Recent progress in magnetically confined plasma research has brought the goal of controlled thermonuclear fusion within reach. Simultaneously, the increased size, temperature, and density of these devices has resulted in rapid changes in diagnostic techniques. In this article, recent developments in instrumentation are reviewed together with trends for the future. The topics discussed include far-infrared laser interferometry and polarimetry, ruby laser television Thomson scattering, excimer and Nd-glass laser scattering, ion temperature determination via large-$\alpha$ scattering with FIR and $\text{CO}_2$ lasers, collective scattering, Schottky diode mixer technology, synchrotron radiation diagnostics and imaging, ion beam probes, x-ray diagnostics and imaging, neutron diagnostics, resonance fluorescence scattering, ultraviolet diagnostics, and internal magnetic field measurement.

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INTRODUCTION

There is presently considerable worldwide interest in the study of high-temperature plasmas. Much of this interest is driven by the goal of achieving controlled thermonuclear fusion which involves the release of net nuclear energy from the fusion of light nuclei such as deuterium and tritium. In order that net energy be released in the process, several criteria must be met. First, the nuclei must collide with sufficient relative velocity to overcome the repulsive Coulomb barrier. The required thermal energy corresponds to $\simeq 10 \text{ keV}$ for D-T reactions and rises to $\simeq 300 \text{ keV}$ for alternate fuels such as $p^+\text{D}$. In addition to the temperature requirement, both the density and confinement time must be sufficiently high so that the reactions have time to occur within the lifetime (or disassembly time) of the plasma.

The above considerations have led to two parallel approaches, the so-called magnetic confinement and inertial confinement concepts. In the former, the plasma confinement is dependent upon an externally imposed magnetic field. Since economic and technological considerations limit these field strengths to $\simeq 50-100 \text{ kG}$, magnetically confined plasmas tend to be relatively tenuous with densities of $\simeq 10^{13}-10^{15} \text{ cm}^{-3}$. The required confinement time for breakeven in the case of D-T fuel is on the order of a second. Plasma dimensions are $\simeq 1-10 \text{ m}$. In the case of inertial confinement fusion, a small target ($\simeq 0.1-5 \text{ mm}$) is ablative heated and compressed to high density ($\simeq 10^{20}-10^{25} \text{ cm}^{-3}$). There is no actual confinement and the fusion reactions must take place on the time scale during which the target's inertia prevents its disassembly. Relevant timescales are $\simeq 0.03-5 \text{ ns}$.

It is obvious from the above discussion that there is a great disparity in the densities, sizes, and time scales associated with the two fusion schemes. It is, therefore, not surpr
prising that the measurement techniques required for magnetic and inertial confinement fusion differ greatly. Consequently, to attempt to discuss diagnostic instrumentation for both magnetic and inertial confinement fusion in a single review is an impossible task. The present review article will, therefore, be restricted to a review of magnetic fusion diagnostics. A companion article on inertial confinement diagnostics written by E. M. Campbell and R. L. Kauffman is being prepared for publication in a subsequent issue of the Review of Scientific Instruments. Furthermore, since there are so many excellent references available, the present article will concentrate on the current limits of the instrumentation and technology together with future development directions.

Before proceeding to a discussion of magnetic fusion plasma diagnostics, it will prove helpful to consider the various confinement schemes since they impact the measurements. Unfortunately, space does not permit a discussion of all of the devices presently under investigation throughout the world. Therefore, we will restrict our attention to the tokamak and magnetic mirror devices which presently constitute the major efforts as well as representing the so-called closed or toroidal and open geometries, respectively. The interested reader is referred to references such as those edited by Teller for details on other devices.

