THE ROLE OF ATOMIC COLLISIONS IN FUSION

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ABSTRACT

Atomic physics issues have played a large role in controlled fusion research. A general discussion of the present role of atomic processes in both magnetic and inertial controlled fusion work is presented.

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INTRODUCTION

Atomic collision processes play a large role in both magnetic and inertial confinement fusion research. There are four principal areas of fusion research where atomic processes are important: (1) the hot central plasma where the fusion reactions are expected to occur, (2) the plasma edge where the plasma interacts with the external environment, (3) plasma heating methods used to produce the hot plasma, and (4) diagnostic techniques used to measure the physical properties of the plasma. In magnetic fusion research, atomic processes such as line radiation and charge exchange have contributed significantly to the energy and particle balance in the hot center of the plasma. Atomic processes involving low Z ions, hydrogen and impurity molecules, and ion-solid collisions are important at the plasma edge. Neutral beam heating has been crucial to the success of the tokamak and mirror programs. Many of the techniques used to diagnose magnetic fusion experiments rely on electron and ion collision processes.

Atomic physics is even more important in inertial confinement research. A knowledge of equations of state and opacities is crucial to designing fusion target pellets and to understanding their behavior. The absorption and stopping of laser light and energetic ions in the outer layers of an imploding pellet, the design and construction of high energy lasers and ion beams, and the use of spectroscopy and laser scattering to diagnose pellet implosions all rely on atomic collisions of one type or another.

These lectures will introduce all of these issues. I will first cover the general issues in fusion such as fusion reactions, confinement requirements, and general reactor schemes. In the second section, I will begin to discuss magnetic confinement and introduce the experimental and theoretical plasma physics issues in magnetic confinement. The third section will cover the role of impurities in the energy balance in magnetic fusion experiments. The fourth section will concentrate on the atomic collision issues important at the plasma edge. The fifth section will discuss plasma heating, and the sixth section will cover plasma diagnostics. The seventh section will cover inertial confinement. The eighth section will discuss the role of plasma physics in atomic physics, and the ninth will list useful review articles.

It is useful to first describe the conditions (electron and ion temperature, density, etc.) that characterize inertially and magnetically confined plasmas. These conditions can be derived from a few simple assumptions.

Most currently envisioned fusion schemes rely on the nuclear reaction $D + T \rightarrow n + He + 17.5 \text{ MeV}$. The neutron has 14.1 MeV of kinetic energy. The mean free path for neutrons in a plasma is much larger than the dimensions of commonly discussed fusion plasmas, and thus the neutron energy is absorbed in a "blanket" structure surrounding the fusion chamber. The charged alpha particle is confined and used to heat the plasma directly, that is to resupply the heat losses of the fusing plasma. The reaction rates of deuterium-tritium (DT) fusion (Greene, 1967) are large only for a temperature $T \gtrsim 10 \text{ keV}$ (Fig. 1). This sets a lower limit on the temperature required at the center of a fusion plasma. A hot plasma will try to cool in a characteristic loss time τ_E due to various energy loss processes including conduction, convection,

expansion, radiation, etc. A second condition for a fusion plasma can be introduced from the requirement that the alpha particles produced from the fusion reaction deposit their heat in the plasma to balance the energy losses. A simple energy balance equation for a DT plasma in which the alpha heating exceeds the energy loss rate is

$$\frac{3}{2} k \frac{\partial(\sum n_i T_i)}{\partial t} = \frac{3}{2} k \frac{(n_D T_D + n_T T_T + n_e T_e)}{\tau_E} \lesssim n_D n_T \langle \sigma v \rangle_{DT} E_\alpha, \quad (1)$$

where the sum is over the different ion species and the electrons, τ_E is the plasma energy confinement time, n_i and T_i are the density and temperature of plasma species i , $\langle \sigma v \rangle_{DT}$ is the fusion reaction rate of deuterium-tritium (DT) (Fig. 1), and E_α is the kinetic energy of the alpha particle (3.5 MeV). Assuming that $n_D = n_T = (n/2)$ and that $T_D = T_T = T_e$ and denoting n_e by n , one obtains the condition $n\tau_E \gtrsim (12kT/\langle \sigma v \rangle E_\alpha)$ for ignition. $n\tau_E$ for DT is a strong function of temperature (Fig. 2), and has a minimum at $T \sim 25$ keV, with $n\tau_E \gtrsim 1.5 \times 10^{14}$ sec/cm³. Thus one goal of controlled fusion research is the attainment of $T \gtrsim 10$ keV and $n\tau_E \gtrsim 10^{14}$ sec/cm³.

We can generalize these considerations to include radiation losses explicitly. Taking bremsstrahlung as an example, we can equate bremsstrahlung and other losses to the alpha heating

$$\frac{3nT}{\tau_E} + bn^2 \sqrt{T} = \frac{1}{4} n^2 \langle \sigma v \rangle E_\alpha,$$

or

$$n\tau_E = \frac{3T}{(1/4)\langle \sigma v \rangle E_\alpha - bT^{1/2}}, \quad (2)$$

for ignition ($b = 3.336 \times 10^{-19}$ eV/(sec-cm³(keV)^{1/2}), (Book, D. L, 1980). Ignition can only occur when the denominator of Eq. (2) is greater than zero. The temperature at which this occurs is called the ignition temperature.

Ignition is a stringent criterion requiring very good confinement. We can relax the confinement requirements by recovering some of the lost energy and using it to heat the plasma and derive what is known as the "Lawson criterion" (Lawson, 1957). If η is the efficiency with which we can recover the lost energy and convert it to plasma heating, we can write

$$\begin{aligned} \frac{3nT}{\tau_E} + bn^2 T^{1/2} &= \frac{1}{4} n^2 \langle \sigma v \rangle E_\alpha \\ &+ \eta \left(\frac{3nT}{\tau_E} + bn^2 T^{1/2} + \frac{1}{4} n^2 \langle \sigma v \rangle E_n \right), \end{aligned} \quad (3)$$

where E_α is the energy of the alpha particles (3.5 MeV), and E_n is the energy of the neutrons (14.1 MeV). This yields a "Lawson" condition for $n\tau_E$ of the form

$$n\tau_E = \frac{3T(1-\eta)}{(1/4) \langle \sigma v \rangle (E_\alpha + nE_n) - (1-\eta) bT^{1/2}} \quad (4)$$

Another way of characterizing the economic potential of a fusion system is to define an energy multiplication factor $Q_E = \text{Energy extracted/Energy input}$. Then

$$Q_E = \eta \frac{P_I + P_F}{P_I} \quad (5)$$

where P_F is the fusion power plus the power generated by neutron reactions in the material surrounding the reactor, and P_I is the power one has to inject into the plasma to balance the energy losses. $Q_E = 1$ corresponds to the Lawson condition in which the reactor is able to run itself, and $Q_E = \infty$ corresponds to ignition. Another commonly defined Q is $Q_p = P_F/P_I$, the power multiplication factor.

The DT reaction requires a source of tritium which is usually assumed to be produced by $n + \text{Li}$ reactions (breeding) in a blanket surrounding the burning DT plasma. This complicates reactor designs. There is therefore an advantage to use only deuterium. The reaction chain would be (Rose and Clark, 1961)

The major disadvantage of this scheme is that the reaction rates are a factor of 10 to 100 lower than DT (Fig. 1) except at very high temperatures. The lower reaction rates for DD require better confinement (higher $n\tau_E$) (Fig. 2) and higher operating temperatures than DT. Fusion with DD requires temperatures of about 40 keV and $n\tau_E \sim 2 \times 10^{15} \text{ sec/cm}^3$ and is thus much harder than fusion with DT. Consequently all designs for fusion experiments in this century are based on DT.

The word "scientific breakeven" is often used with regard to the next generation of fusion experiments. It usually refers to $Q_p \sim 1$; i.e., the fusion power and the heating power are roughly equal. While $Q_p \sim 1$ represents an important milestone, it still falls short of the Q_p of ten or more needed for an economic pure fusion reactor.

At the moment there are two approaches to producing the temperatures and confinement needed for fusion. One approach began in about 1952 and uses the

pressure of magnetic fields to confine the plasma. The plasma is heated by energetic neutral atom beams, ion beams, internal currents, or radio frequency waves. The other approach, at least for peaceful uses, began sometime thereafter (mid 1960's), and is basically an attempt to scale down a hydrogen bomb. The basic idea is to use a laser beam, ion beam, or electron beam to heat the surface of a pellet (Nuckolls et al., 1972). The heated surface of the pellet "boils off," accelerating the outer layers of the pellet inward. The spherical convergence of the pellet compression would produce a high density and high temperature plasma in the center of the pellet for a long enough time to ignite deuterium and tritium. The central pellet plasma would be confined by its own inertia thus leading to the term "inertial confinement."

MAGNETIC CONFINEMENT, THEORY AND EXPERIMENTS

We can use very general considerations to derive the physical conditions that characterize the atomic processes we have to consider for magnetic fusion. Since a temperature of 10 keV is far hotter than can be tolerated by any material, the plasma must be kept out of contact with any containment vessel. The idea of magnetic confinement is to use the pressure of high strength magnetic fields (1-10 Tesla) to confine the plasma. The plasma pressure is limited by the requirement that it cannot be much greater than the magnetic field pressure. Since most engineering structures can support, at best, pressures of about 1000 atmospheres, it is reasonable to expect the plasma pressure to be on the order of 1000 atmospheres or smaller. Since the pressure is the product of the density and temperature ($p = nT$), a temperature of 10 keV and a pressure of 1-1000 atmospheres implies densities ranging from $6 \times 10^{13} \text{ cm}^{-3}$ up to $6 \times 10^{16} \text{ cm}^{-3}$. Realistic systems would be expected to operate with densities of 10^{13} to 10^{15} cm^{-3} . For the confinement condition $n\tau_E \sim 10^{14} \text{ sec/cm}^3$, these densities imply $\tau_E \sim 0.1$ to 1 sec.

A plasma can be confined because it consists of electrons (charge $q_e = -e$) and ions (charge $q_z = Ze$) which are affected by magnetic and electric fields. The charged particles in a magnetic field B will execute gyro-orbits of frequency $\omega = (qB/mc)$ and radius $\rho_L = v_\perp/\omega$, where $v_\perp$ is velocity perpendicular to B (see Chen, 1974). The charged particles will move freely along the field. The magnetic field will do no work on the charged particles since the work done by the field on a particle with charge q is $W = \vec{F} \cdot \vec{v} = q(\vec{E} + (\vec{v}/c) \times \vec{B}) \cdot \vec{v} = e\vec{E} \cdot \vec{v}$. Thus a magnetic field provides confinement in two directions (perpendicular to the field). The problem is then how to provide confinement in the third direction along the field lines.

There are, in general, two approaches to magnetic confinement (Fig. 3). The first approach is to provide confinement by bending the field lines around into a torus so that the lines are essentially endless. The second is try to stop the end losses by reflecting the particles from the ends of the field lines. The first approach is termed toroidal or closed confinement, and the second is called linear or open confinement. Examples of toroidal confinement systems are tokamaks, stellarators, and toroidal pinches (see Teller, 1981). Mirror machines and linear pinches are examples of open confinement systems (Teller, 1981). Confinement along the field lines is essential, since the ion velocity ($v \sim (E/m)^{1/2}$) is about $2 \times 10^8 \text{ cm/sec}$. Without confinement along the field lines a confinement time of one second would require a 2000 km long reactor, much beyond what is reasonable.

