ shows the bremsstrahlung and Thomson scattering noise sources to be

$$I_{brems} = \frac{cT_e}{\lambda^4} E_1 \quad (7.7)$$

$$I_{other} = \frac{cT_e}{\lambda^4} E_2 \quad (7.8)$$

and

$$I_{Thomson} = \frac{cT_e}{\lambda^4} E_3 \left[\frac{\omega_1}{2c_s k_2} \right] \left[\frac{\cos^2 \theta}{\sqrt{\cos^2 \theta - (\omega_{pe}^2 / \omega_1^2)}} \right] \quad (7.9)$$

where E_1, E_2 and E_3 account for the emission and absorption in the appropriate regions of the plasma for each expression, and where θ is the scattering angle, measured with respect to $\hat{x}$. Further, equation (7.9) assumes that $ZT_e/T_i \gg 12$, with Z being the average charge state of the ions, and this is certainly fulfilled in all of the titanium foil experiments reported here, as was discussed in earlier chapters.

As stated in the reference, it is important to note that $I_{Thomson}$ can easily dominate I_{brems} since

$$I_{Thomson} \propto I_{brems} \frac{\omega_1}{c_s k_2} \quad (7.10)$$

and for SBS in particular, $\omega_1 / (c_s k_2) \gg 1$. Finally, expression (7.1) is calculated in the spreadsheet listed in Appendix C, as is described in detail below.

To predict the SBS spectral flux from these plasmas, the calculation proceeds first by determining the resonant scattered-light and ion acoustic frequencies and wave vectors that are produced at each position in the

plasma along the laser axis. These daughter-wave parameters are uniquely determined by the plasma conditions of density, temperature, and plasma flow, the dispersion relations for electromagnetic and ion acoustic waves in an unmagnetized plasma, the matching conditions discussed in chapter 2, and the angle of observation with respect to the pump wave vector. Because these plasmas consist of symmetric density and temperature profiles and antisymmetric flow profiles, there is a unique scattered light frequency predicted for each position in the plasma. This one-to-one correspondence translates in the laboratory as a unique observed wavelength outside the plasma for each position in the plasma where SBS is sufficiently driven to produce observable levels of scattered light.

Once the daughter-wave parameters are determined for each position in the plasma, the calculation proceeds by determining the predicted Rosenbluth gain at each position. To complete the calculation, the appropriate initial light flux or "noise" coefficients is calculated at each position from the coefficients in expression (7.1) above. Once this initial noise level is determined, the resulting scattered light flux originating from each position and its corresponding scattered light frequency are then combined to produce a prediction of the observed spectral lineshape outside of the plasma. These predicted lineshapes are then compared to the absolutely-calibrated spectra obtained in the experiments at the observation angles of interest. Once again, the algorithm for calculation appears in appendix C.

Before proceeding further it is necessary to determine the effect that errors in LASNEX predictions may have on the convective theory predictions. As was discussed in chapter 4, there is as much as a factor-of-two discrepancy between LASNEX-predicted and Raman-inferred peak plasma density and temperature. It is clear from equation (7.1) that the predictions are most

sensitive to density and temperature in the exponential gain factor, G . The gain factor varies weakly with temperature compared to other plasma parameters, so errors in peak temperature would have a weak effect on the predictions. The gain factor varies with the quantity $n_e/\nabla u$ where n_e is the electron density and ∇u is the plasma flow gradient. If the peak electron density is twice that predicted by LASNEX, as the Raman measurements suggest, then the hydrodynamics would dictate that the flow gradient is also larger by some proportional amount. Thus, one would expect that any increase in the gain factor would be less than a factor-of-two.

However, as a worst-case scenario, consider the effect of doubling the gain factor. For large values of the gain factor ($G \gg 1$) the predicted light flux would drastically increase. In contrast, for gain factors on the order of unity, as is the case for the low intensity experiments reported below, doubling the gain factor would have a much less pronounced effect on the predictions. Certainly this alone cannot account for the six order of magnitude difference in flux levels reported in the last chapter. And, although a systematic increase in the gain factor will result in a narrower predicted lineshape, it will soon be shown that such a systematic change in the predictions has almost no impact on the conclusions that are drawn from the comparison. Finally, one would expect to predict a slightly more red-shifted lineshape from an increased plasma density. However, the magnitude of this shift would be on the order of $1-2\text{\AA}$, which is much less than the width of the predicted lineshape and is significantly less than the resolution of the measured spectra. Therefore, it is believed that errors in the peak density and temperature predicted by LASNEX have an insignificant impact on the convective theory predictions that are reported below.

In addition to calculating the predicted lineshapes, instability thresholds are also determined for each position in the plasma. Thresholds for these processes are stability criteria; below a physical threshold, the process may exist but not grow above the noise level. However, once the threshold intensity is exceeded, the process becomes unstable. There are several different thresholds that can be predicted for parametric instabilities. Thresholds can exist either because there is collisional or other forms of damping of the daughter waves (absolute or convective), or because the resonant triad (the pump and two daughter waves) "detunes" as it propagates through an inhomogeneous plasma so significantly that the instability produces insufficient gain to be observable above the background noise (exclusively convective). Detuning of the triad outside the resonant region is typically thought of as a form of damping since it stabilizes SBS and thus limits the total gain of the daughter waves.

The threshold condition for SBS that is growing convectively unstable as it travels through a resonant region is given by^{11, 12}

$$\gamma_0 \geq \sqrt{\nu_1 \nu_2} \quad (7.11)$$

where $\nu_1(\nu_2)$ is the collisional energy damping rate for the scattered light wave (ion acoustic wave). In contrast, the threshold condition for SBS that is growing absolutely in time at a particular location is given by

$$\gamma_0 \geq \frac{1}{2} \left\{ \nu_1 \sqrt{\frac{\nu_2}{\nu_1}} + \nu_2 \sqrt{\frac{\nu_1}{\nu_2}} \right\} \quad (7.12)$$

where $v_1(v_2)$ is the group velocity of the scattered light wave (ion acoustic wave). In both (7.11) and (7.12), the equality indicates threshold for stability and exceeding the inequality indicates instability. Because both of these conditions arise from damping of the daughter waves, it is obviously appropriate to consider these damping thresholds.

In contrast, it is somewhat misleading to consider the threshold arising from convective detuning as a damping threshold, though it is frequently referred to as such. Close scrutiny of the Rosenbluth gain expression indicates that for any non-vanishing gain factor, G , the spectral intensity produced will exceed the initial noise, which is given by the coefficients preceding the exponential in expression (7.1). However, since very small values of the gain factor produce practically unobservable increases in the scattered light flux, a traditional convective "threshold" has evolved in the literature and in practice. Historically, since the Rosenbluth gain factor was originally written^{6, 13} as $G = 2\pi g$, this phenomenological threshold is taken to be $g \geq 1$ ($G \geq 2\pi$). Hence for convective gains exceeding ~ 535 [$\approx \exp(2\pi)$], SBS is said to have exceeded "convective threshold". Of course, this use of the term "threshold" is a misnomer since below this value, the process is still by definition unstable, albeit less detectable. Finally, it is important to note that it is possible for SBS to be above threshold in some locations in these inhomogeneous plasmas and below threshold at other locations. This, of course, will have a significant impact on the spectrum of observed light and may offer insight into the experimental observations.

Note that in general, all experiments presented here were driven with laser intensities which far exceeded (at least by an order of magnitude) the convective threshold intensity. For the low intensity experiments reported, the laser intensity was anywhere between a factor of two and a factor of ten

below the absolute threshold intensity, depending upon the particular experiment, the time during the observation, or the location within the plasma where the SBS was occurring. For the high intensity experiments, the laser intensity clearly exceeded the absolute threshold intensity at some particular times and locations in the plasma. This makes the high intensity experiments a poor test of convective theory, as discussed further in the discussion section of this chapter. We shall soon find that these high intensity experiments failed to test convective theory for additional reasons.

7.3 Measurements of Scattered Light from SBS

For the experiments conducted in this campaign, the temporal evolution of the SBS spectra were recorded at seven angles simultaneously for each experiment using the new instrument *MATRS* (see chapter 5 for a detailed description of *MATRS*). This dissertation is concerned with those observations in the rearward direction: for scattering angles greater than 90° with respect to direct forward of the interaction beam. As explained previously, the temporally-resolved spectra obtained are absolutely calibrated by measuring the absolute levels of time- and spectrally-integrated light at each angle with an appropriately filtered, absolutely-calibrated photodiode (oblique angles) or absolutely-calibrated calorimeter and relatively-calibrated photodiode (direct backscatter). These absolute measurements are then equated to the film intensity of the *MATRS* data, resulting in absolutely calibrated streaked spectra.

It was determined after the completion of the experimental campaign that the direct backscatter calorimeter was operating improperly and had thus produced erroneous results. For this reason, the information collected

with the relatively-calibrated photodiode was used exclusively to determine lower limits of absolute backscattering for calibration of the direct backscatter measurements made with MATRS. A more detailed explanation of this alternate method of analysis now follows.

The original purpose of the backscatter diode was to obtain a relative measure of the contributions of direct backscatter energy originating from plasma scattering and scattering from optics in the beam line. Such a measurement would then allow the experimenter to remove the optical scattering contribution from the total energy recorded with the calorimeter. Since the calorimeter gave erroneous results, this goal was unobtainable. Instead, it was assumed that the contribution of the optics scattering was approximately 1% of the incident beam energy. This is a conservative estimate in light of the optical layout, and is consistent with the lowest measurements in experiments dedicated to this purpose. This assumption allows absolute calibration of the portion of the diode signal that is due to optical reflections, easily identified by its early time-of-arrival. Then, a direct comparison of the signal strengths (on the same diode trace) from the optics scattering and the plasma scattering yields an absolute measurement of the plasma reflectivity, within the error of the assumption. This results in a lower limit in plasma reflectivity because the assumption of 1% optics scattering is a lower limit. Based on this method of calibration, a higher optics scattering (which was often observed in dedicated experiments) would result in a higher plasma reflectivity. Thus, the absolute spectra shown below in the direct backscatter direction are most likely conservatively low estimates of the spectral flux that was actually observed. Figure 7.7 illustrates the origin of this problem in greater detail.

The next section of this chapter will discuss the observations of scattered light at four separate angles (140° , 150° , 160° , and 180° with respect to direct forward of the interaction beam) for experiments that probed low and high peak density, and low and high interaction beam intensity. First, the interaction laser power evolution for each experiment and the measured axial beam profile will be given for purposes of analysis, followed by a diagram depicting the reduced accuracy in the backscatter measurements. Next, iso-intensity contour plots of the streaked spectra at four separate angles for an experiment are shown. Finally, immediately following each set of contour plots, spectral lineouts at two different times during the experiment (2.3 ns and 3.0 ns) will be compared to convective theory predictions (low intensity experiments) or discussed in their own right (high intensity experiments). In all figures, the contours represent successive factors-of-two, the highest-valued contour representing 90% of the peak spectral intensity observed *at that particular angle*, and the measured spectra are averaged over approximately 120 ps.

7.4 Observations

7.4.1 Low Intensity Experiments

Figure 7.8 shows the iso-intensity contours for a low density plasma experiment (#21070916) in which the total laser energy was measured to be less than 157 J. It is clear from figure 7.5.a that the laser intensity at best focus increased from less than 5×10^{13} W/cm² during the first 100 ps of the pulse to over 3×10^{14} W/cm² within the last 200 ps of the pulse. The reason for this “high contrast” (>10) pulse was discussed earlier in chapter 4 but it is

interesting to note that such a pulse offers the potential to observe and record an intensity-scaling of the instability in these experiments since the time evolution of the instability is recorded with MATRS.