A schematic of the Princeton PLT tokamak is shown in Fig. 1. The plasma is confined in a toroidal vacuum vessel with the primary field \( B_0 \) produced by the toroidal field coils. The Ohmic heating primary windings indicated in Fig. 1. induce a toroidal current in the plasma which acts as a single-turn secondary. The resulting poloidal field \( B_\theta \) provides the plasma equilibrium. The toroidal current also provides joule heating of the plasma—hence, the term Ohmic heating current. However, because of the \( T^{-1/2} \) dependence of the collision frequency, supplementary particle or wave heating is generally thought to be required to achieve ignition. The first diagnostics problem is simply access. Figure 2 illustrates this problem with a recent photograph of the PLT device. The heating and diagnostic devices crowded around the machine compete for the relatively small number of ports dictated by the need to provide the proper magnetic field geometry. The next generation of large tokamak devices are intended to demonstrate the scientific feasibility of fusion. These include the Princeton Tokamak Fusion Test Reactor (TFTR) device (which began operation at the end of 1982), the Joint European Torus (JET) tokamak at Culham Laboratory in England, and the JT-60 device in Japan. These devices have major radii of 2.65–3 m, mean minor radii of 0.85–1.6 m, toroidal fields of 27–52 kG, plasma currents of 2.5–4.8 MA with flattop times of 1.6–20 s.

The second fusion scheme mentioned earlier is the magnetic mirror device. This is an example of a so-called open-confinement scheme. Here, the term refers to the fact that field lines leave the confinement region. The magnetic mirror device employs the fact that charged particles with sufficiently large ratio of perpendicular to parallel velocity are reflected as they move into regions of increasing magnetic field. This was one of the first approaches investigated for controlled fusion. However, it was soon found that a simple mirror configuration, as produced by two coils (in which the field decreases radially outward from the axis), is magnetohydrodynamically or MHD unstable. One solution to this problem is shown in Fig. 3 which is a schematic of the Lawrence Livermore Laboratory tandem mirror experiment (TMX). The main body of the plasma is contained in the long central solenoid section. Within the baseball coil regions is a minimum \(- B \) magnetic field that contains a high-temperature plasma maintained and heated by neutral beam injection. The next generation of mirror device will be the Lawrence Livermore National Laboratory Mirror Fusion Test Facility B (MFTF-B). As will be seen later, the large size of this device (the vacuum chamber is approximately 64 m long and 6 m in diameter) presents the diagnostician with...
some significant problems. The axial profiles for one mode of operation of the confining magnetic field, plasma potential, and density are shown in Fig. 4. It is clear that the diagnostics problem is complicated (with respect to the tokamak) because of these strong axial variations.

From the above, it should be obvious that apart from the actual diagnostic instrumentation, there is a need for computer control, data acquisition, processing, and display. The complexity of these systems and the data handling requirements are enormous. For example, the basic Princeton TFTR system employs $\approx 1500$ separate CAMAC modules with a total requirement of $\approx 10$ megawords. However, this is a separate field from the actual diagnostic instrumentation and as such is beyond the purview of the present review.

I. INTERFEROMETRY

The most widely adopted method to determine electron density in fusion plasmas is optical interferometry. A Mach–Zehnder interferometer is typically utilized. Plasma occupies one arm of the interferometer and, since the refractive index of the plasma is density dependent, a measurement of the resultant phase shift allows the determination of electron density.

**TMX MAGNET AND NEUTRAL BEAM CONFIGURATION**

![Diagram of TMX magnet and neutral beam configuration](Image)

**FIG. 3.** Schematic of the Lawrence Livermore National Laboratory tandem mirror experiment (TMX).
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The phase shift is given by

\[ \phi = \frac{\pi}{\lambda_0 n_e z_{fL}} \int_{z_1}^{z_2} n(x) \, dz, \tag{1} \]

where \( \lambda_0 \) is the wavelength of the probe source, \( n_e \) is the electron density, \( z_2 - z_1 \) is the path length inside the plasma, the density \( n_e \) is assumed much less than critical \( n_s \), and cgs units have been adopted.