There are three requirements that any magnetic fusion experiment must meet. The particles must be on confined orbits. The magnetic field created by the external and internal currents must be in a quasi-stable equilibrium (forces must balance in a stable way). The loss rate of particles and energy due to collisions and instabilities must be small.

The net motion of a particle in a magnetic field is given by the motion of the guiding center (the motion of the center of the gyro-orbit). In an inhomogeneous and curved magnetic field, the motion of a particle in a gyration is non-uniform with the result that there is a small net motion of the guiding center of the particle after one gyration. These small net motions add and result in a "drift" of the particle guiding center. A second feature of guiding centers is that there are constants of the motion of the guiding centers (adiabatic invariants) which are useful in describing the particle orbits. One constant is the magnetic moment μ

$$\mu = \frac{W_{\perp}}{B},$$

where $W_{\perp}$ is the kinetic energy perpendicular to the magnetic field B . For a fixed total kinetic energy $E = W_{\perp} + W_{\parallel}$, where $W_{\parallel}$ is the energy parallel to the magnetic field, conservation of μ implies that there is a maximum magnetic field $B = (E/\mu)$ in which the particle orbit can exist. Thus particles can be confined along a field line by making the fields small in the middle and large at both ends (Fig. 3). The particles bounce away from the high field region, hence the term mirror. Of course only part of the particles are confined, those where $(W/W_{\perp}) < B_{\text{end}}/B_{\text{center}}$ ($= R_n$, the mirror ratio) with W and $W_{\perp}$ being evaluated at the middle of the field line. Particles with $(W/W_{\perp}) > R_n$ escape and are lost. Thus scattering events which change the angle of the particle velocity with respect to the field line can change $W/W_{\perp}$ and cause previously confined particles to be lost. In a conventional mirror this turns out to be the dominant loss mechanism.

A new type of mirror machine is being investigated, the tandem mirror (Fowler and Logan, 1977). It is basically a long solenoid with a conventional mirror at each end. This design overcomes some of the problems with the high loss rate of a standard mirror.

An open confinement scheme has a confinement time of about an ion-ion scattering time, which is fairly short. A closed confinement scheme, such as a tokamak or stellarator, bends the field lines around in a torus so that they remain in the plasma forever. The particle orbits, in both cases, are described by drifts and the conservation of μ . Orbits for mirrors and toruses are reviewed in Siambus, et al. (1965) and Furth and Rosenbluth (1969).

In a simple toroidal field, the particles would drift vertically due to the curvature and gradients of the toroidal field. Since ions and electrons would drift in opposite directions, a charge separation would develop, creating an electric field in the vertical direction. Another drift, the $\vec{E} \times \vec{B}$ drift, would cause the plasma to expand outward, and there would be no equilibrium. Equilibrium can be obtained if the magnetic field lines are made helical (Fig. 4) so that the particle drifts average to zero. The current in a tokamak is usually induced by making the plasma the secondary

winding of a transformer. In a tokamak, the helical pitch of the field is then provided by an internal current. In a stellarator, it is provided by external helical windings.

The study of how the forces produced by the plasma pressure and the forces produced by the magnetic and electric fields balance in a plasma is called magnetohydrodynamics (MHD). The magnetic and electric fields are produced by internal and external currents and fields. An important quantity in MHD is the ratio of the plasma pressure to the magnetic field pressure, $\beta = 2nT/(B^2/8\pi)$. As β increases, the plasma typically becomes more unstable. Theoretical estimates of the upper limit on β in tokamaks are four to seven percent (some estimates go as high as twenty percent). Experiments have reached two and one-half percent. Mirror experiments have reached betas of fifty percent and even one-hundred percent. The pressure, nT , is fixed by the fusion power density needed. Raising the magnetic field strength raises the cost of the magnets and thus the cost of a reactor. There is thus a great incentive to use as high a β as possible, which is a clear advantage mirror machines may have over tokamaks. However, both types of machines have a large volume which contains magnetic field but no plasma, so a detailed comparison requires detailed designs. A good review of MHD equilibrium and stability issues for mirrors is Taylor (1963) and for toruses in Solovév and Shafranov (1967).

The biggest advantage of toroidal systems is that their particle and energy losses should be diffusive. A collision which changes the particle velocity should only cause the guiding center to move from the original field line to an adjacent one. Thus energy would be lost through conduction and convection perpendicular to the magnetic field.

In a tokamak, the helical field lines map out closed surfaces called flux surfaces. The plasma temperature and density are reasonably constant along the field lines, so the energy is lost by diffusion perpendicular to the flux surfaces from the center to the plasma edge. A naive estimate of the thermal diffusivity across the field lines is given by $D \sim (r_L^2/\tau_1)$ where r_L is the gyroradius and τ_1 is the ion-ion scattering time (Spitzer, 1963). We can write a simple energy balance equation for a cylinder

$$\frac{\partial E}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} \left(r D \frac{\partial E}{\partial r} \right),$$

where r is the distance from the center of the plasma, and E is the plasma energy density. This can be approximated as

$$\begin{aligned} \frac{E}{\tau_c} &\sim D \frac{E}{(a/2)^2}, \\ \text{or} \\ \tau_c &\sim \frac{a^2}{4D} = \tau_1 \frac{a^2}{4r_L^2}, \end{aligned}$$

where a is the minor radius of the torus and τ_c is the confinement time. Thus the confinement time has been increased by the factor $a^2/4r_L^2$ compared to a

mirror ($\tau_c \sim \tau_i$ for a mirror). However, things are, as always, a little more complicated. Many of the particles are "mirror trapped" in the toroidal field since it increases toward the center of the torus, and the correct gyroradius to use is not the one for the toroidal field, but the one for the poloidal field (the field the short way around the torus) set up by the plasma current. Since $B_p \sim 0.1 \times B_T$, the factor $a^2/4r_L^2$ is not as large as one might at first expect. However, it is still > 50 for most tokamak and stellarator experiments.

The measured loss rates are, unfortunately, greater than the rates calculated using more exact versions of the classical theory above. These "anomalous" losses are probably due to small scale turbulence which either increases the effective step size across the field lines between collisions, or decreases the collision time. The physics of this loss process is the focus of much of the experimental and theoretical work in the fusion community. Baldwin (1977) presents a review of classical and anomalous losses in mirrors. The theory of classical losses for toruses is reviewed in Hinton and Hazeltine (1976) and for anomalous losses in Tang (1978).

Table I lists the major tokamak experiments in the United States. I will concentrate on atomic physics issues in tokamaks, since they are the best developed magnetic confinement experiments. It should be noted that mirror machines offer the promise of superior MHD properties, and have some reactor design advantages as well. If they succeed in achieving the good confinement of tokamaks, they should have similar atomic physics problems. A recent review of mirror research is given in Cohen (1980). Two recent reviews of tokamaks are given in Rawls (1979) and Furth (1975).

To date tokamaks have achieved central temperatures of about 7 keV, densities of $10^{14} - 10^{15} \text{ cm}^{-3}$, and $nT_e \sim 3 \times 10^{13} \text{ sec/cm}^3$. Mirror experiments (TMX) have achieved temperatures of about 200 eV, densities of about 5×10^{13} and $nT \sim 7 \times 10^{10} \text{ sec/cm}^3$ (Coensgen et al., 1980). The growth in the United States tokamak fusion program is illustrated in Table II. Similar sized tokamaks to TFTR are being constructed throughout the world (JT-60, Japan; JET, Europe; and T-15, USSR). The mirror program is building a large experiment at Livermore, California, costing about \$200 million, to be ready in about 1984. The United States magnetic fusion budget is about \$450,000,000 a year, and the worldwide budget is about $\$10^9/\text{year}$.

IMPURITIES

Ideally, the plasma in a fusion experiment consists entirely of hydrogen ions and electrons. Unfortunately, this is often not the case. Ions of elements used in the construction of the vacuum vessel, walls and limiters, and lighter elements usually associated with high vacuum systems are often present in significant quantities. These "impurity" ions degrade the fusion potential of the plasma. They contribute to the plasma pressure without adding significantly to the fusion yield. Since the MHD stability limits are usually set by a maximum $\beta = P/(B^2/8\pi)$ where $P = \sum n_i T_i$, summed over the plasma constituents, impurities "dilute" the reacting plasma reducing the fusion yield for a given β . Fully stripped impurities also increase the bremsstrahlung losses per ion since these losses are proportional to Z^2 .

Dilution and enhanced bremsstrahlung are usually the major concern for fully stripped ions. However, partially stripped ions can lose appreciable energy through line radiation. At densities of 10^{14} cm^{-3} , the plasma is optically thin, and the radiation goes directly from the center of the plasma to the walls thus short circuiting the plasma confinement (Fig. 5).

Energy losses from the plasma edge, such as radiation from partially stripped low Z ions like carbon or oxygen and losses due to the charge exchange of hot ions and cold atoms, do not degrade the confinement since the plasma energy has already reached the edge. In fact, low Z impurity radiation may help reduce the high Z impurity level by reducing the power that strikes the wall and limiters as charged particles. Energetic particles can produce high Z impurities by sputtering wall or limiter materials.

The major energy losses come from electron impact excitation of bound electrons which decay radiatively (Jensen, et al., 1977). From the requirement that $n \sim 10^{14} \text{ cm}^{-3}$ and $T \sim 10 \text{ keV}$, we can specify the conditions appropriate for the atomic processes in the central plasma. The radiative lifetime for most excited states (Einstein $A \sim 10^{10} \text{ sec}^{-1} \Rightarrow \tau_{\text{rad}} \sim 10^{-10} \text{ sec}$) is short compared to the electron excitation collision time ($\nu \sim 10^5 \text{ sec}^{-1} \Rightarrow \tau_{\text{ee}} \sim 10^{-5} \text{ sec}$). This condition in which the ions are in the ground state before excitation is usually referred to as coronal. One can then write an ionization balance equation of the form

$$\frac{\partial n^i}{\partial t} = R^{i+1 \rightarrow i} n^{i+1} + I^{i-1 \rightarrow i} n^{i-1} - R^{i \rightarrow i-1} n^i - I^{i \rightarrow i+1} n^i, \quad (6)$$

where n^i is the density of charge state i , $R^{i+1 \rightarrow i}$ is the frequency for recombination from charge state $i+1$ to i , and $I^{i-1 \rightarrow i}$ is the frequency for ionization from charge state i to $i+1$. It is useful to study the "coronal equilibrium" case with $\partial n^i / \partial t = 0$ for all i from 0 to Z . Then $R^{i+1 \rightarrow i} n^{i+1} = I^{i \rightarrow i+1} n^i$. If we know the recombination and ionization rates, we can solve for the charge state distribution.

The ionization rate can be written as $n_e \langle \sigma v \rangle$ where n_e is the electron density, and $\langle \sigma v \rangle$ is the reaction rate for ionization averaged over a Maxwell-Boltzmann electron distribution for a given electron temperature. The rates are usually summed over all the bound electrons. $\langle \sigma v \rangle$ for a given charge state is then only a function of the electron temperature.