Figure 7.9 shows comparisons of the observed and predicted spectra at two separate times during the experiment. At 2.3 ns the nominal laser intensity at best focus is approximately $5 \times 10^{13} \text{ W/cm}^2$ and at 3.0 ns the intensity has increased to approximately $1 \times 10^{14} \text{ W/cm}^2$. Convective theory predicts that early in time the spectra should be smooth gaussians with red-shifts of approximately 1 nm, with approximately 2 nm full-width at half-maximum (FWHM). Also, these predicted spectra are at or near the predicted thermal levels of scattered light from the plasma, since the exponential gain factors predicted are $G \leq 1.7$. In contrast, the experiments show structure in the spectra, with the dominant wavelength at or very near the interaction laser wavelength. At 140° and 160° the spectral widths are similar, but at 150° and 180° the spectral widths are narrower than predicted. At each angle there is a significant red-shifted component to the spectra which is closely aligned with the peak wavelength in the predicted spectra. The most significant difference between the observed and predicted spectra are their respective amplitudes. Whereas the predicted spectra are at or near thermal levels of plasma emission (4×10^3 to $5 \times 10^4 \text{ W/sr/nm}$) the measured spectra exhibit fluxes that are as much as 10^6 higher (7×10^7 to $1 \times 10^{10} \text{ W/sr/nm}$) and may be even higher since the highest fluxes recorded appear at 180° , where the measurement is only a lower limit.

Later in time, there are qualitatively different observations, as seen in figure 7.10. Now, both the predictions and the measurements show structured spectra. The source of the structure in the predictions arises from the spatial structure of the density and temperature profiles that have

evolved in the plasma, according to LASNEX predictions. The higher laser intensity increases the gain coefficient, thus enhancing any differences in plasma conditions along the laser axis. This increased laser intensity also accounts for the increased spectral width of the predictions — although the FWHM have changed very little, there is now significant flux predicted at twice the FWHM. The observed spectra have also broadened, and again there is significant flux beyond twice FWHM. Thus, since neither the predictions nor the observations appear to have gaussian-shaped spectra, it can be somewhat misleading to simply quote the FWHM of the spectra. Despite the structure of the spectra, the similarities between predictions and observations cease. Whereas earlier in time the predicted spectra were more red-shifted than the observations, now the observations are more red-shifted than those predicted. There are also more significant differences between the spectral widths at different angles. For example, the spectral width at 160° (approximately 3 nm at full-width, tenth-maximum (FWTM)) is nearly 50% narrower than the spectral width at 180° (approximately 4.5 nm FWTM). The difference in amplitudes between theory and experiment has decreased only slightly to 10^5 or less, this despite the fact that the laser intensity has increased by a factor of 5 since the early-time comparisons. This suggests that the instability is more sensitive to the change in plasma conditions than to the pump intensity. Whether this sensitivity is directly related to the plasma conditions, or that the change in conditions has permitted the onset of some other physical phenomenon which can compete with or quench SBS is not clearly understood.

Again, iso-intensity contours (figure 7.11) and spectral lineouts (figures 7.12 and 7.13) are shown, this time for a high density plasma experiment (#21062616) where the total laser energy was less than 350 J and where the

peak density is predicted to be less than 15% of critical density at the start of the interaction pulse. Though the total energy was slightly higher than in the previous example, the contrast of the pulse still exceeded 7 and the peak intensity exceeded $6 \times 10^{14} \text{ W/cm}^2$ (figure 7.5.b).

Early in time (figure 7.12) one observes that once again, the predicted lineshapes are smooth gaussians, red-shifted by approximately 1 nm and with widths of approximately 2 nm FWHM. In contrast, the observed spectra not only differ from those predicted, but also those observed in the previous example. At 150° , 160° and 180° , the amplitudes are roughly an order of magnitude higher than in the previous example. The laser intensity at this time is approximately 2.5 times higher ($1 \times 10^{14} \text{ W/cm}^2$) than at the same time in the previous example. Though the predicted peak fluxes are now as high as 10^5 W/sr/nm , they are still as much as 10^6 lower than measured (see observations at 180°). All reasonable (see below) spectral measurements at this time are clearly blue-shifted by 1-2 nm, depending on angle of observation. However, the measured spectral features differ significantly with angle of observation. At 140° , the spectrum is very broad (almost 5 nm FWTM). There is much structure that resembles harmonic generation since the spacing between peaks is very similar. At 150° the spectrum is quite narrow (4 nm FWTM, but less than 1 nm FWHM) and exhibits far less structure. The measurement at 150° appears to suffer from a bad choice of filtering of the instrument. Though the peak flux exceeds $1 \times 10^9 \text{ W/sr/nm}$ (which would have in hindsight probably led to no change in filtering from what was used) clearly the majority of the signal was well below this value. This can be concluded much easier from viewing the iso-intensity contour (figure 7.11.b). For this reason, it would be invalid to attempt to discuss the quantitative features of this spectrum. It would be equally invalid to attempt

to draw any conclusions of the validity of the prediction based upon this measurement. Finally, at 180° the spectrum is once again quite broad (almost 5 nm FWTM) but with a clearly dominant, narrow wavelength at around 526 nm.

Late in time, it is observed that the measured spectra remain blue shifted, unlike the previous example. And as expected, the predictions are slightly less red-shifted than earlier in time. Once again, the predictions now begin to exhibit slight structure, but not nearly to the extreme that the measurements do. Before proceeding, special comment is warranted for the measurement at 140° . At this particular angle of observation, SBS has been quenched in the experiment to a level below delectability, in fact some 1.3 ns before the completion of the interaction pulse. Though this is unique at this angle for this experiment, this was observed at many angles for many other experiments with different parameters, and can be said to be one of the dominant results encountered throughout this experimental campaign. The significance of this is not understood, but it is believed by the author that the key to understanding the behavior of SBS in these plasmas may lie in understanding the early "turn-off" of SBS in many of the experiments. It certainly deserves more attention and contemplation in future efforts, but is beyond the scope and time scale of this dissertation.

The recorded spectra at the remaining three angles of observation are qualitatively similar in that they all possess very high contrast structure. However, the peak amplitude at 150° is 100 times lower than at 160° and 180° . The discrepancies between measured and predicted amplitudes now is much smaller—in fact less than a factor of ten at 160° .

One notes that convective theory actually predicts higher peak fluxes at 160° than at 180° . This feature of convective theory, which is not

duplicated in homogeneous theory, is not widely appreciated since most have formed their intuition from the latter. This difference is easily understood if one considers the limits on the theory and the placement of the MATRS collectors with respect to the plasma density gradient rather than with respect to the pump wave vector ($\vec{k}_0$). Although the detector at 160° is 20° away from direct backscatter, it is nearly orthogonal to the predicted direction of the density gradient, which in these experiments is taken to be along the "phase gradient". Here, the direction of the phase gradient is that direction along which the resonant triad dephases, and its amplitude is significant since it is a measure of the range over which the triad remains resonant as it convects through the plasma — it is essentially a measure of the length of the convective amplifier. One observes from the Rosenbluth gain factor for a linear ramp profile [equation (7.2)] that the denominator depends upon the components of the daughter-wave group velocities along the phase gradient. Since the group velocity of the scattered light wave recorded by the detector located nearly normal to the phase gradient has a vanishing component along the phase gradient ($v_{1x} \equiv \vec{v}_1 \cdot \vec{K}' \rightarrow 0$), the gain factor diverges at this angle of observation. Incidentally, this is also predicted for observation angles directly opposed to the phase gradient since now the ion acoustic group velocity has a vanishing component along the phase gradient ($v_{2x} \equiv \vec{v}_2 \cdot \vec{K}' \rightarrow 0$). Experimentally, one would expect some other nonlinear effect to dominate the observations at these angles if convective theory were valid, since divergent light flux is physically invalid. Consequently, both qualitative and quantitative features of the spectra should be different at these angles of observation. This is one of many reasons why it is essential to measure the time-resolved spectra at many angles. Simple time- and spectrally-integrated energy measurements do not

in figure 7.14. From Figure 7.15 one observes that at 2.3 ns the SBS peak wavelength is narrow (≤ 1 nm FWHM) and red-shifted at all four angles of observation. The shifts vary in magnitude from approximately 1 nm (150°) to nearly 2 nm (180°). However, in addition to this narrow peak, the spectrum at each angle possesses a lower-level broad signal at 10-50% of the peak. Note that the shift of the peak wavelength does not increase monotonically with increasing angle of observation, as one would expect from simple application of wave vector- and frequency-matching conditions. This suggests that the peak scattered wavelength at each angle of observation originates from different positions in the plasma, since they are necessarily created from regions with distinctly different plasma conditions. Finally, the amplitudes measured are similar, if not slightly lower than those measured in the low intensity experiments. If the experiments were conducted in the linear convective regime, then one would expect the peak fluxes to increase roughly exponentially with increasing laser intensity. Clearly this is not the case at these intensities. A discussion of how the instability may be accessing the saturated regime is saved for the final section of this chapter. A cautionary note about the amplitude stated at 180° : the unphysically low values are a result of the analysis algorithm discussed earlier in this chapter where it was explained that all values shown for 180° are lower limits. Obviously for this case, the lower limit is unreasonably conservative, since it is six orders of magnitude lower than the amplitude measured at 160° .

Late in time, one sees again that the SBS flux has dropped below the threshold of detectability. This can most easily be seen by observing the contours shown in figure 7.14. Since the laser intensity has remained essentially constant and the plasma conditions have evolved linearly in time, one would not predict such behavior from any linear theory alone.

discriminate between features that could indicate the source of the scattered light. The fact that the spectra observed at 160° in this experiment and others is not qualitatively different, nor significantly more intense from spectra at other angles during the same experiments is just one example that suggests convective theory fails to explain the nature of SBS in these plasmas.

7.4.2 High Intensity Experiments

A separate but compatible⁽¹⁾ goal of this experimental campaign was to measure intensity-scaling of SBS in these long scale length, very underdense plasmas. The high intensity ($I_{\text{laser}} \geq 1 \times 10^{15} \text{ W/cm}^2$) experiments in this campaign partially fulfilled this goal, but failed to comply with the linear theory assumptions of weakly driven instability ($\gamma_0/\omega_{ia} \ll 1$) and weakly perturbed plasma ($v_{osc}/v_e \ll 1$), although its entertaining to note that many in the past have invoked linear convective theory to try and explain their observations in such regimes. Further, some of these experiments exceeded absolute instability thresholds at a few select locations within the plasma. Nonetheless, for completeness, these experiments will now be discussed to gain insight into the qualitative features of SBS at these higher intensities. They also indicate by their relative scaling to the low intensity experiments that SBS was saturated even in the low intensity experiments, which would likely further invalidate the linear convective theory.

The following data are shown for a low density plasma experiment (#21071911). Again, intensity contours are presented at four separate angles

⁽¹⁾Recall that the primary goal of this experimental campaign was to test the linear convective theory of SBS.

For completeness, we now consider a similar high density plasma experiment (#21071913). Again, contours for three separate angles appear in figure 7.16. No contour plot appears for SBS at 150° because there was no scattered light recorded at that angle. It is speculated that the MATRS collector at that particular angle was either misaligned, very over-attenuated by filtration or damaged during this experiment. The spectra recorded at 2.3 ns for this particular experiment are characterized by large shifts, as seen in figure 7.17. Note that at 140° and 160° the shifts are very red, exceeding 3 nm, with a spectral width of approximately 2 nm or greater. Again, the amplitudes are similar in magnitude to all other experiments that have been discussed, despite the increased laser intensity in this experiment. What is most interesting about these spectra at this time are the features observed at direct backscatter. At 180° , the signal is a double-peaked or "doublet" spectrum, with both red (3 nm) and blue (-3 nm) shifts with, again, very broad spectral widths (greater than 2 nm FWHM each). Though doublet spectra were observed in other experiments, none were as broad or as intense as this. Also, observations of doublets in other experiments showed an overall red- or blue-shift and a dominant peak; that is, the red shifted peak was either more or less shifted than the blue-shifted peak, and was either more or less intense. In this case, the two peaks are shifted approximately the same from the laser wavelength and are about equally intense.

Later in time one observes that the SBS at 140° has vanished and that the blue component of the doublet at 180° has also disappeared (figure 7.18). The remaining spectra are much narrower and much less red-shifted, consistent with the drop in both plasma density and temperature. Note, however, that the amplitudes at 160° and 180° have remained essentially

constant, despite the fact that the amplitude has clearly dropped precipitously at 140° (figure 7.16.a).