From the above, it can be seen that the integrated line density is readily obtained from a measurement of the resultant phase shift. However, simple interferometers possess some unavoidable drawbacks regarding accurate measurement of this phase shift. The low-frequency difference signal from the mixer is proportional to \( \cos \phi \), where \( \phi \) is the phase shift, and so it is possible to determine whether phase is increasing or decreasing. In addition, the signal is linearly dependent on the product of the amplitudes of the reference and probe beams and so variations in this quantity can be wrongly interpreted as a time varying phase shift. The above difficulties can be easily circumvented using phase-modulation techniques.\(^{15,25-33}\) An example is illustrated in Fig. 5 where the realization of such a system in the microwave region is depicted.\(^{14}\) The single-sideband generator shifts the source frequency \( f_0 \) to \( f_0 + f_m \). This signal is further shifted in frequency (phase modulated) as it traverses the plasma column. The signal is then down converted to \( f_m + (df / dt) \) by the mixer and enters the digital phase comparator together with a reference signal of \( f_m \). The digital phase comparator basically consists of two digital counters; one counting upward and the other downward. The outputs are processed by a digital adder, fed to a D/A converter which provides a signal proportional to the phase shift instead of \( \cos \phi \). Although relatively simple to implement in the microwave region, the single-sideband generator is not trivially extended to the submillimeter and infrared spectral regions. Basically, there are two different techniques presently being commonly employed to generate the frequency off-set beam\(^{15}\): (a) doppler shifts and (b) separate lasers. Two variations of the first approach have been utilized. Veron\(^{15,26}\) reflects a portion of the far-infrared laser output from a rotating grating to generate the frequency shift. Alternatively, one may reflect a beam from a moving reflector to generate the frequency shift.\(^{25-38}\) Although modulation frequencies up to \( \pm 1 \) MHz can be obtained using these mechanical modulation schemes, it is of interest to investigate alternative techniques. The approach first reported by Wolfe \( et \) \( al.\)\(^{27}\) and now used in laboratories throughout the world\(^{39,40}\) is to employ two optically pumped molecular lasers which are tuned to slightly different frequencies (within the gain curve of the medium). Frequency shifts up to several megahertz are easily obtainable with the frequency shift \( \Delta f / f_0 \) given by \( \Delta f = -f_0 / L \), where \( f_0 \) is the laser frequency, \( L \) the cavity length, and \( \Delta L / L \) the fractional difference in cavity lengths. An example of such a system\(^{39}\) is shown in Fig. 6.

The above statements on laser modulators apply primarily to far-infrared lasers. For example, in the case of a CO\(_2\) laser interferometer \( (\lambda_0 \approx 10 \mu m) \) or a He–Ne system \( (\lambda_0 = 3.3 \mu m) \), one has other possibilities. In this case, the most common technique is to employ an acousto-optic cell\(^{41-43}\) to produce the modulated beam(s).

Thus far, we have not addressed the question of how one chooses the source wavelength for a given interferometer system. Here let us quickly address a few of the important issues. First, one obviously wishes to operate below cutoff so
that the beam may pass through the plasma. This sets an upper bound on the operating wavelength. However, a more important constraint is imposed by the beam deflection (refraction) resulting from the density gradient present in any real plasma. In the case of a parabolic density profile, the maximum angular deflection is

\[ \alpha_d = \sin^{-1}\left(\frac{n_r}{n_o}\right) \approx n_r / n_o \propto n_0 \frac{\lambda^2}{\Delta \lambda} \]

This must obviously be small enough that the deflected beam shall pass through the exit window. The machine propagation distances and window dimensions also impose restrictions on the operating wavelength. Beam diffraction determines the minimum achievable spot size and, therefore, spatial resolution requires short wavelength operation. Small spot size also reduces measurement errors due to phase averaging over the beam crosssection. All of the above constraints argue for operation at as short a wavelength as possible. However, there are several additional considerations which argue for longer wavelengths. First, vibrations can change the path length on the time scale of the density measurement. For accurate measurements one wishes the vibrational fringe shift to be small compared with the plasma fringe shift. Requiring an error of less than 1% results in the constraint

\[ \lambda_d(\mu m) > 4.09 \times 10^6 \frac{\bar{D}}{r_0(c m)} n_o(c m^{-3})^{1/2} \]