Ionization cross sections are usually evaluated theoretically and benchmarked by a few measurements. Commonly used rates were recently reviewed by De Michelis and Mattioli (1981). As an example of the rates that are typically used by plasma modelers, I will describe the rates used in our code at Princeton (Post et al., 1977). In our calculations, we use an amalgam of cross sections from Jordan (1970); Jacobs (1972); Sampson and Golden (1971); Lotz (1967), and supplemented by those of Younger (1981).

The dominant recombination processes are radiative and dielectronic. Again we use general prescriptions. In general, dielectronic recombination is

dominant when the ion has three or more bound electrons. Radiative recombination is important when dielectronic recombination is small (0, 1, or 2 bound electrons). Three body recombination is extremely small except at very low temperatures and high densities (1 eV and 10^{14} cm^{-3}). All of the rates (except three body) are roughly proportional to the electron density. Dielectronic recombination becomes small at high densities but at 10^{14} cm^{-3} it is roughly proportional to the electron density. Thus we can write $R^i, I^i = n_e \langle \sigma v \rangle^i$ with the $\langle \sigma v \rangle^i$ being dependent only on the electron temperature.

Using the above rates we can compute the equilibrium fraction of each charge state for any given element as a function of electron temperature for any element. The general nature of our formulae allows us to tackle a high Z element such as tungsten (Post et al., 1979) (Fig. 6) as well as the lower Z elements such as C, O, and Fe.

The next step after determining the ionization balance is to determine how much energy is lost through the excitation of bound electrons through electron impact collisions and their subsequent radiative decays. We use primarily the prescriptions of Seaton (1972) with suitable modifications. For systems with many bound electrons (> 15), $\Delta n = 0$ transitions are extremely important. Using the calculated charge state distribution one can then compute the radiative power loss density as

$$P = n_z n_e \sum_{i=0}^Z f_i \sum_{\text{all excitations}} \Delta E^i \langle \sigma v \rangle^i, \quad (7)$$

where n_z is the impurity density, n_e is the electron density, f_i is the fractional abundance of charge state i , and the sum is over all excitations of ions with charge i . ΔE^i is the excitation energy and $\langle \sigma v \rangle^i$ is the excitation rate. Polynomial fits of these cooling rates have been tabulated (Post et al., 1977).

The line radiation is, of course, largest when there are many bound electrons (three or more). In general, an ion will strip out until the ionization potential is of the order of the electron temperature (Fig. 7). Thus lower Z impurities will have little line radiation at high temperatures. Higher Z impurities will have larger power loss rates, not only due to having higher ionization potentials and higher excitation energies, but also because they have many more bound electrons with many possible excitations. Thus the cooling rate for tungsten is about sixty times that for iron and about five hundred times that for oxygen (Fig. 8) (Jensen, 1977).

Although dielectronic recombination had long been recognized as an important recombination mechanism for astrophysical plasmas (Burgess, 1965; Cox, 1969), it was only applied to magnetic fusion work in 1975-1976 (Merts et al., 1976; Jensen et al., 1977). Until then, tungsten limiters and metal rails, inserted to keep the plasma from hitting the walls, were widely used in tokamaks because of their excellent refractory properties. Estimates of the tungsten density in tokamaks were about 0.1-0.3%. If one neglects dielectronic recombination, tungsten strips down to an average charge of +40 at a temperature of 1 keV, and the radiation is not extremely large. Including dielectronic recombination reduces the average charge to about twenty-five at an electronic temperature of 1 keV, with an increase in the

radiation factor of about ten (Jensen et al., 1977). At this level tungsten could radiate and indeed was radiating most of the energy contained in tokamaks such as ORMAK (Isler et al., 1977) and PLT (Hinnov and Mattioli, 1978). PLT even had hollow electron temperature profiles due to the large radiation losses (Hsuan et al., 1978) (Fig. 9). Attempts to heat PLT with the injection of high energy neutral beams only brought more tungsten into the discharge, resulting in the heating power being radiated away with no temperature increase. Calculations were done of the expected tungsten lines (Isler et al., 1977) ($\sim 50 \text{ \AA}$), and these lines were measured. The suspicions about tungsten which had already been reinforced by the power loss calculations were finally confirmed. The tungsten limiters were replaced with carbon limiters, and PLT was subsequently heated to about 7 keV (Eubank et al., 1979). Very high Z materials are no longer used in tokamaks.

Many tokamaks are heated by the injection of beams of high energy neutral atoms (20-40 keV). The presence of these neutral atoms can lead to an additional recombination mechanism through the reaction

This process is particularly important for systems with just a few bound electrons since radiative recombination is inherently small and dielectronic recombination is very small in these cases due to a small number of allowed excitations. The ionization balance equilibrium equation has to be modified (Krupin et al., 1979; Hulse et al., 1980; Puiatti et al., 1981)

$$\frac{n_{i+1}}{n_i} = \frac{I^{i+1+1}}{R^{i+1+1} + (n_n/n_e) v \sigma^{i+1+1}}, \quad (8)$$

where I and R are the usual ionization and recombination rates (dielectronic and radiative), n_n is the neutral density, v is the neutral velocity, and σ^{i+1+1} is the charge exchange cross section. The main effect of this process is to postpone or eliminate the burnout of an ion to non-radiative states as higher temperatures are approached. Modest values of n_n/n_e can significantly affect the ionic balance (Fig. 10) (Hulse et al., 1981). The enhancement of the lithium-like and more highly recombined states can lead to an enhancement in the total impurity radiation (Hulse et al., 1981) (Fig. 11). The effect has only been a major term in the power balance in beam heated tokamaks where the plasma has a fairly high impurity content already, and high power, low energy beams (leading to a high neutral density, n_0) are used.

Clearly, the presence of impurities will increase the plasma energy losses, and thus make fusion more difficult. We can rewrite the energy balance Eq. (2) to include line radiation and beam heating in the form (Jensen et al., 1977; Jensen, et al., 1978)

$$\frac{(3/2) (n_e T_e + \sum_{ions} n_i T_i)}{\tau_E} + R = P_b + P_a, \quad (9)$$

where R is the radiation loss, P_b is the beam heating, P_α is the alpha heating, T_e is the electron temperature, and T_i is the ion temperature. For a DT plasma, $P_\alpha = 0.2 Q P_b$ where $Q = P_f/P_b$. P_f is the fusion power and Q is the fusion power multiplication factor. Equation (9) can be written as

$$n_e \tau_E = \frac{(3/2) [(1+f_z \bar{Z}) T_e + (1+f_z) T_i]}{H - f_z L_z - L_H}, \quad (10)$$

where $H(Q) = P_\alpha/n^2$, f_z is the fraction of the ions that have charge Z , and $\bar{Z}$ is the average charge state. L_z is the radiation rate (in units such as watts cm^3) per free electron per ion, and L_H is the hydrogen bremsstrahlung rate (in units such as watts cm^3). As one can see from Eq. (10) (analogous to Eq. 2) as f_z is increased, the $n_e \tau_E$ required for a given Q (or H) increases. At a large enough f_z , the denominator of Eq. (9) becomes zero and even perfect confinement is not good enough.

At such a point the radiation losses equal the fusion and beam heating. Since the radiation rates are largest for high Z and decrease (in general) as T_e is increased, higher Z raises the $n_e \tau_E$ required and higher T_e lowers it (Fig. 12). We can define a fatal fraction as the f_z at which the denominator of Eq. (10) becomes zero. This then sets an upper limit on the impurity concentration that can be tolerated. If we take as our criteria that the plasma must be ignited, we find that the fatal fraction scales as $\sim Z^{-2.5}$ (Jensen et al., 1978). Less than 10^{-4} tungsten can be tolerated as compared to 10^{-1} carbon (Fig. 13).

Although impurities have played and continue to play an important role in the central plasma, it is always a bad influence by spoiling confinement. Eliminating high Z impurities is viewed as one of the key problems in fusion work.

However, the key questions of the economic and physics viability of fusion reactors still hinge on questions of MHD stability (maximizing β) and plasma transport. Impurities are viewed as a nuisance which hopefully will go away. Atomic processes do not contribute usefully to better confinement or better stability, and most plasma physicists do not particularly care about improving our knowledge of electron impact collision physics in hot plasmas. For them it is enough to know that tungsten is very bad and carbon is not "too" bad. However better knowledge of atomic collision physics in hot plasmas is still crucially important since we do not know how well we will be able to control impurities. In the other three areas of fusion research, the plasma edge, heating methods, and diagnostics, atomic processes play a beneficial role, and most of the qualitatively new applications of atomic processes will probably be in these areas.

PLASMA EDGE

Atomic processes play an important role in the behavior of the tokamak edge plasma. The particle balance (Fig. 14) in a tokamak consists of neutral hydrogen atoms and molecules going into the plasma and undergoing ionization and charge exchange. Some of the hydrogen atoms gain enough energy by charge exchanging with hot plasma ions to increase their mean free path for ionization [$\lambda = (v_0/n_e \langle \sigma v \rangle)$] to penetrate to the center before they are

ionized (Hughes et al., 1978). The plasma ions then diffuse across the flux surfaces until they reach the edge where they recombine at the wall or at a limiter. The ions and neutral atoms strike the wall and limiter and can sputter off material so that some of it enters the plasma and becomes an impurity. Light impurities such as oxygen are also desorbed from the wall due to electron and ion impacts with the wall.

The central portion of a fusion plasma is quite hot (1-5 keV in current tokamaks and greater than 10 keV in a future reactor experiment). Low Z impurities are completely ionized there and do not radiate strongly. The edge plasma near the vacuum vessel walls is generally much cooler with temperatures typically in the 10-50 eV range and sometimes as low as 1-10 eV (McCracken and Stott, 1979). At such low temperatures, atoms of low Z elements such as carbon and oxygen are usually only partially ionized, and thus can radiate strongly. In addition, the plasma edge is usually quite turbulent, and the ions there are poorly confined. Thus ions can recombine by contact with solid structures such as the wall or limiter, and be reintroduced into the edge plasma as neutrals. This "recycling" of impurity ions leads to the production of low ionization states of the impurity ions born from the ionization of the "recycling" neutral impurity atoms. If these ions can diffuse in time scales short compared to ionization times, particularly the time required to ionize an atom from a neutral atom to a helium-like ion, ions with excitation energies much less than the electron temperature can exist. This lowering of average charge state below "coronal equilibrium" can lead to a substantial enhancement of the radiation rate compared to "coronal equilibrium." Enhancement of the recombination due to charge-exchange recombination with the recycling neutral hydrogen can augment this radiation enhancement as well (Hogan, 1981). For example, if a neutral carbon atom is instantaneously deposited into a 10^{13} cm^{-3} , 50 eV plasma, it will radiate about 10 keV before reaching the helium-like state appropriate to coronal equilibrium. The sum of the single electron ionization potentials for carbon from 0 to +4 is 150 eV. If the carbon +4 ion then hits the limiter, recombines and reenters the plasma, it loses another 10 keV in radiated energy before reaching $q = +4$. Ordinary recombination times are often larger than transport times near the edge. This problem has been studied by solving time dependent equations for the transport of the ionization stages for carbon and oxygen including recombination and ionization (Hawryluk et al., 1979) as a generalization of Eq. 6. The relevant equations are

$$\frac{\partial n^i}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} (r \Gamma^i) + R^{i+1 \rightarrow i} n^{i+1} + I^{i-1 \rightarrow i} n^{i-1} - n^i (R^{i \rightarrow i-1} + I^{i \rightarrow i+1}), \quad (11)$$

where Γ^i is the impurity flux ($\Gamma^i = D (\partial n^i / \partial r)$), for example.