7.5 Discussion

It is quite clear from the previous section that convective theory has predicted very different spectra, both quantitatively and qualitatively, from those which were observed in the experiments. This section will focus on possible reasons for these clear differences.

However, before continuing further, it is important to discuss the validity of the assertion made earlier that these experiments successfully isolated SBS for study. As discussed briefly in chapter 3, many past experiments have been complicated by competing effects occurring simultaneously with SBS. Once such competing effect is stimulated Raman scattering (SRS), which is similar to SBS except that the longitudinally polarized decay wave is an electron plasma wave rather than an ion acoustic wave. Many have speculated that the presence of SRS can modify the behavior of SBS¹⁴⁻²¹. For most of the experiments in this campaign, there was no SRS observed. However, for a select few, very small amounts of SRS were observed to occur for very short period of time (~a few hundred picoseconds). It is unclear what impact the presence of Raman would have on these observations. However, since the observed Raman light was at very low amplitudes and lasted for very short duration, it is believed that any impact would be small.

One possible source of discrepancy is the symmetry assumed in the derivation of the theory. The theory applied here is derived for a one-dimensional linear plasma. The goal of the experiments was to produce a

plasma that was as close to 1-D as possible. LASNEX predicts that near the central 500 μ m the plasma is very linear, yet further out the plasma tends to become quite two-dimensional in nature. This has particular consequences for side- and forward-scatter. For a perfect 1-D linear plasma, daughter waves directed orthogonal to the plasma gradients will observe an essentially infinite, homogeneous plasma along their direction of travel. This tends to make flux predictions at side-and forward-scatter diverge. Clearly these laser-produced plasmas have finite, albeit large, scale lengths normal to their gradients. For this reason, comparisons near side-scatter are at best tenuous.

It is also possible that hot spots in the laser beam could partially explain the observed results. In the previous chapter, it was shown that a hot-spot model alone could not completely explain the enhanced convective scattering, although a few hot spots at over 20 times the nominal intensity could lead to a more reasonable comparison. What is certainly possible is that hot spots at a few times the nominal intensity (of which there are several in Nova beams) or greater could exceed the absolute threshold intensity for SBS at select locations within the plasma. One would expect that such behavior would be characterized by observation of narrow, "spiky" spectra, as was seen in some of the experimental results.

Another possible explanation for the large discrepancies are local fluctuations in plasma conditions. Again, the theory is derived for plasmas that can accurately be described by well-behaved or "smoothly-varying" plasma conditions and more importantly, their respective scale lengths. It is possible that short-scale turbulence exists, but remains unmeasured, in these plasmas. Local enhancement of density perturbations could significantly enhance the Thomson scattering noise contribution. More likely, local "flat-spots" in the plasma flow could cause the instability to become absolute at

those locations, since their locally diverging scale lengths no longer offer any stabilizing effect. Once the process becomes absolutely unstable at these select locations, the resulting scattered-light spectrum would be significantly altered from that resulting from a convective instability. The semi-random spectral spikes observed in some of the experiments would suggest that this may be occurring.

These local flat spots in the flow profile wouldn't necessarily have to exist in the plasma a priori. H.A. Rose and D.F. DuBois have shown³, based on analytic calculations, that even moderately intense hot spots in a typical laser beam can ponderomotively expel plasma within the hot spot. The resulting rarefaction in plasma density can result in a local flat spot in the density profile, and it is speculated that this can in turn cause flat spots in the flow profile. These local flat spots, coupled with the enhanced intensity within the hot spot of the pump could together drive the instability absolute at those locations, even if the hot spots were below threshold without profile modification. According to the reference, localized flat spots in the profile will occur if the inequality

$$\frac{I_{hotspot, W/cm^2}}{T_{e, eV}} > \frac{1.05 \times 10^{12}}{\pi \lambda_{0, \mu} L_{n, \mu}} \frac{n_c (n_c - n_0)}{(1 + F_{cm}^2) n_0^2} \sqrt{\frac{n_c - n_0}{n_c}} \quad (7.4)$$

$$\approx \frac{10^{12}}{\pi \lambda_{0, \mu} L_{n, \mu}} \left(\frac{n_c}{n_0 F_{cm}} \right)^2$$

is fulfilled. Here, all quantities are represented in the units denoted in the subscripts, and $F_{cm} = D_{cm}/2f$, where f is the f -number of the beam and D is

taken here to be the beam diameter. For typical values of $f = 4.3, D = 64\text{cm}, \lambda_0 = 0.527\text{nm}, n_0/n_c \approx 0.05, L_n \approx 1200\mu\text{m}$ and $T_e \approx 3000\text{eV}$, one finds that significant profile modification will occur for hot spots in the laser exceeding approximately $1 \times 10^{13} \text{ W/cm}^2$. Note that the nominal or average intensity for all experiments in this campaign exceed this value. Thus, based on these simple calculations, one may conclude that hot spots in the laser could significantly alter the behavior of SBS.

Once again, however, in order for this process to significantly affect SBS, the flow gradient in the vicinity of the hot spot would have to be significantly modified. For a one dimensional plasma model, conservation of particle flux dictates that any modification in the density profile would have to be accompanied by a modification in the flow profile. However, if one considers a two dimensional model of plasma modification by hot spots, the conclusions can be quite different. Two dimensional calculations recently completed by Berger²² have shown that density rarefaction from hot spots can occur without a significant modification to the flow profile. In these calculations, it was found that for a nominal $f/4$ laser cone, the ratio of the length to width of a typical hot spot is roughly 16, leading to plasma expulsion predominantly transverse to the laser wave vector. This transverse particle flux, in turn, resulted in negligible changes in the flow along the original density gradient, which was taken to be along the laser wave vector. Thus, although it is likely that the density profile can be significantly altered by the presence of hot spots in the laser beam, it is unclear how this would affect the flow gradient, which is the dominant stabilizing mechanism for convectively growing SBS.

Filamentation, the zero-frequency counterpart to SBS, could in principle enhance the vacuum laser intensity, which would in turn change

the spectra predicted by convective theory. Filamentation can occur either from the ponderomotive force caused by the laser itself, where the gradient in light pressure tends to rarefy the plasma locally and consequently modify the local index of refraction in such a way as to focus the light, or it can originate from local heating of the plasma by the beam, which again locally modifies the index of refraction in such a way as to enhance the focusing of the beam. Of course, the relative importance of these two effects depends primarily on the intensity of the beam (ponderomotive force) and the collisionality or ability of the plasma to absorb the laser light (local heating). Expressions for the convective growth rates for both ponderomotive and thermal filamentation are given by Stamper *et al.* ²³

$$\alpha_{PM} = 2.92 \times 10^{-9} \frac{(n_0/n_c) I_0 (W/cm^2) \lambda_0 (\mu m)}{(1 + 1/Z) T_0 (eV)} \quad (cm^{-1}) \quad (7.5)$$

and

$$\alpha_{TH} = \frac{13.2 \ln \Lambda}{\lambda (\mu m)} \left(\frac{Z^2 (n_0/n_c)^3 I_0 (W/cm^2)}{(1 + 1/Z) \phi(Z) [T_0 (eV)]^5} \right)^{1/2} \quad (cm^{-1}) \quad (7.6)$$

where the subscript *PM*(*TH*) indicates ponderomotive (thermal) filamentation, I_0 , λ_0 , n_0 , and T_0 are the nominal laser intensity, laser wavelength, background plasma density, and background temperature, respectively, Z is the average charge state of the ions, $\phi = (Z + 0.24)/(1 + 0.24Z)$, and $\ln \Lambda$ is the Coulomb logarithm. For the low density, low intensity experiments one finds $1/\alpha_{PM} \geq 0.1 cm$ and $1/\alpha_{TH} \geq 2.5 cm$. Since these plasmas are at most a few millimeters in size, it is clear that thermal filamentation is

inconsequential, and that there are at most only a couple of ponderomotive filamentation gain lengths in these plasmas. However, for other experiments with higher laser intensity and/or higher peak density, ponderomotive filamentation can become significant. It is important to note, however, that the above expressions are derived for homogeneous plasmas. Consequently, one would expect these calculations to be extremely conservative estimates of filamentation growth lengths. In fact, a more detailed numerical calculation²⁴ for the case of low plasma density and laser intensity indicates that the nominal laser intensity would be increased at most by a factor-of-two (less than one filamentation gain length). So, as was concluded in chapter 6, though filamentation may be enhancing the laser intensity and consequently SBS, it is likely that filamentation alone cannot explain the observations.

Yet another possibility to be considered is that the instability doesn't have enough time to approach steady state. For an instability to be convective, the resonant triad must convect out of the region in which it is resonant. The concept of convective instability is meaningless until the triad completely crosses the resonant region. If, in a time period shorter than the transit time, the instability undergoes such significant growth that it saturates, then the time-asymptotic spectrum of scattered light is dominated by the physics of instability saturation rather than linear convective growth.

To consider the affect of this assertion, one notes that the transit time for the resonant triad to traverse a given resonant length is given by $\tau \approx L_\phi / v_{g2}$, where v_{g2} and L_ϕ are the ion acoustic group velocity along the phase gradient and the phase-matching scale length, respectively. To arrive at an expression for L_ϕ one notes that the gain factor can be written as $G = K L_\phi$ where K is a gain per unit length, given by

$$K = \frac{\gamma_0}{\sqrt{v_{1g} v_{2g}}}. \quad (7.7)$$

It is then a simple matter to show that

$$L_\phi = \frac{2\pi\gamma_0}{K'\sqrt{v_{1g} v_{2g}}}. \quad (7.8)$$

The key question is whether the instability can grow to significant levels before it has time to exhibit convective behavior. This is best addressed by considering the product $\gamma_0\tau$ which is the number of e-foldings that the instability has undergone. This product varies with position within the plasma, and is given by

$$\begin{aligned} \gamma_0\tau &= \gamma_0 \frac{L_\phi}{v_{2g}} \\ &= \frac{2\pi\gamma_0^2}{K'v_{1g}^{3/2}v_{2g}^{1/2}} \\ &= G \cdot \sqrt{\frac{v_{1g}}{v_{2g}}} \end{aligned} \quad (7.9)$$

For the case of direct backscatter (which minimizes transit time), values for the lowest intensity experiment yields $2 < \gamma_0\tau < 130$ within a 2 mm region about the center of the plasma. Thus, based on this argument alone, before the instability can become convectively limited (time-asymptotic behavior) the instability can grow by as much as 100 or more e-foldings! Of course, nonlinear effects would suppress growth of the instability before this could

occur, thus dominating the spectra observed. This would be relatively simply to verify in a series of experiments that measured the early-time behavior of SBS, but since for these experiments the instability can undergo many e-foldings in a time shorter than the instrumental resolution, it cannot be completely verified here. Thus, it is conceivable that significant growth of the instability occurs in these experiments before convective growth has a chance to dominate the behavior of SBS.

7.6 Conclusion

Observations of SBS have been reported in a series of experiments that were designed to be both relevant to the ICF community a clear test of convective theory. Unlike previous experiments, these were conducted in a weakly perturbed plasma, and the instability was weakly driven, both limits to the applicability of the linear convective theory. Further, these plasmas possessed a significant and predictable flow profile, which is thought to be a significantly stabilizing effect for convectively unstable SBS. Finally, the first measurements of temporally and spectrally resolved scattered light were made at many angles simultaneously to verify the behavior of SBS.

A detailed comparison of the experimental observations with predictions of the convective theory show many differences and very few, if any, similarities. First, levels of observed scattering were as much as six orders of magnitude higher than those predicted. Second, a detailed comparison of the spectral shapes showed that neither the shifts nor the widths were predictable by the linear convective theory. Finally, structure within the observed spectra were occasionally similar to predictions, but differed significantly both in magnitude and wavelength. Many possibilities

were explored to explain these differences, and it is concluded that the linear convective theory fails to predict the observations most likely because of its inability to incorporate sufficient plasma detail, possibly including time-dependent behavior. More detailed studies of possible nonlinear effects may aid in understanding the observations, but they alone cannot resurrect the linear convective theory.