where \( r_0 \) is the plasma radius and \( \bar{D} / l \) is the vibration-induced path length change. This constant may be relaxed by employing an extremely short wavelength second probe source which will essentially only monitor the vibrations since the plasma shift \( \phi \propto \lambda_d. \) In the case of the TFTR far-infrared laser interferometer system, an extremely safe approach has been taken. Here, vibration compensation has been combined with the choice of a relatively long probe wavelength (119 \( \mu m \)), as well as a vibration isolation frame for the entire interferometer and optical system. An additional constraint arises from fringe counting difficulties. Specifically, a practical lower limit for fringe counting is approximately a twentieth of a fringe. Combining all of the above constraints, it is somewhat surprising that one actually obtains an operating window in wavelength space. For the next generation of tokamak and mirror fusion devices (i.e., TFTR, JET, JT-60, and MFTF-B) the optimum probe wavelength is \( \approx 100-200 \mu m. \)

Since the operating wavelength region is so clearly defined, it is of interest to survey the available source, detector, and component technology. The only available probe sources are lasers: electrical discharge, and optically pumped molecular lasers. Two electrical discharge lasers produce output power in the region of interest. The DCN laser produces \( \approx 250 \) mW at 195 \( \mu m \), while the DH2O laser produces \( \approx 50 \) mW at 118.6 \( \mu m \). As we shall see later, power output is important because multichannel interferometer systems are necessary which place constraints on the minimum acceptable laser power. A typical optically pumped laser system was shown in Fig. 6. Here the output of the grating tuned \( \text{CO}_2 \) laser is adjusted to coincide with a vibrational transition in the molecular gas lasing medium. Far-infrared lasing occurs via a rotational transition. Two particularly strong lasing transitions exist in the wavelength region of interest: \( \text{CH}_3\text{OH} \) at 119 \( \mu m \) and \( \text{CH}_2\text{F}_2 \) at 185 \( \mu m \). Power levels as high as 600 mW have been obtained with the former laser, while the latter yielded the highest known conversion efficiency.

The required laser power is obviously related to detector sensitivity and in fact helps to dictate the choice of detector. For example, pyroelectric detectors prove to be too insensitive given the number of interferometer channels and modulation frequency required in the JET tokamak. Although sensitive helium-cooled detectors such as InSb and Ga:Ge are available, one prefers room-temperature devices.
For this reason, a large number of laboratories employ GaAs Schottky diode detectors which exhibit low noise at room temperatures. At the short wavelengths of interest (≈ 100–200 μm) one cannot employ waveguide to channel radiation to the diode as is commonly done in the microwave and millimeter wave portion of the spectrum. In this region one makes use of quasi-optical antenna structures. As will be discussed in more detail in Sec. IV, a particularly useful configuration is the long wire antenna and corner cube reflector. In the FIR interferometer application, an NEP of ≈ 10^{-10} W/Hz^{1/2} is typical. This is actually much worse than is potentially obtainable and is associated with the low IF frequencies employed (≪ 1 MHz) which result in excessive 1/f noise. An alternative choice of room-temperature detector is the bismuth microbolometer shown schematically in Fig. 7. Each detector is formed by evaporating bismuth across the gap between the silver bow arms which serve as the bias leads and as an antenna structure to channel the incident radiation into the microbolometer. The resultant change in resistance manifests itself as a change in bias current. The small size of the bolometer (≪ 3 μm) results in adequate frequency response (≤ 3 MHz) for interferometry applications. Although the NEP of ≈ 10^{-9} W/Hz^{1/2} is inferior to that of the corner cube Schottky diode detector, the microbolometer has the virtue of low cost due to its ease of fabrication. However, as we shall see later, of more importance is the fact that, as shown in Fig. 7, integrated imaging microbolometer arrays have been fabricated.