As we have said before, impurity radiation from the edge may be beneficial. If we examine the energy balance in a tokamak (Fig. 5), we see that while radiation from the center is detrimental to energy confinement, radiation from the edge is not, since the energy would, in any event, be quickly lost from the edge by convection and conduction. In fact losing energy from the edge by having high energy particles strike the wall or

limiters can lead to the sputtering of metallic impurities which then get to the plasma center and radiate from there. Losing the energy from the edge by photons does not lead to much sputtering, and is a very attractive way of removing the thermal plasma energy losses. Low Z elements such as oxygen and carbon often work well in current tokamaks since they are fully stripped at the plasma center and can radiate strongly at the edge.

The interaction of the edge plasma with the vacuum vessel and limiters is important because it can lead to the introduction of unacceptable levels of moderate Z and high Z impurities into the plasma. The most commonly studied mechanism for impurity production is chemical and physical sputtering of metal surfaces by ions and neutrals. Extensive data in the form of the yield per atom as a function of the projectile energy has been obtained and catalogued (McCracken and Stott, 1979; Roth et al., 1979). While it is expected that sputtering will be the dominant impurity production mechanism on future large experiments, it has not been fully established that it is the dominant mechanism on current experiments.

In addition to the impurity control problem, another problem with fusion reactors is exhausting the helium ash which results from the DT fusion reaction. On an INTOR-sized experiment (INTOR, 1980), a reactor feasibility experiment, the helium is produced at the rate of one percent (of total particles)/sec. Thus helium must be exhausted and replaced with D and T if the fusion reactor is to burn for longer than about twenty seconds.

One approach to solving the impurity control and helium pumping problems for INTOR has been to use a poloidal divertor (Fig. 15), a device in which the edge plasma is swept away from the main plasma to a chamber where it can be controlled. A large number of complicated atomic processes take place when a plasma strikes a metal surface. These effects have been modeled (Heifetz et al., 1981; Petravic et al., 1981) including a large number of reactions (Table III). The plasma flows into the divertor (Fig. 16), strikes the neutralizer plate, and recombines. The resulting neutrals move around the divertor, undergoing charge exchange until they are either ionized or go down the pump or return to the main plasma. The reflection properties of the 1-500 eV neutrals and ions are a crucial ingredient in the model. What is used now is a collection of theoretical and experimental results (Eckstein and Verbeek, 1979). There is almost no data on low energy collisions (1-30 eV), and the theoretical models are very poor in that energy range.

A key feature of the above mentioned models is the number of molecules formed by low energy hydrogen atoms striking a surface. The assumption is made that in steady state the atoms that are not all reflected as energetic atoms, but some of them penetrate significantly into the crystal lattice and are trapped, eventually diffuse to the surface, and are desorbed as molecules. The molecules undergo reactions 3, 4, and 5 in Table III. If they are ionized (H_2^+), they further undergo reactions 5a and 5b. For moderate electron temperatures the molecule is first ionized, then dissociated into H^+ and H^0 (Fig. 17). This leads to a very localized ionization source near the neutralizer plate, which in turn leads to lowering the plasma temperature and raising the plasma density in the divertor (Petravic et al., 1981). This mode of divertor operation has the potential of reducing the sputtering by reducing the energy of the ions to below the sputtering threshold. It has also been observed experimentally on D-III (Mahdavi et al., 1981; Nagami et al., 1980).

An additional function of a divertor for a reactor experiment would be to provide helium pumping. Early estimates of the pumping speeds needed for reactor experiments were large (500,000 l/sec) and caused a serious design problem (INTOR, 1980). It was proposed that a poloidal divertor could enhance the pumping of helium compared to hydrogen (Shimomura et al., 1979). The argument was that helium would slow down due to collisions with the wall, and be quickly ionized, whereas hydrogen could recover the energy lost by wall collisions by charge exchange with the hot plasma ions, and thus keep a long ionization mean free path. Detailed calculations (Callen et al., 1980; Seki et al., 1980; Heifetz et al., 1981) showed that this was not the case. Charge exchange of hydrogenic ions and neutrals can reduce the directional transport of hydrogen relative to helium because it "isotropized" the hydrogen neutrals. The mean free path for charge exchange for hydrogen is about one-half the mean free path for ionization.

However, "helium enrichment" is not necessary since, if a divertor can be operated in a high density regime (Petravic et al., 1981; Mahdavi et al., 1981; Nagami et al., 1980), the neutral pressure at the pumping duct can be as high as 0.1 torr, and large gas throughputs can be obtained with modest pumping speed systems.

Another area that is beginning to receive attention is the transport of molecular impurities such as CH_4 in the plasma edge (Langer, 1981). The production of methane by the chemical sputtering of carbon limiters and steel surfaces by hydrogen is thought to be a significant source of carbon impurities in current experiments. The chain of events is

The original CH_4 is very slow, and ionizes at the very edge of the plasma. It continues to be dissociated by repeated electron collisions from the 5-20 eV edge plasma. However, CH_x^+ is an ion and can begin to thermalize with the relatively hot (5-20 eV) ions in the edge plasma. Thus a 0.05 to 0.1 eV CH_4 can possibly produce a significant fraction of energetic C^0 (about 10-20 eV) which can then penetrate more deeply into the plasma than a 0.05-0.1 eV C^0 . Calculations for a variety of edge profiles show that the effect is significant and produces an enhancement in the neutral penetration for C^0 of 2-3 cm on the average with a significant tail penetrating up to 10 cm further than C^0 which was launched from the wall with an energy characteristic of the wall temperature. Detailed calculations of this type will become more common as more data becomes available. A good understanding of the plasma edge will grow in importance as the fusion community begins to design realistic reactor level experiments which require the handling of large heat and particle

fluxes. Essential to that understanding will be detailed knowledge of the hydrogen and impurity collision processes in low temperature (1-200 eV) edge plasmas with densities in the $10^{10} - 10^{14} \text{ cm}^{-3}$ range.

PLASMA HEATING

The third area of fusion research involving atomic processes is plasma heating. Since a DT plasma requires a temperature of about 10 keV for heating by alpha particles to be effective, the plasma must first be heated to about 10 keV. Most fusion experiments have large internal currents, which supply power to the electrons in the plasma by drawing the current against the plasma resistance. However, since the plasma resistivity is proportional to $T_e^{-3/2}$ (Spitzer, 1963), this method is most effective at low temperatures. Other heating methods include compressing the plasma, and launching radio waves into the plasma. However, the most successful heating method used for tokamaks and mirrors so far has been the injection of high energy (20-60 keV/amu), high current (20-100 amps) neutral hydrogen beams into the plasma. The neutral atoms cross the confining magnetic field and are "ionized" by collisions with the plasma (Fig. 18). The energetic ions are confined by the magnetic field and heat the background plasma electrons and ions by elastic coulomb collisions (Jassby, 1977).

The beam atoms are "ionized" by electron and ion impact ionization and by charge exchange (Fig. 19) (Freeman and Jones, 1974). Electron ionization is a relatively small fraction of the stopping cross section. Below 40 keV/amu, charge exchange is dominant, and above 40 keV/amu proton impact ionization is dominant. Above about 80 keV/amu the stopping cross section is approximately inversely proportional to the beam energy.

Some early work (Morierte, 1974) suggested that the ion impact ionization cross section, σ_q , for $H^0 + A^{+q} \rightarrow H^+ + A^{+q} + e^-$ should scale as $\sigma_q \approx q^2 \sigma_q$. Given the impurity levels of the then current experiments, it was feared that this effect would raise the stopping cross section so that neutral beams would be ionized at the plasma edge and not heat the center (Hogan and Howe, 1976; Duchs et al., 1977). Furthermore, it was feared that the intense heating of the edge would lead to a large impurity influx, thus stopping the beam even more effectively. For this reason among others, there was great interest in the actual values of the electron loss cross sections for hydrogen and multiply charged ions. As a result of intense theoretical (Olson and Salop, 1976; Olson and Salop, 1977) and experimental work (Phaneuf et al., 1978; Berkner et al., 1978; Bayfield et al., 1976), the dominant collision process was discovered to be charge exchange instead of ionization. The scaling was $\sim Z^{1.2}$, instead of Z^2 . The results have been put in a general formula (Olson et al., 1978) for the combined processes of charge transfer and ionization for $1 \leq q \leq 50$ and $50 \text{ keV/amu} < E < 5,000 \text{ keV/amu}$:

$$\sigma_{\text{loss}} = 4.6 q \frac{(1 - e^{-y})}{y} \times 10^{-16} \text{ cm}^2,$$

where

$$y = \frac{E}{32q}$$

and E is in keV/amu. With this scaling, impurities have very little effect on neutral beam penetration (Fig. 20) for beams of 100 keV/amu or less. This is a consequence of the Z scaling of σ_q for a neutral beam atom. The mean free path would be

$$\begin{aligned}\frac{1}{\lambda} &= n_e \frac{\langle \sigma v \rangle}{v_o} + n_H \alpha_H + n_q \sigma_q \\ &\approx n_e \frac{\langle \sigma v \rangle}{v_o} + n_H \alpha_H + n_q q \alpha_H \\ &\approx n_e \left(\frac{\langle \sigma v \rangle}{v_o} + \alpha_H \right),\end{aligned}$$

where $\langle \sigma v \rangle$ is the electron impact ionization rate, v_o is the neutral velocity, n_H is the hydrogen ion density, and n_q is the density of ions of charge q . This is the same answer obtained for the case with no impurities. However, if $\sigma_q \approx Z^2 \alpha_H$, then

$$\frac{1}{\lambda} \approx n_e \left(\frac{\langle \sigma v \rangle}{v_o} \right) + Z \alpha_H \approx \frac{Z}{\lambda_H},$$

for typical cases, and since Z can be as large as three or four, impurities could lead to very poor penetration of a neutral beam into a tokamak plasma.

The charge exchange of hydrogen ions with partially ionized impurity ions, $H^+ + A^{+q} \rightarrow H^0 + A^{+q+1}$, could serve as an important loss mechanism for fast beam injected ions if the cross sections were greater than 10^{-19} cm^2 and the density of partially ionized impurity ions were large ($\sim 0.1 - 1\%$ of n_e). However, estimates of these cross sections seem to be lower (in the 10^{-21} cm^2 range).

Understanding atomic processes is also crucial to understanding neutral beam systems. Neutral hydrogen beams are made by forming a low temperature plasma (1-10 eV), extracting the protons and accelerating them to $\sim 20-60$ keV/amu while keeping the beam space charge neutralized, and then neutralizing the protons by charge exchange in a gas cell of H_2 . The unneutralized ions are then bent out of the beam by magnets and the neutral beam proceeds on to the plasma (Fig. 18).

The atomic physics of molecular hydrogen plays a key role in the ion source. H^+ , H_2^+ , and H_3^+ are all present (Raimbault and Girard, 1981). The H_2^+ and H_3^+ are undesirable, and considerable effort has been expended to reduce their concentration. They are formed from $H_2 + e^-$ and $H_2^+ + H_2$ reactions, and the successful approaches for reducing H_2^+ and H_3^+ concentrate on keeping the H^+ in the ion sources away from the walls to reduce the H_2 density and avoiding the production of H_2^+ .