7.7 References

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	Low Intensity Laser ($\sim 4 \times 10^{14}$ W/cm ²)	High Intensity Laser ($\sim 4 \times 10^{15}$ W/cm ²)
Low Density Plasma ($n_{\text{peak}} \leq 0.05 n_c$)	21070916 21111511 21111516	21070914 21071911
High Density Plasma ($n_{\text{peak}} \leq 0.15 n_c$)	21062616 21111514 21071915	21062614 21062618 21071913 21111508

Table 7.1 Nova shot numbers corresponding to all of the Titanium Foil experiments conducted to study SBS. Numbers in bold face indicate acquisition of streaked-spectra with MATRS. Those in normal face produced no streaked-spectra, but did produce time- and wavelength-integrated SBS scattering fluences.

Figure 7.1 LASNEX-predicted axial profiles of (a) electron density (b) electron temperature (c) ion temperature and (d) flow velocity for low density plasma experiments early in time (~ 2.3 ns). Predictions are symmetric about the Z -axis. LASNEX run #211207a.

Figure 7.2 LASNEX-predicted axial profiles of (a) electron density (b) electron temperature (c) ion temperature and (d) flow velocity for low density plasma experiments late in time (~ 3.0 ns). Predictions are symmetric about the Z -axis. LASNEX run #211207a.

Figure 7.3 LASNEX-predicted axial profiles of (a) electron density (b) electron temperature (c) ion temperature and (d) flow velocity for high density plasma experiments early in time (~ 2.3 ns). Predictions are symmetric about the Z -axis. LASNEX run #211106a.

Figure 7.4 LASNEX-predicted axial profiles of (a) electron density (b) electron temperature (c) ion temperature and (d) flow velocity for high density plasma experiments late in time (~ 3.0 ns). Predictions are symmetric about the Z -axis. LASNEX run #211106a.

Figure 7.5 Interaction Laser Temporal Profiles for experiments conducted with (a) low density plasma/low laser intensity (#21070916) (b) high density plasma/low laser intensity (#21062616) (c) low density plasma/ high laser intensity (#21071911) and (d) high density plasma/high laser intensity (#21071913). Note that start times are defined in this figure to be the first non-zero value and are assumed to occur at 1.8 ns, which is indicated from streak records obtained during the experiments.

Figure 7.6 Best fit of axial dependence of NOVA beam radius, measured at the second harmonic ($2\omega_0$) of $0.527 \mu\text{m}$. Offset of “best focus” from $Z=0$ results from measurement inaccuracies, and reflects typical axial alignment tolerances on Nova.

Figure 7.7 History of directly backscattered light for a typical experiment (#2111514) recorded with a fast photodiode. Absolute levels of direct backscattered light were estimated by assuming that the early KDP reflection contained 1% of the laser energy incident on target. Comparing the ratio of areas under the KDP reflection and SBS signals above, one can arrive at an estimate of SBS scattering. Since the 1% value for KDP reflectivity was a measured lower limit, this places a lower limit on the observed levels of SBS.

Figure 7.8 Iso-intensity contour plots of streaked spectra at (a) 140° (b) 150° (c) 160° and (d) 180° for a low density/low intensity experiment (#21070916). Contours are successive factors-of-two, with the peak contour representing 90% of the peak intensity observed at that particular angle of observation. Probe wavelength was 527 nm.

Figure 7.9 A comparison of measured and predicted SBS fluxes early in time (2.3 ns) during a low density/low intensity experiment (#21070916). Comparisons are made at observation angles of (a) 140° (b) 150° (c) 160° and (d) 180°. Solid lines represent observations with values on the left axes and broken lines represent predictions with values on the right axes. Gray vertical lines represent the interaction laser wavelength. Note different normalization constants for observations and predictions, as indicated in the axes' labels.

Figure 7.10 A comparison of measured and predicted SBS fluxes late in time (3.0 ns) during a low density/low intensity experiment (#21070916). Comparisons are made at observation angles of (a) 140° (b) 160° (c) 180° and (d) 150°. Solid lines represent observations with values on the left axes and broken lines represent predictions with values on the right axes. Gray vertical lines represent the interaction laser wavelength. Note different normalization constants for observations and predictions, as indicated in the axes' labels.

Figure 7.11 Iso-intensity contour plots of streaked spectra at (a) 140° (b) 150° (c) 160° and (d) 180° for a high density/low intensity experiment (#21062616). Contours are successive factors-of-two, with the peak contour representing 90% of the peak intensity observed at that particular angle of observation. Probe wavelength was 527 nm.

Figure 7.12 A comparison of measured and predicted SBS fluxes early in time (2.3 ns) during a high density/low intensity experiment (#21062616). Comparisons are made at observation angles of (a) 140° (b) 150° (c) 160° and (d) 180°. Solid lines represent observations with values on the left axes and broken lines represent predictions with values on right axes. Gray vertical lines represent the interaction laser wavelength. Note different normalization constants for observations and predictions, as indicated in the axes' labels.

Figure 7.13 A comparison of measured and predicted SBS fluxes late in time (3.0 ns) during a high density/low intensity experiment (#21062616). Comparisons are made at observation angles of (a) 140° (b) 150° (c) 160° and (d) 180°. There was no observed spectra at 140° by this time in the experiment. Solid lines represent observations with values on the left axes and broken lines represent predictions with values on right axes. Gray vertical lines represent the interaction laser wavelength. Note different normalization constants for observations and predictions, as indicated in the axes' labels.

Figure 7.14 Iso-intensity contour plots of streaked spectra at (a) 140° (b) 150° (c) 160° and (d) 180° for a low density/high intensity experiment (#21071911). Contours are successive factors-of-two, with the peak contour representing 90% of the peak intensity observed at that particular angle of observation. Probe wavelength was 527 nm.

Figure 7.15 Measured SBS fluxes early in time (2.3 ns) during a low density/high intensity experiment (#21071911). Spectra are measured at observation angles of (a) 140° (b) 160° and (d) 180°. Gray vertical lines represent the interaction laser wavelength. No comparisons to predictions from convective theory are made since these experiments were conducted in a strongly driven regime, which violates an assumption made in the derivation of the linear theory. The unusually low values recorded in (d) are a result of the method used to calibrate the streaked spectra at that particular angle and are, based on a comparison with (a)-(c), obviously unrealistically low.

Figure 7.16 Iso-intensity contour plots of streaked spectra at (a) 140° (b) 150° (c) 160° and (d) 180° for a high density/high intensity experiment #21071913. Contours are successive factors-of-two, with the peak contour representing 90% of the peak intensity observed at that particular angle of observation. Probe wavelength was 527 nm. Note that the detector at 150° was improperly filtered, and that (d) has different axes limits.

Figure 7.17 Measured SBS fluxes early in time (2.3 ns) during a high density/high intensity experiment (#21071913). Spectra are measured at observation angles of (a) 140° (b) 160° and (c) 180°. No spectra were observed at 150°. Gray vertical lines represent the interaction laser wavelength. No comparisons to predictions from convective theory are made since these experiments were conducted in a strongly driven regime, which violates an assumption made in the derivation of the linear theory.

Figure 7.18 Measured SBS fluxes late in time (3.0 ns) during a high density/high intensity experiment (#21071913). Spectra are measured at observation angles of (a) 160° and (b) 180° . No spectra were observed at 140° or 150° at this time in the experiment. Gray vertical lines represent the interaction laser wavelength. No comparisons to predictions from convective theory are made since these experiments were conducted in a strongly driven regime, which violates an assumption made in the derivation of the linear theory.

Appendix A

A Derivation of the Coupled Wave Equations

The fundamental equations which describe stimulated Brillouin scattering can be derived from three sets of equations. The first set is comprised of the two-fluid equations of continuity and momentum

$$\frac{\partial n_j}{\partial t} + \vec{\nabla} \cdot (n_j \vec{u}) = 0 \quad (\text{A.1})$$

$$n_j \left(\frac{\partial \vec{u}_j}{\partial t} + (\vec{u}_j \cdot \vec{\nabla}) \vec{u}_j \right) = \frac{n_j q_j}{m_j} \left(\vec{E} + \frac{\vec{u}_j \times \vec{B}}{c} \right) - \frac{1}{m_j} \vec{\nabla} p_j \quad (\text{A.2})$$

where n_j , m_j , q_j , $\vec{u}_j$ and p_j are the density, mass, charge, velocity and pressure of particles of species j , and $\vec{E}$, $\vec{B}$, and c are the vacuum electric and magnetic fields, and the speed of light in vacuum, respectively. They describe the dynamics of charged particles. Maxwell's equations (in vacuum, cgs units)

$$\vec{\nabla} \cdot \vec{E} = 4\pi\rho \quad (\text{A.3})$$

$$\vec{\nabla} \cdot \vec{B} = 0 \quad (\text{A.4})$$

$$\vec{\nabla} \times \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t} \quad (\text{A.5})$$

$$\vec{\nabla} \times \vec{B} = \frac{4\pi}{c} \vec{J} + \frac{1}{c} \frac{\partial \vec{E}}{\partial t} \quad (\text{A.6})$$

where

$$\rho = \sum_j n_j q_j \quad (\text{A.7})$$

and

$$\vec{J} = \sum_j n_j q_j \vec{u}_j \quad (\text{A.8})$$

are the charge and current densities, respectively, describe the dynamics of electromagnetic radiation in vacuum. Finally, the isothermal and adiabatic equations of state

$$p_j = n_j \theta_j \quad (\text{A.9})$$

and

$$\frac{p_j}{n_j^3} = \text{constant}, \quad (\text{A.10})$$

respectively, where θ_j is the temperature of species j (in energy units), are constitutive equations which relate the pressure and density of species j .

We assume that the plasma is initially uniform. For the purposes of this derivation, it is convenient to write the electric and magnetic fields in terms of the electrostatic potential, ϕ , and the vector potential, $\vec{A}$,

$$\vec{E} = -\frac{1}{c} \frac{\partial \vec{A}}{\partial t} - \nabla \phi \quad (\text{A.11})$$

$$\vec{B} = \nabla \times \vec{A} \quad (\text{A.12})$$

Substituting equations (A.11) and (A.12) into equation (A.6), and invoking the Coulomb gage, ($\nabla \cdot \vec{A} = 0$), we find

$$\left(\frac{1}{c^2} \frac{\partial^2}{\partial t^2} - \nabla^2 \right) \vec{A} = \frac{4\pi}{c} \vec{J}_t \quad (\text{A.13})$$

where $\vec{J}_t$ is the transverse component of the current density, defined by $\vec{\nabla} \cdot \vec{J}_t = 0$. The transverse current density can be written as

$$\vec{J}_t = -n_e e \vec{u}_t = -n_e e \vec{v}_{osc} \quad (\text{A.14})$$

where the ions are neglected since they respond very weakly to the electric field oscillations due to their much greater inertia. In the nonrelativistic limit, $(\vec{v}_{osc}/c) \ll 1$, $\vec{v}_{osc} = e\vec{E}/m_e\omega = e\vec{A}/m_e c$ since, from the equation of motion for an electron in an external electric field,

$$\frac{\partial \bar{u}_t}{\partial t} = -\frac{e}{m_e} \bar{E}_t = \frac{e}{m_e c} \frac{\partial \bar{A}}{\partial t} \quad (\text{A.15})$$

Substituting for the current density in equation (A.13) leads to

$$\left(\frac{\partial^2}{\partial t^2} - c^2 \nabla^2 \right) \bar{A} = -\frac{4\pi e^2}{m_e} n_e \bar{A} \quad (\text{A.16})$$

which is a wave equation which describes the propagation of a light wave in a plasma. To study this effect to first order, we linearize by taking $\bar{A} = \bar{A}_{laser} + \tilde{\bar{A}}$ and $n_e = n_{0e} + \tilde{n}_e$, where the quantities with tildes ($\sim$) denote small perturbations (first order terms) on the background (zeroth order terms) quantities. Substituting, and retaining terms only to first order, one arrives at the expression describing the scattering of a light wave in a plasma by a small density perturbation

$$\left(\frac{\partial^2}{\partial t^2} - c^2 \nabla^2 + \omega_{pe}^2 \right) \tilde{\bar{A}} = -\frac{4\pi e^2}{m_e} \tilde{n}_e \bar{A}_{laser} \quad (\text{A.17})$$