The optical components employed in this spectral region deserve mention. The long propagation paths in fusion interferometers (up to 30–40 m) would result in unacceptable atmospheric absorption (≈ 0.1 m^{-1} absorption coefficient) if the beam path were not filled with dry nitrogen. This often leads to the use of a hollow dielectric waveguide for beam propagation since it easily allows for a controlled nitrogen atmosphere. Reported total averaged loss coefficients of ≈ 2.4 × 10^{-2} m^{-1} have been reported at 195 μm using such a system. In the case of focusing elements one normally employs reflective optics at these wavelengths in order to minimize material losses. A variety of beam splitter elements are available including meshes, mylar, and quartz.

It has been noted earlier that a multichannel interferometer system is required rather than a single beam interferometer. This can be seen by noting from Eq. (1) that the phase shift yields the chord-averaged line density and not the density directly. Therefore, one must unfold the interferometric information to determine \( n_e(r) \). The process is illustrated schematically in Fig. 8. Consider a beam passing through an inhomogeneous cylindrical plasma of radius \( R \). The phase shift is given by

\[
\phi(y) = \frac{2 \pi}{\lambda} n_e(r) \int_0^y n_e(r) dr/(r^2-y^2)^{1/2},
\]

where the prime denotes differentiation with respect to \( y \).

From Eq. (3), we see that one requires a large number of interferometer channels (or chordal paths) to accurately unfold the density profile even in the case of azimuthal symmetry. Using discrete detector and optical components, it is extremely difficult and expensive to employ more than 5–10 separate channels. This has motivated the recent developments of interferometer systems using integrated detector arrays. However, before examining this work, it is of interest to review the present state-of-the-art in discrete multichannel systems. Figure 9 displays a schematic of the 195 μm DCN FIR laser interferometer presently under construction for the JET tokamak. The ten-channel system

![Fig. 7. Bow-tie microbolometer 1.2-mm detector array. The bow arms are evaporated silver and the substrate is fused quartz (seen as black in the photograph). The narrow line that crosses the gap between the bow arms is evaporated bismuth which forms a microbolometer at each gap (after Ref. 70).](image)

![Fig. 8. Schematic of interferometer probe beam geometry (after Ref. 34).](image)
provides measurements along seven vertical chords and three "horizontal" chords, thereby providing a measure of plasma asymmetries. A dielectric waveguide is employed because of the long propagation distances. A schematic of the eight-channel CO₂ interferometer for the Los Alamos ZT-40 device is shown in Fig. 10. In this system the frequency shifted beam is provided by a Ge Bragg cell driven at 40 MHz. The high modulation frequency employed also permits this system to be utilized in the study of density fluctuations as will be discussed in Sec. III. Two techniques have been employed for phase detection. In the first, quadrature phase detection of the 40-MHz carrier yields the sine and cosine of the phase \( \phi(t) \) which is itself obtained by computing the arctangent. In the second, FM demodulation provides a direct measurement of the phase.

Let us now turn to the recent development of infrared and far-infrared multichannel systems which employ integrated detector arrays. The advantage of such systems is that the optical design is greatly simplified. An example of such a CO₂ laser interferometer system is shown in Fig. 11. The heart of the system is a 15-element PbSnTe detector array cooled to 77 °K. The other important feature to note from Fig. 11 is that the combination of the cylindrical ZnSe lens "A" and the parabolic (in one direction) mirror "B" produce a rectangular or sheet-like beam for efficient illumination of the plasma. The optical system is, therefore, seen to be much simplified over a corresponding 15-channel discrete detector system. The modulation and detection system is similar to those described earlier. An example of the time-resolved density profile in a resistive arc plasma is shown in Fig. 12.