The reason H_2^+ and H_3^+ are undesirable is that they are extracted and accelerated and neutralized along with H^+ . Thus the beam can have an appreciable fraction of neutrals with one-half and one-third of the extraction and acceleration voltage. These lower energy neutral atoms do not penetrate

to the center of the plasma as well as the full energy atoms, and thus are not as effective in heating the plasma.

The fraction of an H^+ (or D^+ or T^+) beam that can be neutralized by charge exchange in a gas cell declines rapidly above 50 keV/amu (Fig. 21), chiefly due to the decline in the charge exchange cross section H^+ (fast) + $H^0 \rightarrow H^0$ (fast) + H^+ compared to the atom impact ionization cross section H^0 (fast) + $H^0 \rightarrow H^+$ (fast) + ?. This renders H^+ based neutral beam systems unfeasible above about 100 keV/amu due to the low conversion efficiency for turning H^+ beams into H^0 beams. It may be possible to recover the power in the unneutralized component of the H^+ beam, or inject it through a second neutralizer cell. Such schemes, however, lead to significant complications in any practical beam system, and have not yet even been demonstrated to work on any large beam system. A fairly large column density of neutral hydrogen (on the order of $10^{16}/\text{cm}^2$) has to be maintained in the neutralizer gas cell. Keeping this neutral gas out of the main vacuum system for the plasma requires extensive high speed cryogenic hydrogen vacuum pumping systems. These difficulties increase as the beam energy is raised and in practice a large neutral beam at energies of 80-100 keV/amu is a very complicated and inefficient system, involving separators, energy recovery, and huge pumping speeds (20 megaliters of liquid He cryopumping for INTOR (INTOR, 1980)) (Fig. 22) (Jassby, 1977). However, neutral beams do have the advantage that they have had the most success in heating tokamaks (Eubank et al., 1979) and mirrors (Coensgen et al., 1980). A summary of the parameters of typical tokamak neutral beam systems is given in Table IV.

However, the complicated operation and low efficiency of H^+ based neutral beams, especially at the energies above the 100 keV desirable for large experiments to achieve good penetration and heating, has led to interest in beam systems based on negative ions. H^- beams can be neutralized with about sixty percent efficiency at high energies in a H_2 gas cell (Fig. 21). The chief difficulty has been the construction of high current H^- sources (10 amps or greater). A wide variety of approaches are under investigation (Sluyters, 1980). The most successful have been the use of double charge exchange of H^+ with alkali and alkaline earth vapors, $H^+ + Na \rightarrow H^- + \dots$, and the conversion, on an alkali coated surface, of H^+ to H^- . Another promising approach is the volume production of H^- in an energetic molecular hydrogen gas (Bacal, 1981). The exact details are not clear but it is thought that the volume production process involves the electron impact dissociation of vibrationally excited H_2 into H^0 and H^- . High energy beams would be useful for mirror machines as well as tokamaks.

Marginal gains in the neutralization efficiency of H^+ beams might be obtained, allowing one to push the maximum feasible energy of such beams up to ~ 120 keV/amu, by the use of He or Ne as a neutralizer gas instead of H_2 (Grisham and Post, 1981a). One gains efficiency at the cost of increasing the pumping difficulty with the use of helium.

The neutralization of H^- by photodetachment has received some study (Sluyters, 1980) but looks marginally acceptable due to high laser powers required, and the relative inefficiency of these lasers.

The search for better beams has led to a new idea, the injection of high energy neutral or singly charged atoms of Li, B, C, O, Na, Si, etc. (Grisham

et al., 1981b; Dawson and MacKenzie, 1980). In the neutral beam approach (Fig. 23), negative ions of Li, C, O, or Si would be accelerated to 1-2 MeV/amu in an RF accelerator, then neutralized by either gas collisions or photo-detachment, and finally injected into the tokamak or mirror experiment (Post et al., 1981a) where they ionize in steps until they are fully stripped. The ions then heat the plasma by coulomb collisions. It appears that it may be feasible to make 500 mA to 2 amp O^- sources, for example, and the necessary RF accelerator technology is being developed for other purposes. The key issues are the formation of the negative ions Li^- , C^- , O^- , and Si^- , the one electron loss cross sections of the negative ions for gas collisions and photodetachment, and the electron loss cross sections for all charge states of the atoms at relative velocities of $\sim 10^9$ cm/sec. Recent experiments (Grisham et al., 1981c; McCullough et al., 1981) have shown that about fifty percent of a 10^9 cm/sec Li^- beam can be neutralized in a gas cell (Fig. 24) but the maximum neutral fractions for C^- , O^- , and Si^- are thirty, twenty-five, and twenty percent. The Brookhaven (Grisham et al., 1981c) results give a hint that the cross sections for the loss of two or more electrons in a single collision may be large. The injection of positive ions depends on their being injected on unconfined orbits and being ionized in the plasma to a higher charge state, thereby shrinking the gyroradius, and changing the unconfined orbit to a confined one. This effect is also important for the neutral light atom beam. The calculations (Grisham, et al., 1981b) (Fig. 25) of this process involve a succession of one electron loss collisions. Significant multi-electron loss cross sections would change the penetration and heating calculations. Also uncertain is the role of two step processes in which an atom or ion is first excited by a collision then ionized by a second collision resulting in a possible enhancement in the ionization rate.

The chief advantage of the light atom neutral beam approach lies in the very large beam energy that could be used (1-2 MeV/amu). The upper limit on the beam energy is set by the requirement that the neutral atoms be ionized in the plasma, and the ions thus formed be confined while they slow down and heat the plasma. The high energy would reduce the current required, thereby making possible a relatively simple DC beam system with small pumping requirements. One oxygen beam could deliver 1 amp at 30 MeV for a total power of 30 MW. With such low currents, the impurity contamination would be smaller (0.1% per sec) than the helium ash buildup. Since the beams are low Z, the central radiation would be low. Normal H^+ beams typically have a 0.5-1% contamination level of oxygen and other impurities. Due to the high energies of the low Z neutral beams compared to 50-100 keV/amu for conventional H^+ based beams, the injected neutral oxygen current per megawatt of beam power is lower for a 1 MeV/amu neutral oxygen beam than for a 75 keV/amu H^+ based H^0 neutral beam.

There are problems of impurity production and hydrogen ion density increases in tokamaks during RF heating. These are probably caused by the RF degrading the confinement of ions at the plasma edge. If loss of confinement is the problem, the poorly confined ions strike the limiter and walls, knocking impurities off and desorbing hydrogen from the walls. While this may be a serious problem, it is not intrinsically different from the general plasma wall interaction problem.

DIAGNOSTICS

The field of plasma diagnostics has made as much or more progress in the last ten to fifteen years as other areas of fusion research. Now, much more than in the past, basic plasma parameters (n_e , T_e , T_i) are measured by a variety of methods. Many of the techniques now can provide spatially resolved and time dependent data during a single discharge (Fig. 26). The plasma parameters quoted for fusion experiments can be relied upon, a marked change from fifteen to twenty years ago. Each important parameter is usually measured by several techniques (Rawls, 1979). A review of diagnostic techniques has been done by the TFR group (DeMichelis, 1978), and I will just summarize the general field and discuss several atomic physics issues. Table V lists many of the quantities measured, the techniques commonly employed, and a few techniques being developed for use in the next generation of tokamaks. I will discuss laser techniques, then neutral particle diagnostics, particle beam techniques, and UV, visible, and x-ray spectroscopy.

The most fundamental parameters in a plasma are T_e , n_e , and T_i , the electron temperature, electron density, and the ion temperature respectively. The basic diagnostic for n_e and T_e uses Thomson scattering. Light from a Q-switched ruby laser is collected after being scattered by 90° by the plasma electrons (Bretz, 1978). The doppler broadening of the scattered light gives the electron temperature and the scattered intensity gives the electron density. The phase shift of a microwave beam, due to the index of refraction for microwaves caused by free electrons, provides another measurement of the electron density. Other interferometry techniques using HCN (Vernon et al., 1977), alcohol (Wolfe, 1976), and CO_2 (Baker and Lee, 1978) lasers have recently been developed primarily for use with high electron density experiments.

Just as Thomson scattering and microwave interferometry are the fundamental diagnostics for T_e and n_e , neutral particle analyzers are the primary tool for measuring the ion temperature. The technique is to use charge exchange of the confined hot ions with hydrogen neutral atoms (either background or introduced by a beam) to neutralize the hot ions and allow them to escape the confining magnetic field. Once outside the plasma, the neutrals can be ionized and analyzed. The ion temperature can then be taken from the slope of the energy spectrum of the hot neutrals.

A second use of neutral particle diagnostics is the measurement of the fast ion velocity distribution produced either from RF heating of the plasma or from the slowing down of the injected fast ions which are introduced during neutral beam heating. These measurements are important since they provide information about the details of how the plasma heating technique is working.

Extensions of this idea to the use of double charge exchange to measure energetic He^{++} distributions have been proposed. The first technique of this type would involve the use of neutral helium beams of 20-100 keV to neutralize the He^{++} ions used in ICRF 3He^{++} minority heating (Post et al., 1981b). In the second technique of this type, the distribution of fast alpha particles would be measured by double charge exchange between the fast alpha particles and a high energy neutral lithium beam (~ 6 MeV) (Post et al., 1981c). The double charge exchange cross section of Li and He^{++} has recently been measured by McCullough et al. (1981) (Fig. 27). Three remaining issues relevant to

this diagnostic technique involve the atomic physics of Li and He. A substantial and known portion of the neutral lithium beam must penetrate into the center of the plasma. The importance of two-step processes (successive collisions between the 6 MeV Li^0 beam and the plasma) leading first to an excited state of Li^0 and then to the ionization of the loosely bound excited state, must be evaluated. The double charge exchange cross section for excited Li^0 and He^{++} has to be evaluated. The excited state distribution of He^0 produced by the double charge exchange must be known, since only the ground state and low lying excited states of He^0 have a good chance of escaping the plasma before being ionized again.

High energy heavy ion beams have also been used to measure the electrostatic potential in the plasma (Jobes and Hosea, 1973). In the ST tokamak, a 100 keV Th^+ beam was injected into the plasma, where some of it was ionized to Th^{++} . By detecting the Th^{++} beam current as a function of position and energy, the electrostatic potential at the point of the second ionization was determined. Similar techniques are used on mirror experiments, where the electrostatic potential is a crucial parameter.

A neutral lithium beam has been used on Pulsator (McCormick et al., 1977) to measure the poloidal field by the Zeeman splitting of the collisionally excited resonance line ($2s - 2p$, 6708 Å). The polarization of the unshifted component of the Zeeman triplet line is used to measure the angle between the toroidal field and the poloidal field, from which the current can be inferred.

Spectroscopy (visible, UV, and x-ray) is one of the major diagnostic tools for fusion plasma research. It is used both to measure the plasma conditions (n_e , T_e , T_i , etc.) and to measure the plasma constituents. Its main use in measuring plasma parameters is to measure the plasma ion temperature by measuring the line broadening due to doppler shifts from the thermal motion of the radiating ions (DeMichelis, 1978). Two new uses and extensions of this technique have recently been made.