To derive an expression for the low frequency ion density fluctuations (ion acoustic waves), we return to the two-fluid equations. First, consider the equations for the ions. Since the ion disturbance propagates faster than the ion heat flow, $\omega/k \gg v_i \propto Q_i$, it is appropriate to apply the adiabatic equation of state. In a similar fashion to that above, we may linearize the ion fluid equations by taking $n_i = n_{0i} + \tilde{n}_i$, $\bar{u}_i = \tilde{\bar{u}}_i$, $\phi = \tilde{\phi}$, and $p_i = p_{0i} + \tilde{p}_i = p_{0i} + 3\theta_i \tilde{n}_i$. This last expression can readily be verified by substituting into the adiabatic equation of state

$$\frac{p_{0i}}{n_{0i}^3} = \frac{p_{0i} + \tilde{p}_i}{(n_{0i} + \tilde{n}_i)^3} \quad (\text{A.18})$$

and noting that at very early or late times, the ions will also be isothermal so that $p_{0i} = n_{0i}\theta_i$. Substituting all of these expressions into the ion fluid equations gives

$$\frac{\partial \tilde{n}_i}{\partial t} + n_{0i} \tilde{\nabla} \cdot \tilde{\mathbf{u}}_i = 0 \quad (\text{A.19})$$

$$\frac{\partial \tilde{\mathbf{u}}_i}{\partial t} = -\frac{Ze}{M} \tilde{\nabla} \tilde{\phi} - \frac{3\theta_i}{M} \tilde{\nabla} \tilde{n}_i \quad (\text{A.20})$$

To combine these two coupled partial differential equations, we take the time derivative of equation (A.19) and the divergence of equation (A.20) and eliminate the common term $\partial(\tilde{\nabla} \cdot \tilde{\mathbf{u}}_i)/\partial t$ to arrive at

$$\frac{\partial^2 \tilde{n}_i}{\partial t^2} - \frac{n_{0i}Ze}{M} \nabla^2 \tilde{\phi} - \frac{3\theta_i}{M} \nabla^2 \tilde{n}_i = 0 \quad (\text{A.21})$$

To complete the analysis, an expression for $\tilde{\phi}$ in terms of plasma parameters must be found. Consider the electron fluid equations, where we now note that the electron fluid is best described by the isothermal equation of state since the electrons can equilibrate faster than the speed of the ion wave, $\omega/k \ll v_e \propto Q_e$. After separating the velocity into transverse and longitudinal components, and substituting the isothermal equation of state for the pressure gradient, the longitudinal component of the momentum equation becomes

$$\frac{\partial \bar{u}_l}{\partial t} = \frac{e}{m_e} \bar{\nabla} \phi - \frac{1}{2} \bar{\nabla} \left(\bar{u}_l + \frac{e \bar{A}}{m_e c} \right)^2 - \frac{\bar{\nabla} (n_e \theta_e)}{n_e m_e} \quad (\text{A.22})$$

where we have used the identity $\bar{\nabla}(\bar{u})^2 = 2(\bar{u} \cdot \bar{\nabla})\bar{u}$. The electron inertia may be neglected ($m_e \rightarrow 0$) since we are studying a low frequency ion wave. Then, linearizing equation (A.22) in a manner identical to equations (A.16), (A.18)-(A.20)

$$\bar{\nabla} \bar{\phi} = \frac{e}{m_e c^2} \bar{\nabla} (\bar{A}_L \cdot \bar{\bar{A}}) + \frac{\theta_e}{en_{0e}} \bar{\nabla} \bar{n}_e \quad (\text{A.23})$$

Finally, substituting equation (A.23) into equation (A.21) yields an expression for an ion acoustic wave driven by the ponderomotive force evolving from the interaction between the laser and scattered light waves

$$\frac{\partial^2 \bar{n}_i}{\partial t^2} - c_s^2 \nabla^2 \bar{n}_i = \frac{n_{0i} Z e^2}{m_e M c^2} \nabla^2 (\bar{A}_L \cdot \bar{\bar{A}}) \quad (\text{A.24})$$

where we have noted that $Z n_{0i} = n_{0e}$, $\bar{n}_e / n_{0e} \approx \bar{n}_i / n_{0i}$, and where we have used the definition of sound speed

$$c_s = \sqrt{\frac{Z \theta_e + 3 \theta_i}{M}} \quad (\text{A.25})$$

Appendix B

An IDL Algorithm for Visualization and Quantitative Analysis of MATRS Data

In order for the data that is recorded on film by MATRS to be of any practical use to the experimentalist, it must be accurately transformed into numerical information. This is accomplished by digitizing the photographs with a photodensitometer, as was described in chapter 5. Once the data is converted and stored digitally in this fashion, it is then possible to manipulate and extract quantitative information useful for comparison to theory, for example.

A series of algorithms, written in the Interactive Data Language (IDL) programming language, were produced for just this purpose. IDL is a FORTRAN-like programming language which operates on many different computer platforms, and in this case it was executed on a VAX mainframe computer interactively with various DEC workstations and Ramtek graphics display terminals. Hardcopy output of the data can be obtained either in

black and white or color, and the data files created with the algorithms can also be transferred to a Macintosh computer for further formatting and printing. With the exception of the contour plots in this dissertation, which were produced directly in IDL, the remainder of data figures were produced with KaleidaGraph on the Macintosh from IDL-created data files that were transferred.

As an example of the method in which the MATRS data was transformed, an example of the most fundamental algorithm used for visualization and quantitative analysis is given here. No effort is made to explain the algorithm in a line-by-line or even global fashion, but the language is sufficiently similar to FORTRAN that anyone proficient with FORTRAN will easily be able to follow the format of the algorithm.

The code now follows as it is executed, on the next page, where the only change made is the font in which it appears, to match the format of this dissertation.


```

;70916PROCESS2.PRO
;THIS ROUTINE TAKES THE STREAK RECORD,SEPARATES THE VARIOUS MATRS CHANNELS
;INTO SEPARATE SUBIMAGES, CREATES CONTOUR PLOTS, ABSOLUTELY CALIBRATES THEM,
;TIME AND WAVELENGTH CALIBRATES THEM, AND PREPARES THEM FOR LINEOUT ANALYSIS
;
;THIS VERSION RUNS ON IDL V.2
;
;WRITTEN 14 AUGUST,1991
;MODIFIED 10 SEPTEMBER 1991
;BY KEITH S. BRADLEY

!FANCY =0
DATE=STRMID(!STIME,0,11)

;!XMIN=510
;!XMAX=550
;!YMIN=-1
;!YMAX=14

!XMIN=0 & !XMAX=400 & !YMIN=0 & !YMAX=400

IMG99 = GETINT('21070916SOS4_BLG.INT')
COMMON INTINFO,DIR,EXT,INS,IPPS,X99,Y99,HEAD,TITLE,XLAB,YLAB,ZLAB
!MTITLE = '21070916, MATRS INTO SOS-IV; '+DATE
!XTITLE='PIXELS'
!YTITLE='PIXELS'

;TVSCL,ALOG(IMG99)
PRINT,'THE MAXIMUM VALUE OF THE IMG99 IS', MAX(IMG99)

DLAM1=-13.147
DLAM2=-9.462
DLAM3=-6.574

;THESE ARE THE WAVELENGTH OFFSETS
;DETERMINED FROM THE FIDU-CALIBRATION
;SHOT #21072900. THESE ARE REFERENCED

```

```

DLAM4=-2.690      ;3.287
DLAM5=0
DLAM6=3.880      ;3.187
DLAM7=6.670      ;6.275
DLAM8=9.362

DIME1=0
DIME2=-3.864
DIME3=-0.336
DIME4=6.671      ;2.898
DIME5=-0.065
DIME6=7.453      ;2.940
DIME7=4.042      ;0.252
DIME8=-3.108

FIX1=0.711
FIX2=0.643
FIX3=0.690
FIX4=0.669
FIX5=0.762
FIX6=0.766
FIX7=0.770
BL6_PIGTAIL=1.046
;D3_ADVANCE =

BL6_ADVANCE = 3.129

; TO CHANNEL 5 SINCE FOR THIS SHOT DAY,
; THE BL6 BACKSCATTER WAS CONNECTED TO
; CHANNEL 5 AND THE STOCK REFLECTION FOR
; THIS CHANNEL WAS USED WITH SOS_PROCESS
; TO DEFINE 527NM AND 0.0NS

; THESE ARE THE TIME OFFSETS REFERRED TO
; CHANNEL 5, FOR THE SAME REASONS THAT
; ARE LISTED ABOVE. TO SEE HOW THESE
; OFFSETS ARE DETERMINED, REFER TO THE
; ROUTINES 72900CON.PRO AND FIDU_LOC.PRO
; LOCATED IN THE [.IDL.29JULY91] SUBDIR.

; THE FILTER FIXTURES USED WITH THE NEW
; INPUT SLIT FOR SOS-IV EACH INTRODUCE A
; SHORT TIME DELAY FOR THEIR RESPECTIVE
; CHANNELS. I DID NOT RECORD WHICH FIXTURE
; WAS CONNECTED TO WHICH CHANNEL ON THIS
; SHOT DAY, BUT THIS ONLY INTRODUCES AN
; ERROR OF +/- 65PS, WHICH IS SMALL COMPARED
; TO THE ERROR INTRODUCED WHEN PROCESSING
; THIS FILM CHANNEL D3 HAS A DIFFERENT OPTIC,
; CAUSING THE LIGHT IN THIS CHANNEL TO
; ARRIVE AT THE SOS-IV PLATFORM EARLIER THAN
; THE OTHER CHANNELS

; ON THIS SHOT DAY, A NEW, SHORTER FIBER WAS
; USED ON THE BL6 DIRECT BACKSCATTER
; CHANNEL TO FURTHER SPREAD OUT THE DATA
; ON THE PAGE. SINCE BL6 WAS USED WITH
; SOS_PROCESS TO ESTABLISH WAVELENGTH AND

```

```

; TIME FIDUCIALS, ALL OTHER CHANNELS MUST
; BE ADJUSTED FOR THIS TIME ADVANCE.  THUS,
; 3.129 NS MUST BE ADDED TO ALL OTHER
; CHANNELS.