The microbolometer arrays shown earlier in Fig. 7 are ideally suited for far-infrared and millimeter wave phase imaging measurements. The optical configuration for such a system is shown in Fig. 13. The need for the substrate lens is most easily seen if one makes use of reciprocity and considers an array of transmitting antennas on the substrate. The transmission efficiency into the substrate exceeds that into free space by the cube of the index of refraction. Thus most of the energy is radiated into the substrate. That fraction which exceeds the critical angle is trapped in the substrate and couples the individual detector elements degrading the imaging performance. A solution employed in the microwave region is to make the substrate extremely thin (typically \( \lambda_a / 10 \) in...
II. POLARIMETRY

Fusion plasma operate with high densities and temperatures (or drift velocities) and, therefore, the plasma diamagnetic field can become comparable to, or even larger than, the vacuum magnetic field. Although a determination of the internal magnetic field distribution is extremely difficult, it constitutes one of the most important measurements required on present devices. For example, in the case of a tokamak the radial current profile (or equivalently the poloidal field distribution) is required. In this section, we will discuss the use of Faraday rotation of the plane of polarization of an FIR beam to determine the local magnetic field.\textsuperscript{15,80,84,85-93}

Consider the propagation of a plane-polarized wave along a magnetic field. The plane of polarization rotates because the wave numbers of the \( R \) and \( L \) waves are not equal.\textsuperscript{15,18} The resulting rotation angle is given for \( \alpha_{pe} \), \( \omega_{ce} \ll \omega \) by

\[
\alpha \, [\text{degrees}] = 1.5 \times 10^{-12} \lambda_\nu^2 [\text{cm}] \int_0^1 n_e [\text{cm}^{-3}] B_\parallel [\text{kG}] dl, \quad (4)
\]

where \( L \) is the chord length in cm and \( B_\parallel \) is the component of the magnetic field along the propagation direction. The phase shift given by Eq. (1) is of course simultaneously present. One notes from Eq. (4) that the rotation angle is the chord average of the product of density and magnetic field.

The inversion process then requires an extremely accurate knowledge of \( n_e \) and thus the earlier statements concerning the need for multichannel measurements are even more true in this case.

The arrangement for tokamak poloidal field determination\textsuperscript{15} is shown in Fig. 14. There is also a competing polarization change due to the linear birefringence associated with the toroidal magnetic field. However, since this is \( \propto \lambda_\nu B_\parallel \) by operating at sufficiently short wavelength one can minimize this error. Several different techniques are currently being used at various laboratories to simultaneously measure phase shift and Faraday rotation. Figure 15 shows the arrangement successfully used on the TFR tokamak by modifying one of the HCN laser interferometer channels.\textsuperscript{90} More
FIG. 15. Experimental arrangement for Faraday rotation measurements in the FIR on the TFR 600 Tokamak (after Ref. 90). In the absence of the plasma all of the probe beam impinges on detector Dp. The presence of a poloidal field component causes the electric field vector to rotate giving rise to a partial reflection of the probe beam into the Faraday rotation detector Dp.

The complete details of the beam recombination in the two wire grid polarizing beam splitters are shown in Fig. 16.

An alternative technique, originally developed by Kunz and Dodel,87 is to apply a small high-frequency modulation \( \Delta \omega \) of the plane of polarization of the plasma probe beam in addition to the normal frequency modulation \( \omega_m \). The relevant output signal from the interferometer detector can then be shown to be given by

\[
u_\text{s}(t) \propto \cos(\alpha(t) + \Delta \omega \sin(\Delta \omega t)) \cos(\omega_m t + \phi(t)).
\]

With the proper signal processing, both \( \alpha(t) \) and \( \phi(t) \) can be obtained. Both Kunz and Dodel87 and the ISX group90 have employed a ferrite modulator to provide the required polarization rotation. Although successful, the disadvantage of this technique is that high modulation frequencies are not obtainable and that material losses become large at short wavelengths. An alternative technique has been described by Mansfield and Johnson92 in which the polarization of an optically pumped FIR laser is modulated by modulating the polarization of the CO2 pump laser. An alternative scheme72 now under test for use on the UCLA Microtor tokamak makes use of the polarization sensitivity of the bow-tie antennas shown earlier in Fig. 7. By staggering the antenna orientation as shown in Fig. 17, one can simultaneously measure the phase shift and polarization rotation with the separation performed easily in the low-frequency electronics portion of the system.
III. THOMSON SCATTERING