The first has been to use the doppler shift to measure the rotation velocity of a tokamak plasma. The PLT plasma has been observed to rotate at speeds of about 10^7 cm/sec (Suckewer, 1981) during high power neutral beam heating when the beams are injected in only one direction. This measurement was done using forbidden lines of FeXX. The second application of spectroscopy has been to use forbidden lines of highly ionized ions to measure the ion temperature in the hot central plasma. It is usually difficult to obtain sufficient resolution for the short wavelength resonance lines (1-100 Å) of the highly stripped ions present in the center of a hot discharge to measure doppler shifts and broadening. However, the use of forbidden lines with wavelengths in the near UV (eg., 2665 Å for FeXX where optical techniques are useful) has made it easier to use doppler broadening measurements in very hot plasmas. A good summary is given by Suckewer (1981). The technique is being extended to higher electron temperatures (3-5 keV) by the use of Kr and Mo injected into the plasma in trace amounts (Cohen, 1981).

The slope of the energy spectrum of the x-ray continuum radiation is often used to measure the electron temperature. The continuum radiation (usually at energies of $2-5 \times T_e$) is due to radiative recombination and bremsstrahlung, and the spectrum has a temperature dependence of (DeMichelis and Mattioli, 1981)

$$\frac{1}{T_e^{1/2}} \exp[-(h\nu/T_e)] .$$

For $h\nu \gg T_e$, the x-ray intensity spectrum gives T_e .

The presence of high energy electrons (runaways) accelerated by the toroidal electric field can be measured by the enhancement of the high energy x-ray spectrum above the recombination and bremsstrahlung spectrum.

Another technique for measuring T_e is to use the intensity ratio of a dielectronic satellite line to the K_{α} -line of helium-like FeXXV using a high resolution crystal spectrometer at 1-2 Å (Bitter et al., 1979). This ratio is proportional to $\exp(E_i/T_e)/T_e$ where E_i is the ionization energy of lithium-like iron. The electron temperature thus deduced is accurate enough to confirm other measurements of the electron temperature.

The major use of spectroscopy, however, has been to measure the constituents of the plasma. The line emission of different impurity charge states allows the measurement of both the relative position and the absolute density of the impurity ions as a function of time. By comparing the observed positions of each impurity charge state with calculated positions assuming various transport models, the plasma transport can be studied (Fig. 28). Deducing the plasma transport rates requires highly accurate ionization, recombination, and excitation rates. It also requires including the effects of charge exchange recombination, since charge exchange recombination can shift the ionization equilibrium.

The measurement of the plasma transport and impurity transport in general is crucial since impurity radiation losses pose a severe threat to fusion reactor experiments. Much of the transport analysis has been complicated by the lack of knowledge of the behavior of the fully stripped impurity. Techniques to measure the fully ionized impurity density are under development. They involve the measurement of the distribution of the excited state radiation of hydrogen-like ions produced by charge exchange of the fully stripped ions and hydrogen neutrals (Isler, 1977; Afrosimov et al., 1978; Isler et al., 1981). The development of this technique requires a good understanding of the excited state distribution of the ion produced by the charge exchange. With a more complete picture of all of the charge states, progress in understanding impurity transport will hopefully follow.

Finally, spectroscopic measurement of a few of the dominant resonance lines of an impurity is used in conjunction with models for all the line transitions to determine the total impurity line radiation energy losses.

Another technique used to measure the energy losses due to radiation and neutral particles is bolometry, which measures the total energy falling on a detector. It provides a total loss rate, which is checked by the spectroscopic measurements.

The ion temperature is also measured by the neutron emission, since the fusion reaction rate is a strong function of T_i ($\propto T_i^{3-4}$ at 1-3 keV).

INERTIAL CONFINEMENT

The invention of lasers and high energy particle beams in the 1960's opened the possibility of compressing small pellets of DT and other fusionable materials to high enough densities and temperatures so that the fusion alphas would be stopped in the pellet and thus ignite the pellet in a very much scaled down version of a hydrogen bomb. The basic idea is that energy (laser, beams) is focused on the pellet surface (Fig. 29). The surface is rapidly heated and ablated. The continued ablation drives the pellet material inward with a spherically convergent shock wave. The central plasma is compressed, hopefully adiabatically. The central pressure ($P_{\text{V}} = \text{const}$) increases until the pressure of the compressed plasma equals the pressure of the driving ablator. Then the pellet center expands and the pressure drops. The ignition conditions are that T exceed a few keV, and ρR be greater than 1 gm/cm^2 . ρR is the product of the the compressed fuel density and pellet radius required for alpha containment. Typical pellet dimensions (consistent with yields less than one ton of TNT) are in the range of 0.01 - 1 cm for the uncompressed pellet and 10^{-1} to 10^{-3} cm respectively for the compressed pellet. Since the temperature must be in the 1-10 keV range for fusion to occur, the thermal velocity of the particles is $\sim 10^8 \text{ cm/sec}$. The pellet will thus fly apart in 10^{-9} - 10^{-10} seconds. This time can roughly be interpreted as a confinement time, τ_c , so $n_e \tau_c \sim 10^{14} \text{ sec/cm}^3$ implies that $n_e \sim 10^{23}$ - $10^{25} / \text{cm}^3$. Thus the conditions for inertial confinement fusion are $T >$ a few keV, $\tau_c \sim 10^{-9}$ sec, and $n_e \sim 10^{24} \text{ cm}^{-3}$.

Atomic physics is somewhat more intimately involved in the physics of inertial confinement than in magnetic confinement, but the issues can still be roughly divided into those affecting the central plasma, the edge plasma, heating and diagnostics (Table VI).

Conditions in the central plasma depend crucially on the equation of the state of the plasma in the pellet and the opacity of the central pellet plasma. The equation of state is given by

$$P = \frac{\rho}{A} (Z^* + 1) T,$$

where P is the pressure, ρ is the mass density, A is the atomic weight of the plasma ion, Z^* is the number of free electrons per ion, and T is the temperature. The key quantity to calculate is Z^* , which is also the effective charge of the ion. The heat capacity also requires Z^* . The other item of interest in calculating the hydrodynamics is the opacity. The transport of thermal x-rays is an important factor in the behavior of the imploding and exploding pellet. Particularly important is line radiation from moderate Z and high Z materials used in multi-layered pellet designs. These data are included as essential ingredients in large hydrodynamic calculations such as LASNEX (Zimmerman and Kruer, 1975).

The high densities ($\sim 10^{24} / \text{cm}^3$) and short times ($\sim 10^{-9}$ sec) encountered in inertial confinement open up a class of atomic phenomena that are different than conventional astrophysical and laboratory plasmas. The interatomic spacing, defined as the ion sphere radius, R_1 , (Weisheit, 1982), where

$$\frac{4\pi R_1^3}{3} = \frac{\langle Z_1 \rangle}{N_e} ; R_1 = \left(\frac{1.61 \times 10^{24} \langle Z_1 \rangle \text{cm}^{-3}}{n_e} \right)^{1/3} a_0 ,$$

where $a_0 = 0.529 \times 10^{-8} \text{cm}$ (the Bohr radius), and $\langle Z_1 \rangle$ is the average number of free electrons per ion, is the order of the size of the orbits for bound electrons. This implies that the pressure of nearby ions will affect the energy levels, collision rates, etc., of an ion. Binary collision models which treat an ion as an isolated system undergoing an impact collision will have to be modified. A measure of the importance of this effect is the ion coupling parameter given by the ratio of the coulomb potential at the ion sphere radius divided by the temperature (Dewitt et al., 1973)

$$\Gamma = \frac{\langle Z_1 e \rangle^2}{R_1 T} = \frac{\langle Z_1 \rangle^{5/3}}{T(\text{eV})} \left(\frac{n_e}{8 \times 10^{19} \text{cm}^{-3}} \right)^{1/3} .$$

When $\Gamma \gg 1$, coulomb interactions control the particle motions, but when $\Gamma \ll 1$ thermal motions are dominate. ICF (inertial confinement fusion) plasmas lie in the $10^{-2} \lesssim \Gamma \lesssim 1$ range (Fig. 30).

The first step in calculating collision rates, line positions, etc., is to calculate the potential a bound electron sees on a partially stripped ion. The contribution due to adjacent ions is usually done with a modified Debye-Huckel model (Weisheit, 1982). The contribution of the bound electrons is typically computed using a Thomas-Fermi type model.

There is appreciable Stark broadening of emission lines due to the transient electric fields produced by the thermal motion of the ions and electrons (Hooper, 1968).

The first major effect of nearby ions is that the ionization continuum is lowered. Electrons of the n^{th} quantum shell (n = principal quantum number) have an average orbit radius

$$r_n = a_0 \frac{n^2}{Z} .$$

The density effects are significant (More, 1982) when $r_n \gtrsim R_1$ or $n_1 \gtrsim 1.6 \times 10^{24} \text{cm}^{-3} (Z^*/n^2)$. Electron wave functions of adjacent ions overlap, leading to a broadening of atomic states into energy bands and "pressure ionization." The energy levels are also displaced due to the screening by "free" electrons of the core potential seen by a bound electron.

Screening affects not only the energy levels but also the eigenfunctions of bound electrons. Thus screening changes the strengths of radiative transitions from the values for isolated ion (Fig. 31). The strength of the screening can be characterized by the Debye length D ,

$$D = \left[4\pi e^2 \sum_i Z_i^2 n_i / T_i \right]^{1/2}$$

$$= \left\{ \sum_i Z_i^2 n_i / (1.97 \times 10^{22} \text{ cm}^{-3} T_i (\text{eV})) \right\}^{-1/2} a_0.$$

In general, the transition strengths tend to decrease rapidly as the shielding length becomes short (Weisheit, 1982).

Screening can significantly reduce the cross sections for electron impact excitation near threshold. For comparison Weisheit (1982) shows the $4s \rightarrow 2p$ excitation cross section of the isoelectronic configurations (Fig. 32) O^+ and N. The energy is scaled by the threshold energy of each transition. The finite range of the N^0 coulomb potential, illustrating the effects of screening, causes the nitrogen cross section to vanish near the threshold whereas the O^+ potential is long range and the cross section rises steeply near the threshold. The same effect occurs for the $1s \rightarrow 2s$ transition in one electron atoms as the screening is increased.

Dielectronic recombination rates are also reduced due to the removal of high n orbitals as illustrated by Fig. 33 which uses unscreened radiative and collisional rate coefficients (Weisheit, 1975).

The edge plasma is extremely important to inertial confinement since the energy needed to compress the pellet is deposited there. Typical densities run from $10^{14} - 10^{22} \text{ cm}^{-3}$ and typical temperatures run from 10 eV to 1 keV. The key issues are the absorption and stopping of laser light and the absorption and stopping of fast ions and electrons. Magnetic fields in the megagauss range are generated and can affect the collision rates. As in the magnetic fusion case, the interaction of the pellet debris with the reactor vessel is important. The emphasis is on the survival of wall to pulsed heat loads and particle loads (about 1,000 joules/cm² in 10^{-9} seconds).

The exact mechanisms for the absorption of laser light are not understood. The light propagates through the low density plasma at the pellet edge until it reaches the "critical" density where the light frequency equals the plasma frequency $\nu_p = 8.97 \times 10^3 n_e^{1/2} / \text{sec}$ (Spitzer, 1962) (about 10^{19} cm^{-3} for CO_2 lasers, about 10^{21} cm^{-3} for Nd glass).