```

```

NULL_PIXELS= WHERE(IMG99 LE 1.1*MIN(IMG99))      ; THIS SETS BACKGROUND
IMG99(NULL_PIXELS)=0                             ; TO ZERO

```

```

PEAK99 = 0.9*MAX(IMG99)
LEV99 = [PEAK99/512,PEAK99/256,PEAK99/128,PEAK99/64,PEAK99/32,PEAK99/16$
,PEAK99/8,PEAK99/4,PEAK99/2,PEAK99]
;FOR ENTIRE IMAGE

```

```

THICK10=[3]
IX.THICK=3
IY.THICK=3
CONTOUR,SMOOTH(IMG99,3),LEVELS=LEV99,C_THICK=THICK10

```

```

!XMIN=524
!XMAX=532
!YMIN=-0.5
!YMAX=4
; !XMIN=0 &!XMAX=0 & !YMIN=0 & !YMAX=0

```

```

;IMGD1 = IMG99()
;XD1 = X99()- DLAM1
;YD1 = Y99()- DTIME1 + BL6_ADVANCE
; DETECTOR D1 @ 18.5 DEG FROM FORWARD

```

```

;IMGD2 = IMG99(204:313,132:221)
;XD2 = X99(204:313) + DLAM7
;YD2 = Y99(132:221) + DTIME7 + BL6_ADVANCE
;LEVD2 = [44,22,11,5.5,2.75]
; DETECTOR D2 @ 38.3 DEG FROM FORWARD

```

```

;IMGD3 = IMG99()
;XD3 = X99() - DLAM3
;YD3 = Y99() - DTIME3 - D3_ADVANCE + BL6_ADVANCE
; DETECTOR D3 @ 54.4 DEG FROM FORWARD

IMGD4 = IMG99(84:161,187:295)
; TVSCL,ALOG(IMGD4+.0000000001)
XD4 = X99(84:161) - DLAM4
YD4 = Y99(187:295) - DTIME4 ;+ BL6_ADVANCE
; DETECTOR D4 @ 23.6 DEG FROM BACK

IMGD4(0:17,0:33)=0
IMGD4(70*:0:*)=0
; ERASING CROSSTALK
; ERASING CROSSTALK

SHEAR,IMGD4,XD4,YD4,IMGD4SHEARED
; CORRECTS FOR SHEAR IN IMAGE

IMTTITLE = '21070916, 160 DEGREES FROM FORWARD;'+DATE
IXTTITLE = 'WAVELENGTH (NM)'
IYTTITLE = 'TIME (NSEC)'
PEAKD4 = 0.9*MAX(IMGD4)
LEVVD4 = [PEAKD4/128,PEAKD4/64,PEAKD4/32,PEAKD4/16,PEAKD4/8,PEAKD4/4$
,PEAKD4/2,PEAKD4]
CONTOUR,SMOOTH(IMGD4SHEARED,3),XD4,YD4,LEVELS=LEVVD4,C_THICK=THICK10

; NEXT LINES ABSOLUTELY CALIBRATE THE DATA FROM THE EBM DATA
;
;
EBMD4=1.9

ETOTALD4=TOTAL(IMGD4)
EXTALKD4=TOTAL(IMGD4(0:17,0:31))

; EBM#125 HAD LOW COUNTS, SO THIS
; VALUE, IN J/SR, IS A FITTED VALUE
; FROM ADJACENT EBM'S
; TOTAL ENERGY OF SUB-IMAGE
; ERRONEOUS CONTRIBUTIONS
; FROM ADJACENT MATRS

```

```

ESIGNALD4=ETOTALD4-EXTALKD4
; CHANNELS
; SUBTRACT OUT ERRONEOUS
; ENERGY CONTRIBUTION

PRINT,'THE D4 SIGNAL IS',ESIGNALD4,'OR',ESIGNALD4*100/ETOTALD4,$
'PERCENT OF THE TOTAL ENERGY FOR THAT CHANNEL'
IMGD4_ABS=IMGD4SHEARED*EBMD4*1E9/ESIGNALD4/(XD4(70)-XD4(69))$
/(YD4(70)-YD4(69))

IMGD5 = IMG99(182:237,119:231)
;
; TVERAS
; TVSCL,ALOG(IMGD5+0.0000000001)
XD5 = X99(182:237) - DLAM7
YD5 = Y99(119:231) - DTIME7  ;+ BL6_ADVANCE

IMGD5(34:*,0:57)=0
IMGD5(47:*,0:*)=0
IMGD5(0:12,62:*)=0

; REMOVING CROSSTALK

SHEAR,IMGD5,XD5,YD5,IMGD5SHEARED

IMTITLE = '21070916, 150 DEGREES FROM FORWARD; '+DATE
PEAKD5 = 0.9*MAX(IMGD5)
LEVD5 = [PEAKD5/128,PEAKD5/64,PEAKD5/32,PEAKD5/16,PEAKD5/8,PEAKD5/4,$
PEAKD5/2,PEAKD5]
; CONTOURXY,SMOOTH(IMGD5,3),XD5,YD5,LEVD5
CONTOUR,SMOOTH(IMGD5SHEARED,3),XD5,YD5,LEVELS=LEVD5,C_THICK=THICK10

; NEXT LINES ABSOLUTELY CALIBRATE THE DATA FROM THE EBM DATA
; SEE D4 PROCESSING ABOVE FOR EXPLANATION OF ALGORITHM

```

```

EBMD5=0.72          ; EBM#104 HAD A READING OF 0.72 J/SR
ETOTALD5=TOTAL(IMGD5)
EXTALKD5=TOTAL(IMGD5(0:14,82:*)))+TOTAL(IMGD5(40:*,0:37))
ESIGNALD5=ETOTALD5-EXTALKD5
PRINT,'THE D5 SIGNAL IS',ESIGNALD5,'OR',ESIGNALD5*100/ETOTALD5,$
'PERCENT OF THE TOTAL ENERGY FOR THAT CHANNEL'
IMGD5_ABS=IMGD5SHEARED*EBMD5*1E9/ESIGNALD5/(Y99(70)-X99(69))/(Y99(70)-Y99(69))

```

```

IMGD6 = IMG99(156:223,198:292)          ; DETECTOR D6 @ 48 DEG FROM BACK
;TVERAS
;TVSCL,ALOG(IMGD6)
XD6 = X99(156:223) - DLAM6
YD6 = Y99(198:292) - DTIME6  ;+ BL6_ADVANCE

IMGD6(41:*,0:23)=0
IMGD6(0:11,37:*)=0          ; REMOVING CROSSTALK

```

```

SHEAR,IMGD6,XD6,YD6,IMGD6SHEARED

```

```

IMTITLE = '21070916, 140 DEGREES FROM FORWARD, '+DATE
PEAKD6 = 0.9*MAX(IMGD6)
LEVVD6 = [PEAKD6/64,PEAKD6/32,PEAKD6/16,PEAKD6/8,PEAKD6/4,PEAKD6/2,$
PEAKD6]
; CONTOURXY,SMOOTH(IMGD6,3),XD6,YD6,LEVVD6
CONTOUR,SMOOTH(IMGD6SHEARED,3),XD6,YD6,LEVELS=LEVVD6,C_THICK=THICK10
; NEXT LINES ABSOLUTELY CALIBRATE THE DATA FROM THE EBM DATA
;
;

```

```

EBMD6=0.27          ; EBM#102 HAD A READING OF 0.27 J/SR
ETOTALD6=TOTAL(IMGD6)
EXTALKD6=TOTAL(IMGD6(43:*:0:22))+TOTAL(IMGD6(0:15,42:71))
ESIGNALD6=ETOTALD6-EXTALKD6
PRINT,'THE D6 SIGNAL IS',ESIGNALD6,'OR',ESIGNALD6*100/ETOTALD6,$
'PERCENT OF THE TOTAL ENERGY FOR THAT CHANNEL'
IMGD6_ABS=IMGD6SHEARED*EBMD6*1E9/ESIGNALD6/(X99(70)-X99(69))/(Y99(70)-Y99(69))

IMGBL6 = IMG99(108:181,21:135)          ; DETECTOR BL6 AT DIRECT BACKSCATTER
;TVERAS
;TVSCL,ALOG(IMGBL6)
XBL6 = X99(108:181)-DLAM5
YBL6 = Y99(21:135)-DTIME5      ;+ BL6_ADVANCE

IMGBL6(66:*:*)=0          ; REMOVING CROSSTALK

SHEAR,IMGBL6,XBL6,YBL6,IMGBL6SHEARED

IMTITLE = '21070916, DIRECT BACKSCATTER,'+DATE
PEAKBL6 = 0.9*MAX(IMGBL6)
LEVBL6 = [PEAKBL6/256,PEAKBL6/128,PEAKBL6/64,PEAKBL6/32,$
PEAKBL6/16,PEAKBL6/8,PEAKBL6/4,PEAKBL6/2,PEAKBL6]
CONTOURXY,SMOOTH(IMGBL6,3),XBL6,YBL6,LEVBL6
CONTOUR,SMOOTH(IMGBL6SHEARED,3),XBL6,YBL6,LEVELS=LEVBL6,C_THICK=THICK10
;      NEXT LINES ABSOLUTELY CALIBRATE THE DATA FROM THE EBM DATA
;
;
EBMBL6=5.73          ; SINCE THE BACKSCATTER CALORIMETER
; SEEMED TO BE OPERATING ERRATICALLY, THIS
; NUMBER IS A MINIMUM VALUE DETERMINED

```

; FROM THE DIODE AND AN ASSUMPTION THAT
 ; THE KDP REFLECTION WAS AT LEAST 1% OF THE
 ; INCIDENT BEAM ENERGY. THIS IS THEREFORE A
 ; LOWER LIMIT.

ETOTALBL6=TOTAL(IMGBL6)
 EXTALKBL6=0
 ESIGNALBL6=ETOTALBL6-EXTALKBL6
 PRINT,'THE BL6 SIGNAL IS',ESIGNALBL6,'OR',ESIGNALBL6*100/ETOTALBL6,\$
 'PERCENT OF THE TOTAL ENERGY FOR THAT CHANNEL'
 IMGBL6_ABS=IMGBL6SHEARED*EBMBL6*1E9/ESIGNALBL6/(X99(70)-X99(69))/(Y99(70)-Y99(69))

END

Appendix C

A Microsoft Excel™ Spreadsheet Algorithm for Predicting SBS Temporally-Resolved Spectra

The following list is a printout of the expressions used to calculate predicted SBS lineshapes as they appear in chapter 6 (spectral widths and peak flux evolutions) and chapter 7 (spectral shapes at multiple times).

This spreadsheet calculation was performed on a Macintosh computer using Microsoft Excel™. One of the features of Excel is that it can print, in the order they are executed, the expressions that appear in each cell of the spreadsheet. That capability was exploited to produce this list. Though this is a useful feature of the software, the author would like to note that the feature would be more useful if it allowed the user to format the list. The inability to format this list is the reason why the list appears in such elongated fashion, and the author apologizes for the inconvenience this

causes. Nonetheless, the author felt it was important that this information appear in this Appendix for the reader.

Finally, it should be noted that this spreadsheet was written primarily by Steven H. Batha, with significant input and modification by the author for the specific purposes described earlier in this dissertation.

The list of expressions follows on the next page. The first list is the subroutine used to calculate the scattered light frequency from the wave vector and frequency matching conditions at a particular position in the plasma for a given set of plasma conditions. This subroutine is called by the main spreadsheet in column O. The listing of the main spreadsheet follows immediately thereafter.

Cell: A97
Formula:=ARGUMENT("Wlaser")

Cell: A98
Formula:=ARGUMENT("Wplasma")

Cell: A99
Formula:=ARGUMENT("Klaser")

Cell: A100
Formula:=ARGUMENT("Debye")

Cell: A101
Formula:=ARGUMENT("Scatt_Angle")

Cell: A102
Formula:=ARGUMENT("Attack_Angle")

Cell: A103
Formula:=ARGUMENT("Flow_Angle")

Cell: A104
Formula:=ARGUMENT("U")

Cell: A105
Formula:=ARGUMENT("Detector_Theta")

Cell: A106
Formula:=ARGUMENT("Telectron")

Cell: A107
Formula:=ARGUMENT("Tion")

Cell: A108
Formula:

Cell: A109
Formula:

Cell: A110
Formula:

Cell: A111
Formula:=SET.VALUE(DegToRad,PI()/180)

Cell: A112

Formula: =SET.VALUE(Wsmax,1.05*Wlaser)

Cell: A113

Formula: =SET.VALUE(Wsmin,0.95*Wlaser)

Cell: A114

Formula: =SET.VALUE(Wstest,0.5*(Wsmax+Wsmin))

Cell: A115

Formula:

Cell: A116

Formula: =SET.VALUE(E_Difference,Wsmax-Wsmin)

Cell: A117

Formula: =SET.VALUE(E_Err_Lim,0.0000000001*Wmax)

Cell: A118

Formula: =SET.VALUE(E_Counter,0)

Cell: A119

Formula: =SET.VALUE(Phi,0)

Cell: A120

**Formula: =WHILE(AND(ABS(E_Difference)>E_Err_Lim,
E_Counter<E_Count_Lim))**

Cell: A121

Formula:

Cell: A122

Formula: =SET.VALUE(Ksbs,SQRT(Wstest^2-Wplasma^2)/SOL)