A. Review

Thomson scattering of electromagnetic radiation from plasma electrons represents one of the most powerful diagnostic tools utilized in controlled thermonuclear research. It permits spatially and temporally resolved measurements of electron temperature and density, ion temperature, magnetic field direction, and electron density fluctuation spectra.9,11,13,14,16,17,34,94-96

Electrons accelerated by the electric field of a laser radiate with an intensity proportional to the square of the density fluctuation amplitude \(S_n\). In thermal plasmas \(S_n \propto n_e^{3/2}\) resulting in the scattered power scaling linearly with electron density. If the fluctuations have frequency \(\omega\) and wave number \(k\), the scattered wave \((\omega_s, k_s)\) must satisfy energy and momentum conservation so that

\[
\omega_s = \omega_0 \pm \omega \quad \text{and} \quad k_s = k_0 + k.
\]

The wave number \(k\) of the fluctuations being studied is determined by the incident wavelength and choice of scattering angle, via the Bragg law

\[
k = 2k_0 \sin(\theta_s/2).
\]

In order to determine the scattering regime under investigation, it is useful to define a characteristic parameter

\[
\alpha = 1/\lambda k_D = \left[ \lambda_0 / 4\pi k_D \sin(\theta_s/2) \right].
\]

When \(\alpha < 1\), scattering is from individual, uncorrelated electrons and the scattered spectrum determines the electron temperature. In addition, the electron density may be determined from a measurement of the absolute scattered power. When \(\alpha > 1\), plasma waves and thermal ion fluctuations may be studied. The scattered spectrum in a thermal plasma results in a determination of ion temperature.

B. Electron temperature and density measurements

As discussed earlier, by performing small \(-\alpha\) Thomson scattering one can determine electron density. Because of the low level of scattered power at magnetic fusion plasma densities \((10^{13}-10^{15} \text{ cm}^{-3})\), one requires an extremely high power source. Therefore, this measurement depended critically on the development of high power, temporally, and spatially coherent low divergence lasers, as well as sensitive detectors. Specifically, the development of the ruby laser \((\lambda = 6943 \text{ Å})\) together with the availability of sensitive photomultiplier tubes in the visible portion of the spectrum led to the rapid adoption of this as a standard diagnostic tool.9,11,13,14,16,17,34,94-96

More recently, there has been considerable interest in the development of alternative sources in the 0.4-\(\mu\)m region. Basically, a large motivation for this work has been the need to obtain time-resolved electron temperature and density measurements. The ruby laser does not permit the high repetition rate pulses required for such a measurement. Therefore, experimenters are currently pursuing studies with Nd–glass \((1.06 \mu\text{m}),96\) frequency-doubled Nd–glass,100 and Raman-shifted XeCl excimer \((459 \text{ nm})\) and HgBr excimer \((502 \text{ nm})\)101 lasers.

The plan of this section is as follows: First, a discussion of ruby laser Thomson scattering systems is contained in Sec. III B 1. Since this is such a well-established technique, the discussion is primarily concerned with the various Thomson scattering (TVTS) systems which have provided spatially resolved measurements of electron temperature and density. The reader is referred to the many excellent references available for more general information on this topic.9,11,13,14,16

The remaining portions of this section are concerned with the development of the above-mentioned alternate lasers.

1. Ruby laser scattering

The first definitive measurement of tokamak electron temperature was made using a ruby laser Thomson scattering apparatus.97 The relative high output power (typically \(\gtrsim 4 \text{ J}, < 20\text{-ns pulse duration}\) together with the high efficiency of photomultiplier tubes \((\approx 15\% \text{ quantum efficiency})\) readily led to the adoption of this as a standard fusion plasma diagnostic. The earlier systems all provided only information concerning a single spatial region during each pulse. Typically, the scattered light was collected and then dispersed and analyzed by a diffraction grating spectrometer whose output was introduced into photomultiplier tubes for each wavelength channel.9,11,13,14,16