Ion beams are stopped by collisions with the partially ionized ions and electrons in the pellet. Most experimental data has been gathered only for cold materials. For cold materials the stopping is due primarily to ionization and excitation of bound electrons. Stopping power calculations involve adding the contributions of the bound electrons of the partially ionized atoms and the stopping due to the free electrons (Hauer, 1981). A portion of the absorbed laser energy gives rise to a high energy electron population. As with heavy ions, fast electrons are stopped by impact ionization and excitation collisions with bound electrons, and scattering from free electrons.

The main approaches to inertial confinement are defined by the type of beam used to compress the pellets (lasers, light ions, heavy ions, and electrons). The original approach was to use high energy lasers (Nuckolls et

al., 1972). The primary difficulties with lasers are their low electrical efficiency and poor coupling of the laser power to the pellet. A facet of the coupling problem is the generation of high energy electrons which penetrate to the pellet core and heat it before the final compression stage, thus increasing the power required to reach the required high density. Current experiments indicate that the laser-plasma coupling is improved, and the hot electron preheat is reduced as the laser wavelength is reduced. Large CO₂ lasers (10 μ) are being used at the Los Alamos National Laboratory. Neodymium glass lasers (1 μ) are the most common, with programs at Lawrence Livermore National Laboratory, KMS, Inc., and the University of Rochester, in the United States, at the Lebedev Institute in USSR, at Osaka University in Japan, at LeMeil in France, and at the Max Planck Institute in Garching, West Germany. There is intense interest in going to shorter wavelengths using exotic lasers such as KrF lasers or by frequency doubling. Livermore is constructing a frequency doubled laser ($\sim 0.5\mu$) beginning with a Neodymium glass laser ("Novette").

Given the apparent simplicity of the stopping of high energy ions in a plasma, the short range of the ions, and the possibility of high electrical efficiency for ion beams, there is an active program to use both light and heavy ion beams for compressing pellets. The light ion beam research is being done in the United States at Sandia Laboratories in Albuquerque, New Mexico, and at the Kurchatov Institute in Moscow. The heavy ion beam fusion work is being done at the Argonne National Laboratory, and the Lawrence Berkeley Laboratory in the United States, and the Rutherford Laboratory in the United Kingdom.

One candidate beam is ~ 1 GeV Cs⁺. Recent work on the reaction Cs⁺ + Cs⁺ + Cs⁰ + Cs⁺², Cs⁺ + Cs⁺⁺, etc., has had a bearing on the feasibility of such beams, since these reaction could destroy the beam.

Since lasers involve electronic transitions in molecules and atoms, atomic processes will continue to play a large role in the search for high power, efficient lasers. The thrust toward shorter wavelengths ($< 1\mu$) will be an important research area. Significant improvements in both these areas will have to be made for lasers to be feasible as drivers for an inertial confinement fusion reactor. At this moment ion beams seem more promising for reactor applications than lasers. Thus there will be an increasing emphasis on the production and acceleration of efficient, high energy, and high power ion beams.

Diagnostic techniques for inertial confinement experiments rely heavily on atomic processes. Spectroscopy in the UV, x-ray, and visible is used (Yaakobi et al., 1977). One common technique is to seed the DT gas with neon or argon and look at the emission lines for line profile changes, line shifts, and other high density effects (Bristow et al., 1979). Many pellets have layers of high Z materials, and time and space resolved line radiation, and x-ray continua from these high Z layers allows one to determine the time history of the temperature and density of the layers in the compressing and exploding pellet. Spectral information coupled with rate equation modeling of ionization and recombination in the strong temperature and density gradients at the ablation layer where the laser is deposited, provides a measure of the laser plasma interaction. Optical laser probes are used with interferometry, faraday rotation, and polarization to measure the density profiles in the

ablating surface of the pellet.

One example of this is the use of K_{α} radiation produced by the impact excitation of fast electrons as a measure of the power absorbed which goes into the production of fast electrons. The simple ratio of the total stopping power to the stopping power due to K_{α} excitation is given by (Hauer, 1981).

$$\frac{dE/dx|_{\text{tot}}}{dE/dx|_{K_{\alpha}}} = C \frac{\ln(4E/I)}{\ln(E/E_{K_{\alpha}})},$$

where E is initial electron energy, I is the average bound excitation energy, and $E_{K_{\alpha}}$ is the K_{α} excitation energy.

This simple analysis is complicated by the heating and stripping of the ions being excited and the subsequent shifts of the emission lines (for A^0 , A^{+1} , A^{+2} , A^{+3} , etc.). Resonance line absorption into K shell excitations, the contribution of the free electrons to the stopping power, and the variation of the ratio of K_{α} line emission to Auger emission with charge state all further complicate the analysis.

Bremsstrahlung radiation is also used as a tool to measure the hot electron temperature and transport using detailed bremsstrahlung cross sections (Pratt et al., 1977). Extensive modeling of these two emissions and their interaction with the plasma is required to recover quantitative information about the behavior of the fast electrons.

The best conditions achieved so far by the inertial confinement program are $n \sim 3 \times 10^{24} \text{ cm}^{-3}$, $T \sim 1 \text{ keV}$, and $\tau_c \sim 1.5 \text{ nsec}$ (Hauer, 1981). Pellets of initial diameters of two-hundred to four-hundred microns have been compressed to ten to twenty microns. The major funding commitment is for large laser systems at Livermore and Los Alamos in the United States (Table VII). The United States budget is \$215 million/year.

THE ROLE OF PLASMA PHYSICS IN ATOMIC PHYSICS

The programmatic needs of the fusion community have provided motivation and resources for many of the current research topics in atomic physics. The interest in charge exchange of multi-charged ions with hydrogen was spurred on by the concerns of the fusion community that neutral beams would not penetrate plasmas with even relatively small impurity levels. In addition, fusion research has, as a by-product, produced new experimental and theoretical tools to explore atomic physics in new regimes. Plasmas with temperatures as high as 7 keV and densities up to $3 \times 10^{24} \text{ cm}^{-3}$ are produced in laboratories. These have led to the identification of the energy levels of highly stripped ions (such as W^{+50}) (Reader and Luther, 1980). Large spectroscopy groups have been set up at the major fusion laboratories with the attendant improvements in spectrometers and detector technology. These groups have observed effects before seen only in astrophysical plasmas before, such as dielectronic satellite lines and recombination edges. The work on excitation-autoionization is of great interest to the fusion community. The work on the atomic physics of dense matter is a qualitatively new area that will almost certainly grow in interest and importance.

A lot of extremely useful work is being done by the various data center groups throughout the world to assemble the data relevant to fusion and put it in a digestible form so that plasma physicists can easily use the data. In the United States there are data centers at the Oak Ridge National Laboratory in Oak Ridge, Tennessee; at the Joint Institute Laboratory for Astrophysics in Boulder, Colorado, and at the National Bureau of Standards in Washington, D. C. In Japan there are data centers at the Plasma Physics Institute at Nagoya University in Nagoya, and at the Japanese Atomic Energy Research Institute at Tokai. In Europe there are groups at the International Atomic Energy Agency in Vienna, Austria; at Queen's University, Belfast, Northern Ireland, and at the University of Paris in France. In the USSR there is a group at the Kurchatov Institute in Moscow.

REVIEWS

Good textbooks and review articles are essential for people outside a field to efficiently learn what they need to know about the field. I have provided a short key to textbooks and reviews in Table VIII.

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TABLE I (Rawls, 1979)

Design Parameters and "Typical" Plasma Properties for Major U.S. Tokamaks

Tokamak	a (cm)	R ₀ (cm)	B _T ^a (kG)	I _p ^a (kA)	T _e (0) ^b (eV)	T _i (0) ^b (eV)	n _e ^b (cm ⁻³)	τ _E ^b (ms)
ST	13	109	50(40)	130(70)	2500	600	4 × 10 ¹³	10
ATC ^c	17	90	20(15)	100(60)	1100	250	1.5 × 10 ¹³	5
ORMAK	23	80	25	340(170)	1500 ^d	1800 ^d	3 × 10 ¹³	10
ALCATOR A	10	54	120(80)	400(200)	1200	800	6 × 10 ¹⁴	20
DOUBLET II-A	14 × 40	66	10 ¹ 8	320(150)	400	200	3 × 10 ¹³	4
PLT	45	130	45(32)	1400(550)	4000 ^d	7000 ^d	8 × 10 ¹³	100
ISX	25	90	18(12)	160(110)	1300 ^d	1800 ^d	6 × 10 ¹³	20
DOUBLET III	45 × 150	143	40(20)	5000(2000)	1000	400	1 × 10 ¹⁴	60
ALCATOR C	17	64	140(70)	1000(400)	900	700	5 × 10 ¹⁴	30
PDX	45	145	24(24)	500(500)	3000 ^d	6000 ^d	1 × 10 ¹⁴	80

^aDesign values (normal operating values in parentheses).

^bMaximum obtained - not necessarily in same discharge.

^cBefore compression.

^dWith neutral beam heating.

TABLE II

Examples of the United States Tokamak System

	PLT	+	TFTR	+	FED(estimated)
R (major radius)	140 cm		250 cm		520 cm
a (minor radius)	40 cm		85 cm		130 cm
Date	1978		1984		1991
Cost (millions of 1981 dollars)	30		350		1500

TABLE III
Reactions Used in Divertor Model (Heifetz, 1981)

TABLE IV
Neutral Beam Systems

Energy	Ion	Power	Experiment	Date
40 keV	D ⁰ , H ⁰	3 MW	PLT	(1978)
50 keV	D ⁰ , H ⁰	8 MW	PDX	(1981)
80 keV	D ⁰ , H ⁰	7 MW	D-III	(1983)
120 keV	D ⁰	20 MW	TFTR	(1984)
150 keV	D ⁰	60 MW	FED	(1990)

TABLE V
Plasma Diagnostics

Plasma	Quantity	Present Techniques
Electron density profile	$n_e(r)$	Thomson scattering, microwave interferometry, laser interferometry.
Electron temperature profile	$T_e(r)$	Thomson scattering, cyclotron radiation, x-ray spectra.
Ion temperature profile	$T_i(r)$	Charge-exchange spectra, doppler broadening of impurity lines including forbidden lines, neutron emission.
Ion density profile (hydrogen and impurities)	$n^q(r)$	Line spectra (visible, UV, x-ray), charge-exchange recombination.
Neutral beam ion distribution	$f_{\text{beam}}(v)$	Charge-exchange spectra.
Electron velocity distribution	$f_e(v)$	x-ray
Plasma current profile	$j(r)$	Orbit shift of injected beam, Zeeman splitting, Faraday rotation.
Magnetic instabilities	MHD	Loops, x-ray
Plasma turbulence	Fluctuations	Probes, CO_2 laser scattering, microwave scattering
Radiation losses	$P_{\text{rad}}(r)$	Bolometers, spectroscopy
Edge plasma, plasma-wall interaction		Electric probes, ion doppler broadening hydrogen light, surface analysis.
Electrostatic potential	$\phi(r)$	Ionization of heavy probe beams
Sample future work would better versions of above plus:		
Ion temperature profiles	$T_i(r)$	Forbidden lines of injected impurities for $T_e > 3 \text{ keV}$.
Impurity distribution	$n^{+q}(r)$	Multichannel spectrometers
Fast alpha distribution	$f_\alpha(v)$	Double charge-exchange with neutral Li beam

TABLE VI
Inertial Confinement Atomic Physics Issues
(adapted from (Hauer, 1981))