Cell: A123

**Formula: =SET.VALUE(Kia,SQRT(Ksbs^2+Klaser^2-2*Klaser*
Ksbs*COS(Scatt_Angle*DegToRad)))**

Cell: A124

**Formula: =SET.VALUE(Csound,SQRT(((3*Tion/Mass_ion)+(Zbar*
Telectron/(Mass_ion*(1+Kia^2*Debye^2))))))**

Cell: A125

**Formula: =SET.VALUE(Temp_Angle,(-1*(ABS(Klaser)*COS(
Attack_Angle*DegToRad)+ABS(Ksbs)*COS(**

Detector_Theta*DegToRad))/ABS(Kia)))

Cell: A126

Formula:=SET.VALUE(Kia_Angle_X,IF(OR(Temp_Angle=-1,Temp_Angle=1),180,ACOS(Temp_Angle)))

Cell: A127

Formula:

Cell: A128

Formula:=SET.VALUE(Mstar,(U/Csound)*COS(Kia_Angle_X*DegToRad))

Cell: A129

Formula:

Cell: A130

Formula:=SET.VALUE(bracket,Mstar+1/SQRT(1+Kia^2*Debye^2))

Cell: A131

Formula:=SET.VALUE(Wsnew,Wlaser-Csound*Kia*bracket)

Cell: A132

Formula:=SET.VALUE(E_Difference,Wsnew-Wstest)

Cell: A133

Formula:=IF(E_Difference<0,SET.VALUE(Wsmax,Wstest),SET.VALUE(Wsmin,Wstest))

Cell: A134

Formula:=SET.VALUE(Wstest,0.5*(Wsmax+Wsmin))

Cell: A135

Formula:=SET.VALUE(E_Counter,E_Counter+1)

Cell: A136

Formula:

Cell: A137

Formula:=NEXT()

Cell: A138

Formula:=IF(E_Counter=E_Count_Lim,SET.VALUE(Wstest,0))

Cell: A139

Formula:

Cell: A140

Formula: =RETURN(Wstest)

Cell: B2
Formula:180
Value: 180
Names: Scat_Ang

Cell: B3
Formula:50
Value: 50
Names: Attack_Ang

Cell: B4
Formula:0
Value: 0
Names: Polar

Cell: B5
Formula:=ABS((\$H\$22-\$H\$2)/(\$D\$22-\$D\$2))
Value: 710000000
Names: Grad_u, Velocity_Gradient

Cell: B6
Formula:=MAX(G2:G82)
Value: 0.054798
Names: Peak_Density

Cell: B7
Formula:=IF(MAX(Intensity)<1000000000000000,"LOW",
"HIGH")
Value: "LOW"
Names:

Cell: B8
Formula:10000000
Value: 10000000
Names: L_X

Cell: B9
Formula:100000
Value: 100000
Names: L_Y

Cell: B10
Formula:0.014
Value: 0.014
Names: Step_Size

Cell: B11

Formula:=56400*SQRT(Peak_Density*Nc)

Value: 8.30897E+14

Names: Wpeak

Cell: B12

Formula:527

Value: 527

Names: Wave_0

Cell: B13

Formula:=1.1E+21/(\$B\$12*0.001)^2

Value: 3.9607E+21

Names: Nc

Cell: B14

Formula:=2*PI()*SOL/(\$B\$12*0.0000001)

Value: 3.57429E+15

Names: Wo

Cell: B15

Formula:22

Value: 22

Names: Z

Cell: B16

Formula:20.5

Value: 20.5

Names: Zbar

Cell: B17

Formula:Titanium

Value: "Titanium"

Names:

Cell: B18

Formula:0.0002

Value: 0.0002

Names: Thick

Cell: B19

Formula:29979245800

Value: 29979245800

Names: SOL

Cell: B20

Formula:7

Value:7

Names:

Cell: B21

Formula:47.9

Value:47.9

Names: Mass_bar

Cell: B22

Formula:=SOL^2*Mass_bar*511*1836

Value:5.0002E-14

Names: Mass_ion

Cell: B23

Formula:=PI()/180

Value:0.017453293

Names: DegToRad

Cell: B24

Formula:=(Z/Mass_bar)/1836

Value:0.000250158

Names: Delta

Cell: B25

Formula:200

Value:200

Names: Spot_Diameter

Cell: B26

Formula:=0.25*PI()*Spot_Diameter^2

Value:31415.92654

Names: Area

Cell: B27

Formula:0

Value:0

Names: Detector_Phi

Cell: B28

Formula:=IF(Scat_Ang<20,D1_1,IF(Scat_Ang<50,D2_2,IF(Scat_Ang<70,D3_3,IF(Scat_Ang<145,D6__140,IF(Scat_Ang<155,D5__150,IF(Scat_Ang<165,D4_160,BL6__180))))))

Value: 50
Names: Detector_Theta

Cell: B29
Formula:
Value:
Names:

Cell: B30
Formula: 143.04
Value: 143.04
Names: D1_1

Cell: B31
Formula: 149.65
Value: 149.65
Names: D2_2

Cell: B32
Formula: 153.65
Value: 153.65
Names: D3_3

Cell: B33
Formula: 68.27
Value: 68.27
Names: D4_160

Cell: B34
Formula: 36.92
Value: 36.92
Names: D5__150

Cell: B35
Formula: 61.93
Value: 61.93
Names: D6__140

Cell: B36
Formula: 50
Value: 50
Names: BL6__180

Cell: D2
Formula: -0.1

Value: -0.1
Names: Zeta

Cell: E2
Formula: =Zeta*COS(Attack_Ang*DegToRad)
Value: -0.064278761
Names: X

Cell: F2
Formula: =Zeta*SIN(Attack_Ang*DegToRad)
Value: -0.076604444
Names: Y

Cell: G2
Formula: 0.011688
Value: 0.011688
Names: Density, N_Nc

Cell: H2
Formula: -71000000
Value: -71000000
Names: VelocityZ

Cell: I2
Formula: =VelocityZ/Cs
Value: -3.588832791
Names: Mach

Cell: J2
Formula: 0
Value: 0
Names: Theta_UZ

Cell: K2
Formula: =\$J\$2
Value: 0
Names: Theta_GradU

Cell: L2
Formula: 32409936297882.6
Value: 3.24099E+13
Names: Intensity

Cell: M2
Formula: =Wo*SQRT(Density)

Value: 3.8642E+14

Names: Wpe

Cell: N2

Formula: =SQRT(Wo^2-Wpe^2)/SOL

Value: 118526.7258

Names: Ko

Cell: O2

Formula: =TWSolver4!Solve.SBS.Den(Wo,Wpe,Ko,Debye,
Scat_Ang,Attack_Ang,Theta_UZ,VelocityZ,
Detector_Theta,Te,Ti,Mass_ion)

Value: 3.59115E+15

Names: Ws

Cell: P2

Formula: =Wo-Ws

Value: -1.68575E+13

Names: Wia

Cell: Q2

Formula: =SQRT(Ws^2-Wpe^2)/SOL

Value: 119092.3313

Names: Ks

Cell: R2

Formula: =SQRT(Ko^2+Ks^2-2*Ko*Ks*COS(Scat_Ang*DegToRad)
)

Value: 237619.0571

Names: Kia

Cell: S2

Formula: =Cs*Kia/SQRT(P)

Value: 3.02738E+12

Names: Wia_No_Shift

Cell: T2

Formula: =Kia/Ko

Value: 2.004771966

Names:

Cell: U2

Formula: =Ks/Ko

Value: 1.004771966

Names:

Cell: V2

Formula: $=10000000*2*PI()*SOL/Ws$

Value: 524.5261654

Names: Wavelength

Cell: W2

Formula: $=10*(Wavelength-\$B\$12)$

Value: -24.73834593

Names: Shift

Cell: X2

Formula: $=X3$

Value: 0

Names: Delta_Wavelength

Cell: Y2

Formula: $=90+(Attack_Ang-90+ASIN(Ks*SIN(Scat_Ang*DegToRad)/Kia)/DegToRad)$

Value: 50

Names: Kia_Angle

Cell: Z2

Formula: $=IF(AND(X=0,Y=0),1,IF(X>0,-1/SQRT(1+(Y/X)^2*(L_X/L_Y)^4),1/SQRT(1+(Y/X)^2*(L_X/L_Y)^4)))$

Value: 8.391E-05

Names: CosPhiX

Cell: AA2

Formula: $=ACOS(CosPhiX)/DegToRad$

Value: 89.99519231

Names: Int_PhiX

Cell: AB2

Formula: $=IF(Int_PhiX \geq 90, Int_PhiX - 180, IF(Int_PhiX \leq -90, 180 + Int_PhiX, Int_PhiX))$

Value: 89.99519231

Names: PhiX

Cell: AC2

Formula: $=ACOS(-1*(ABS(Ko)*COS(Attack_Ang*DegToRad)+ABS(Ks)*COS(Detector_Theta*DegToRad))/ABS(Kia))/DegToRad$

Value: 130

Names: Kia_Angle_X

Cell: AD2

Formula: =Kia*COS(Kia_Angle_X*DegToRad)

Value: -152738.5857

Names: Kia_Grad_U

Cell: AE2

Formula: =Kia*COS(Kia_Angle_X*DegToRad)

Value: -152738.5857

Names: Kia_Grad_N

Cell: AF2

Formula: =Ks*COS(Detector_Theta*DegToRad)

Value: 76551.07498

Names: Ks_Grad_U

Cell: AG2

Formula: =Ks*COS(Detector_Theta*DegToRad)

Value: 76551.07498

Names: Ks_Grad_N

Cell: AH2

Formula:

Value:

Names:

Cell: AI2

Formula: 0.58375

Value: 0.58375

Names: Ti

Cell: AJ2

Formula: =979000*SQRT(1000*Ti/Mass_bar)

Value: 3417656.273

Names: Vt_Ion

Cell: AK2

Formula: =1320*Z*SQRT((Density*Nc/Z)/Mass_bar)

Value: 6.08658E+12

Names: Wpi

Cell: AL2

Formula: 2.0959

Value: 2.0959

Names: Te

Cell: AM2

Formula: $=41900000 \cdot \text{SQRT}(\text{Te} \cdot 1000)$

Value: 1918223918

Names: Vt

Cell: AN2

Formula: $=Z \cdot \text{Te} / \text{Ti}$

Value: 78.98895075

Names: ZTeTi

Cell: AO2

Formula: $=\text{SQRT}((3 \cdot \text{Ti} / \text{Mass_ion}) + (\text{Zbar} \cdot \text{Te} / (\text{Mass_ion} \cdot \text{P})))$

Value: 19783590.97

Names: Cs

Cell: AP2

Formula: $=(\text{Ks} / \text{Ws}) \cdot \text{SOL}^2$

Value: 29805182741

Names: Vs

Cell: AQ2

Formula: $=\text{Vs} \cdot \text{COS}(\text{Detector_Theta} \cdot \text{DegToRad})$

Value: 19158402171

Names: Vs_Grad_U

Cell: AR2

Formula: $=\text{Vs} \cdot \text{COS}(\text{Detector_Theta} \cdot \text{DegToRad})$

Value: 19158402171

Names: Vs_Grad_N

Cell: AS2

Formula: $=\text{VelocityZ} \cdot \text{COS}(\text{Theta_UZ} \cdot \text{DegToRad}) + (\text{Kia_Grad_N} / \text{Kia}) \cdot ((3 \cdot \text{Ti} / \text{Mass_ion}) + (\text{Zbar} \cdot \text{Te} / (\text{Mass_ion} \cdot \text{P}^2))) / \text{Cs}$

Value: -76939929.14

Names: Via_Gradient

Cell: AT2

Formula:

Value:

Names:

Cell: AU2

Formula: $=1 + \text{Kia}^2 \cdot \text{Debye}^2$

Value: 2.411234107

Names: P

Cell: AV2

Formula: $=743*\text{SQRT}(\text{Te}*1000/(\text{Density}*\text{Nc}))$

Value: 4.9994E-06

Names: Debye

Cell: AW2

Formula: $=31-\text{LN}(\text{SQRT}(\text{Nc})/\text{Te})$

Value: 6.874629759

Names: LogLam

Cell: AX2

Formula: $=0.00000000009*(\text{Density}*\text{Nc})*\text{Z}*\text{LogLam}/(\text{Te}^{1.5})$

Value: 2.07668E+11

Names: Nuei

Cell: AY2

Formula: $=(30000000/\text{LogLam})*\text{SQRT}(0.5*\text{Mass_bar}*(\text{Ti}*1000)^{1.5}/((\text{Z}^3)*\text{Density}*\text{Nc}))$

Value: 6.11062E-13

Names: Col_Time, Ion_Collision_Time

Cell: AZ2

Formula: $=\text{Vt_Ion}*\text{Ion_Collision_Time}$

Value: 2.0884E-06

Names: Ion_mfp

Cell: BA2

Formula: $=\text{Kia}*\text{Ion_mfp}$

Value: 0.496243682

Names: Kia_mfp

Cell: BB2

Formula: $=1+0.4*(\text{ZTeTi}^{1/3})*\text{Kia}*\text{Debye})^{5/3}$

Value: 7.038504619

Names: g_1

Cell: BC2

Formula: $=(1+(\text{Kia}*\text{Debye})^2)^{-1.5}$

Value: 0.267079733

Names: g_2

Cell: BD2

Formula: $=\text{ZTeTi}^{0.1}*\text{EXP}(-1*(0.05*(\text{ZTeTi}-1)/\text{SQRT}(\text{Kia}*\text{Debye}+1))^{4})$

Value: 1.63632E-21

Names: alpha

Cell: BE2

Formula:=g_2+alpha*(g_1-g_2)

Value: 0.267079733

Names: g

Cell: BF2

Formula:=0.3+5*EXP(-0.25*ZTeTi)

Value: 0.300000013

Names: p_ion

Cell: BG2

Formula:=1+0.7*EXP(-1*ZTeTi^4)

Value: 1

Names: q_ion

Cell: BH2

Formula:=1.5-0.5*EXP(-1*(0.1*(ZTeTi-1))^2)

Value: 1.5

Names: s_ion

Cell: BI2

Formula:=3*SQRT(ZTeTi)*EXP(-2*SQRT(ZTeTi))

Value: 5.0846E-07

Names: A1_ion

Cell: BJ2

Formula:=SQRT(0.125*PI())*ZTeTi^1.5*EXP(-0.5*ZTeTi)

Value: 3.09847E-15

Names: A2_ion

Cell: BK2

Formula:=SQRT(0.125*Delta*PI())

Value: 0.00991145

Names: Ge

Cell: BL2

Formula:=q_ion*(A1_ion^(-p_ion)+A2_ion^(-p_ion))^(-1/p_ion)

Value: 3.06329E-15

Names: Gi

Cell: BM2

Formula:=(Ge*g)+((0.5*Kia_mfp/SQRT(ZTeTi))+Gi*g*Kia_mfp^2)/(1+0.25*ZTeTi*Kia_mfp^s_ion+Kia_mfp^2)

Value: 0.006072888
Names: lon_Damp_Wia

Cell: BN2
Formula: =lon_Damp_Wia*Wia_No_Shift
Value: 18384929758
Names: lon_Damp_Casanova

Cell: BO2
Formula: =0.5*Nuei*(Wpe/Ws)^2
Value: 1202246764
Names: Damp_slw

Cell: BP2
Formula:
Value:
Names:

Cell: BQ2
Formula: =25.6*(Wave_0*0.001)*SQRT(Intensity)
Value: 76805031.08
Names: Vosc

Cell: BR2
Formula: 0.27637
Value: 0.27637
Names: S_Density

Cell: BS2
Formula: 7.5837
Value: 7.5837
Names: S_Te

Cell: BT2
Formula:
Value:
Names:

Cell: BU2
Formula: =Wpe^2*Kia_Grad_N/(2*S_Density*SOL^2*Ko*Ks_Grad_N)
Value: -5059.865804
Names: Kappa_Light

Cell: BV2

Formula: $= -\text{COS}(\text{ABS}(\text{Theta_GradU} * \text{DegToRad})) * \text{Kia_Grad_N} * \text{Velocity_Gradient} / \text{Cs}$

Value: 5481532.449

Names: Kappa_Ion1

Cell: BW2

Formula: $= (0.5 * \text{Kia} / \text{P}^{1.5}) * ((1 / \text{S_Te}) + ((\text{P} - 1) / \text{S_Density}))$

Value: 166216.0294

Names: Kappa_Ion2

Cell: BX2

Formula: $= \text{Cs} * (+\text{Mach} * \text{COS}(\text{Theta_GradU} * \text{DegToRad}) + (\text{Kia_Grad_N} / \text{Kia}) / \text{P}^{1.5})$

Value: -74396358.72

Names: Via_Drake

Cell: BY2

Formula: $= (\text{Cs} / \text{Via_Gradient}) * (\text{Kappa_Ion1} + \text{Kappa_Ion2})$

Value: -1452207.548

Names: Kappa_Ion

Cell: BZ2

Formula: $= \text{Kappa_Light} + \text{Kappa_Ion}$

Value: -1457267.413

Names: KappaPrime

Cell: CA2

Formula: $= \text{Via_Gradient} * \text{Kappa_Light} + \text{Cs} * (\text{Kappa_Ion1} + \text{Kappa_Ion2})$

Value: 1.12122E+14

Names: KapPrViaX

Cell: CB2

Formula: $= 0.25 * \text{Wpi} * \text{Kia} * \text{Vosc} * \text{COS}(\text{Polar} * \text{DegToRad}) / \text{SQRT}(\text{Ws} * \text{Cs} * \text{Kia})$

Value: 2.13734E+11

Names: Gamma

Cell: CC2

Formula: $= \text{ABS}(\text{KapPrViaX} * \text{Vs_Grad_N})$

Value: 2.14808E+24

Names: KPrimeV1V2

Cell: CD2

Formula: $= 2 * \text{PI}() * \text{Gamma}^2 / \text{KPrimeV1V2}$

Value: 0.133622108
Names: Rosenbluth_Gain

Cell: CE2
Formula: =SQRT(Damp_slw*Ion_Damp_Casanova)/Gamma
Value: 0.021996491
Names: G_EAW

Cell: CF2
Formula: =2*Rosenbluth_Gain*(ACOS(G_EAW)-G_EAW*SQRT(1-G_EAW^2))/PI()
Value: 0.129880082
Names: EAW_Gain

Cell: CG2
Formula:
Value:
Names:

Cell: CH2
Formula: =SIGN(Via_Gradient*Vs_Grad_N)
Value: -1
Names: Sign_V1V2

Cell: CI2
Formula: =IF(Sign_V1V2<0,0.5*SQRT(ABS(Vs_Grad_N*Via_Gradient))*((Damp_slw/ABS(Vs_Grad_N))+Ion_Damp_Casanova/ABS(Via_Gradient)),1000000000000000000)
Value: 1.45094E+11
Names: Gamma_Absolute_Damp

Cell: CJ2
Formula: =Gamma/Gamma_Absolute_Damp
Value: 1.473074179
Names: G_G_Absolute_Damp

Cell: CK2
Formula: =IF(G_G_Absolute_Damp<0.005,1000000000000000000,Intensity*(G_G_Absolute_Damp)^-2)
Value: 1.49358E+13
Names: Threshold_I_A_Damp

Cell: CL2

Formula:=IF(Sign_V1V2<0,SQRT(ABS(Vs_Grad_N*Via_Gradient*KappaPrime)),1000000000000000000)

Value: 1.46563E+12

Names: Gamma_Absolute_Gradient

Cell: CM2

Formula:=Gamma/Gamma_Absolute_Gradient

Value: 0.145830789

Names: G_G_Absolute_Gradient

Cell: CN2

Formula:=IF(G_G_Absolute_Gradient<0.005,1000000000000000000,Intensity*(G_G_Absolute_Gradient)^-2)

Value: 1.52398E+15

Names: Threshold_I_A_Gradient

Cell: CO2

Formula:=SQRT(ABS(Vs_Grad_N*KapPrViaX))

Value: 1.46563E+12

Names: Gamma_Convective_Gradient

Cell: CP2

Formula:=Gamma/Gamma_Convective_Gradient

Value: 0.145830789

Names: G_G_Convective_Gradient

Cell: CQ2

Formula:=IF(G_G_Convective_Gradient<0.005,1000000000000000000,Intensity*(G_G_Convective_Gradient)^-2)

Value: 1.52398E+15

Names: Threshold_I_C_Gradient

Cell: CR2

Formula:=SQRT(Damp_slw*Ion_Damp_Casanova)

Value: 4701406417

Names: Gamma_Conv_Damp

Cell: CS2

Formula:=Gamma/Gamma_Conv_Damp

Value: 45.46179621

Names: G_G_Convective_Damp

Cell: CT2

Formula: $\text{Intensity} * (\text{G_G_Convective_Damp})^{(-2)}$

Value: 15681405880

Names: Threshold_I_C_Damp

Cell: CU2

Formula:

Value:

Names:

Cell: CV2

Formula: $\text{IF}(\text{Intensity} < \text{Threshold_I_C_Damp}, 0, 1 - (\text{G_G_Convective_Damp})^{-2})$

Value: 0.999516154

Names: Fg

Cell: CW2

Formula: 0.9

Value: 0.9

Names: Transmission

Cell: CX2

Formula: $(\text{Vt}/\text{Vt_Ion}) * \text{EXP}(-0.5 * (\text{Zbar} * \text{Te} + 3 * \text{Ti}) / \text{Ti})$

Value: 1.30307E-14

Names: Q_Tac

Cell: CY2

Formula: $\text{Te} * ((1 + \text{Zbar} * \text{Q_Tac}) / (1 + \text{ZTeTi} * \text{Q_Tac}))$

Value: 2.0959

Names: Tac

Cell: CZ2

Formula: $1 - \text{Transmission}$

Value: 0.1

Names: b_1

Cell: DA2

Formula: $0.5 * (\text{Wo}/\text{Wia_No_Shift}) * (\text{SOL}/\text{ABS}(\text{Vs_Grad_N})) * (\text{Tac}/\text{Te})$

Value: 923.7504739

Names: b_2

Cell: DB2

Formula: $0.00000000000000016 * \text{SOL} * (10000 / (0.001 * \text{Wave_0}))^3 * \text{Te} * \text{Area} / (0.001 * \text{Wave_0})$

Value: 4.09467E+12

Names: Po

Cell: DC2

Formula:=(1-Transmission)*Po+Po*Transmission*((b_1+b_2)*
EXP(EAW_Gain*Fg)-b_2)

Value: 4.72721E+14

Names:

Cell: DD2

Formula:

Value:

Names:

Cell: DE2

Formula:=SOL*(Tac*1.602E-19*1000)/(Wavelength*(
Wavelength*0.0000001)^3)

Value: 132979.6369

Names: BB

Cell: DF2

Formula:=COS(Detector_Theta*DegToRad)^2/SQRT(COS(
Detector_Theta*DegToRad)^2-(Wpeak^2/Ws^2))

Value: 0.688969006

Names: Refraction

Cell: DG2

Formula:=0.5*Ws/ABS(Wia_No_Shift)

Value: 593.1120053

Names: Photon_Conservation

Cell: DH2

Formula:=BB*Transmission*Refraction*Photon_Conservation*(
EXP(EAW_Gain)-1)

Value: 6782892.242

Names: Flux_Langmuir

Cell: DI2

Formula:=BB*((1-Transmission)+(1-Transmission)*Transmission*(
EXP(EAW_Gain)))

Value: 26926.01799

Names: Flux_bremsstrahlung

Cell: DJ2

Formula:=Flux_Langmuir+Flux_bremsstrahlung

Value: 6809818.26

Names: Total_Scattered_Flux

Cell: DK2

Formula: =ABS(Delta_Wavelength)*Total_Scattered_Flux

Value: 0

Names: Scattered_Power

Cell: DL2

Formula: =IF(AND((Intensity>Threshold_I_A_Damp),(Intensity>Threshold_I_A_Gradient)),Intensity-MAX(Threshold_I_A_Damp,Threshold_I_A_Gradient),Scattered_Power)

Value: 0

Names: Total_Scattered_Power

END

DATE
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10/01/93