ICF Related Process	Atomic Processes Involved	Theory and Modeling	Diagnostic Techniques
Absorption of laser light	Ionization and re-combination dynamics in presence of large density gradients and electric fields	Rate equation models coupled to hydro codes or other plasma modeling ionic species and energy spectrum	(1) Spatially resolved spectral line observation (2) Charged particle
Transport and deposition of fast ions and electrons	(1) Collisional ionization and excitation (2) bremsstrahlung (3) electron scattering	Stopping power formulations needed for input to comprehensive hydro modeling	Inner shell and bremsstrahlung radiation used as indicator of particle desposition
Thermal, particle, and radiation transport	(1) Bound contributions to absorption (2) thermal and electrical conduction in partially ionized plasma	(1) Opacity calculations needed for radiation treatment in comprehensive models (2) Calculations of transport coefficients	(1) Line radiation from layered targets as indicator of thermal transport (2) Absorption spectra
Production of high density imploded	(1) Quantum effects in the equation of state (EOS) of dense plasmas (2) line radiation from higher Z plasmas	(1) Thomas Fermi Dirac EOS as starting point (2) spectral modeling with rate equations coupled to hydro codes	(1) Observation in dense plasma of (a) line profiles (b) line shifts (c) continuum edge (2) high intensity shock studies

TABLE VII

Large United States Inertial Confinement Experiments

Name	Type	Wavelength	Site	Date	Power	Energy
Helios	CO ₂	10μ	LASL	1978	30 TW	6 kJ
Shiva	Ni glass	1μ	LLL	1978	30 TW	6 kJ
Antares	CO ₂	10μ	LASL	1982	50 TW	40 kJ
Nova	Ni glass	1μ	LLL	1985	60 TW	30 kJ

TABLE VIII

Textbooks and Review Articles

Topic	Reference
Plasma Physics	Chen, (1974); Spitzer, (1962)
Tokamaks	Rawls, (1979); Furth, (1975)
Mirrors	Cohen, (1980)
Impurity Atomic Physics	DeMichelis et al. (1981); Suckewer, (1981)
Plasma Wall Interactions	McCracken, (1979)
Neutral Beam Heating	Jassby, (1977); Raimbrult, (1981)
Tokamak Diagnostics	DeMichelis, (1978); Rawls, (1979)
Dense Matter Atomic Physics	Weisheit, (1982); More, (1982)
Inertial Confinement	Hauer, (1981); Miley, (1979)
Atomic Data Needs of Fusion	Physica Scripta (1981)

FIGURE CAPTIONS

- Fig. 1. Reaction rates for a variety of fusion reactors averaged over a Maxwellian ion temperature distribution (Greene, 1967). (762346)
- Fig. 2. $n\tau_E$, the product of the plasma density and energy confinement time for ignition and energy breakeven (Lawson criteria) for DT and DD plasmas (Jassby and Towner). (773873)
- Fig. 3. Schematic drawing of charged particle confinement by a magnetic field. (786452).
- Fig. 4. Schematic illustration of a tokamak showing the toroidal and poloidal magnetic fields. (754023)
- Fig. 5. Schematic diagram of energy flows in a tokamak. The center of the torus is on the left. (782058)
- Fig. 6. Charge state distribution as a function of electron temperature for tungsten in coronal equilibrium (Post, et al., 1979). (782223)
- Fig. 7. The temperature range for which helium-like ions of atomic number Z are the dominate charge state. For a given Z, hydrogen-like and fully stripped ions predominate at temperatures above the band. Ions with three or more bound electrons are dominate below the band. Only ions below the band can radiate strongly. (792394)
- Fig. 8. Emissivities of various elements in coronal equilibrium (Jensen, et al., 1977). (772182)
- Fig. 9. Hollow electron temperature profiles in the PLT tokamak caused by large amounts of radiation from tungsten ions (Hsuan, et al., 1978). (786412)
- Fig.10. Ionic species abundances in "coronal" equilibrium modified by beam neutrals for iron as a function of electron temperature. Results are shown for four values of the neutral fraction n_n/n_e of 20 keV/amu hydrogen, using the Olson-Salop (Olson and Salop, 1977) classical trajectory (OSCT) cross sections (Hulse, et al., 1980). (792652, 792651)
- Fig.11. Iron neutral-beam modified "coronal" equilibrium. Hydrogen neutrals at 20 keV/amu are considered, using the Olson-Salop (Olson and Salop, 1977) classical trajectory cross sections. The curves are parameterized by the neutral fraction n_n/n_e . Shown are (a) the average charge state $\langle Z \rangle$, and (b) the overall radiation rate coefficient P/n_en_n , both as functions of the electron temperature (Hulse, et al., 1980). (792653)
- Fig.12. The $n_en_e\tau_E$ due to plasma transport losses required for ignition as a function of the impurity fraction n_z/n_i for four different impurities and two temperatures. n_i is the total ion density (Jensen, et al., 1977). (772101)
- Fig.13. Maximum allowed impurity concentration versus Z for ignition at various plasma temperatures $T_e = T_i$, assuming zero non-radiative losses (Jensen, et al., 1978). (773294)
- Fig.14. Schematic diagram of the particle flows in a tokamak. The center of the torus is on the left. (782057)
- Fig.15. The magnetic field geometry for a poloidal divertor. A null in the poloidal field is created by a coil carrying a current parallel to the plasma current. Field lines outside this null point encircle both the main plasma and the divertor coil. The plasma edge is thus "diverted" to a region away from the main plasma. In addition to the fields shown here, there is also a strong toroidal magnetic field. (803288)

- Fig.16. Schematic representation of the operation of a poloidal divertor. In addition to the fields shown here, there is also a strong toroidal magnetic field. (81P0177)
- Fig.17. Equilibrium fraction of H^0 and H^+ produced from $H_2 + e^- \rightarrow H^0, H^+$ following molecular ionization and dissociation channels (Heifetz, et al., 1981). (802112)
- Fig.18. Neutral injection by the acceleration of H^+ and neutralization by charge exchange, and the subsequent stopping of the beam in the plasma by coulomb scattering and nuclear reactions (Jassby, 1977). (753441)
- Fig.19. Reaction rates for atomic hydrogen for electron impact ionization, proton impact ionization, and charge exchange with protons as a function of lab hydrogen energy. The energy scale for the electron ionization rate is the electron temperature (Freeman and Jones, 1974). (762150)
- Fig.20. Effective trapping cross section versus neutral hydrogen beam energy, defined as

$$\sigma_{eff} = \frac{Z-Z_{eff}}{Z-1} (\sigma_{cx}^H + \sigma_{p-ion}^H) + \frac{\langle \sigma V \rangle_{e-ion}}{Z(Z-1)} (\sigma_{cx}^{imp} + \sigma_{ion}^{imp}) ,$$

for a deuterium neutral beam traversing a deuterium plasma with an impurity of charge Z . σ_{cx}^H is the hydrogen charge exchange cross section, σ_{p-ion}^H is the proton impact ionization cross section, σ_{cx}^{imp} is the impurity-hydrogen charge exchange cross section, and σ_{ion}^{imp} is the ion impact ionization cross section. The cross sections are taken from (Freeman and Jones, 1974) and (Olson, et al., 1978). (792225)

- Fig.21. Equilibrium neutralization efficiency for D^- , D^+ , D_2^+ , D_3^+ in a gas cell of H_2 as a function of the energy per deuteron (Jassby, 1977). (763705)
- Fig.22. Components of (a) ideal and (b) practical neutral-beam injector systems for a tokamak reactor (Jassby, 1977). (763692)
- Fig.23. Outline of a high energy neutral injection system based on negative ions (Grisham, et al., 1981b). (802298)
- Fig.24. The charge state fraction of a 7 MeV Li^+ beam after passing through a CO_2 gas cell as a function of pressure in the gas cell for an experiment at the Brookhaven National Laboratory (Grisham et al., 1981c). (81P0121)
- Fig.25. Gyro-orbits for a 32 MeV neutral oxygen atom injected into a large plasma. The view is a cross section of a tokamak. The neutral oxygen atom is ionized at x, and successively ionized to +8 (we only show +1, +2, +3, and +4) while it drifts toward the plasma center (Grisham, et al., 1981b). (81X0685)
- Fig.26. Time dependent electron temperature and density profiles in PLT measured by Thomson scattering (Bretz, et al., 1978).
- Fig.27. Cross sections for charge exchange of He^{++} and Li^0 as a function of energy (McCullough, et al., 1981).
- Fig.28. Observed iron charge state distribution in PLT compared with coronal equilibrium distribution (Hinnov) and (Stodiek, et al., 1980). (803942)

- Fig.29. Schematic illustration of an inertial confinement pellet. (81P0191)
- Fig.30. Density and temperature regimes of some natural and man-made plasmas. The two dashed lines show (N_e, θ) -values corresponding to different constant values of the Coulomb parameter Γ for a hydrogen plasma (taken from Weisheit, 1982). (802419)
- Fig.31. The ratio S_{DH}/S of Debye-Huckel-screened to unscreened transition strengths for the principal series in a hydrogenic ion of nuclear charge $(Z + 1)e$ for various normalized screening lengths $(Z + 1)D/a_0$. As an aid to the eye, curves have been drawn to connect values for a given normalized screening length (taken from Weisheit, 1982). (802418)
- Fig.32. (a) The computed electron impact excitation cross section [108] for the transition $(2p^3)^4S \rightarrow (2p^3)^2P$ in N and O^+ . When the nitrogen cross section is scaled to agree at high energies with the oxygen cross section, the dashed curve shows its low energy behavior. For nitrogen, the threshold energy is $\Delta E = 3.58$ eV; for oxygen, it is $\Delta E = 5.02$ eV. (b) The computed electron impact excitation cross section for the transition $1s \rightarrow 2s$ in a one-electron ion of nuclear charge $(Z + 1)e$. The Debye-screened ion results are the Born calculations of Hatton et al., (1981), and the unscreened ion result ($D \rightarrow \infty$) is the Coulomb-Born calculation of Oh, et al. (1978). The threshold energy is $\Delta E = 10.2 (Z + 1)^2$ eV (taken from Weisheit, 1982). (802423)
- Fig.33. Rate coefficients $\{K_9(n) [cm^3 sec^{-1}]\}$ for dielectric recombination to excited states $(1snl)$ of Ne^{+8} in a 100 eV plasma. The logarithm of the plasma electron density $N_e (cm^{-3})$ is shown by each curve; at densities $N_e \geq 10^{20} cm^{-3}$, continuum lowering removes all but the few states indicated (taken from Weisheit, 1982). (81T0028)

Fig. 1

Fig. 2

Fig. 3

Fig. 4. PPL-754023

TOKAMAK ENERGY FLOWS

Fig. 5. PPL-782058

Fig. 9. PPL-786412

Fig. 10a. PPL-792651

Fig. 10b. PPL-792652

Fig. 11

Fig. 12

Fig. 13

Fig. 14

POLOIDAL DIVERTOR MAGNETICS

Fig. 15

Fig. 16

Fig. 17

Fig. 18

Fig. 19

Fig. 20

Fig. 21

(a) IDEAL

(b) PRACTICAL

Fig. 22

Fig. 23

Fig. 24

81X0685

Fig. 25.

Fig. 26

Fig. 27

Fig. 28

81P 0191

Fig. 29

Fig. 30

Fig. 31

Fig. 32. PPL-802423

Fig. 33.