The absolute calibration of the laser scattering system is often performed using Rayleigh scattering from high neutral gas pressure (\(\approx 1-760 \text{ Torr}\) chamber filling. However, in some systems which employ a high finesse filter to reject the unshifted stray laser light this technique is not feasible. Fortunately, in this case an alternative calibration technique exists.102 One may instead utilize the anti-Stokes rotational Raman lines from molecular nitrogen, for example. This results in a series of discrete lines spread between 6 \(\text{Å} < \Delta \lambda < 100 \text{ Å}\) with a peak at \(\approx 25 \text{ Å}\). Howard et al.102 measured \(\approx 10^4 \text{ cathode photoelectrons using a 2-J laser}\). In addition, they suggest the use of \(H_2 (\Delta \lambda \approx 170 \text{ Å})\) in the calibration of systems to be used in the measurement of higher electron temperature.

In addition to photomultiplier tubes, systems have been constructed using avalanche photodiodes as channel detectors.103 These have several advantages. In addition to their low cost, they exhibit extremely high quantum efficiency \((\approx 65\% \text{ at 6943 Å})\). As we shall see later, even at 1.06 \(\mu\text{m}\) they still have relatively high efficiency \((\approx 30\%)\) which has permitted scattering studies at longer wavelengths.99,103 However, these detectors also have several problems which must be appreciated in order to obtain optimum performance. First, they are extremely temperature sensitive. Gain sensitivity \(\Delta G / \Delta T \approx 5\%-10\%/\text{°K}\) has been reported.103 A second point to understand is that for very low light levels, the diode noise dominates over photon statistics and the photomultiplier tube is superior. However, except for these exceedingly low levels, the higher quantum efficiency of the photodiode results in a higher SNR. The crossover point occurs at \(\approx 10^{-10} \text{ W} \approx (10 \text{ photons in a 30-ns pulse})\).

Thus far, the discussion has been concerned only with systems which provide a temperature measurement at a single spatial position and single time in a plasma discharge.
Obviously, one wishes to follow the temporal evolution of the spatial electron temperature profiles. A first important step in satisfying this need was made by Bretz et al. Specifically, they developed a multichannel system which now routinely provides electron temperature at 28 radial positions in the PLT plasma. Basically, the system consists of a microchannel-plate (MCP) image intensifier optically coupled to a cooled silicon-intensified-target (SIT) television tube. Twenty wavelength intervals are monitored and a dynamic range of $10^3$ is reported. A detailed description of this system follows.

A schematic of the entire PLT TVTS system is shown in Fig. 18. The laser beam is directed vertically upward into the plasma via a lens and prism beam splitter combination which may be scanned on a shot-to-shot basis to provide horizontal temperature profiles. Scattering from positions on a vertical line (see dashed line, Fig. 18) are collected and introduced into the spectrometer-camera system. A more detailed view of the optics is shown in Fig. 19. The scattered radiation enters a fiber-optic image dissector which provides 56 elements 1.6-mm wide by 6.3-mm long distributed evenly along the laser beam image. In addition, on a parallel arc there are 20 identical elements which sample the plasma background light. The fiber-optics portion consists of bundles of 75-μm-diam clad glass fibers with a transmission of only $\approx 35\%$.

The output from the spectrometer is focused onto the extended red S-20 photocathode of the MCP intensity via a camera lens. A more detailed schematic of the camera system is shown in Fig. 20. The output (P-20 phosphor) of the 40-mm-diam MCP image intensifier (Galileo Electro-optics Corp. Model 9040-1131) is fiber optically coupled to the SIT tube. The intensifier is gated ($\approx 100 \, \text{ns}$) and possesses a shutter ratio of $\approx 10^5$. The cooled SIT tube is also gated resulting in a combined shutter ratio of $\approx 10^8$. The MCP-SIT combination has a measured response which is linear over a range of about 500:1.
Variations of the system described by Bretz et al. are now in use on several tokamaks. Technological advances made since the original work have led to further refinements. In the PDX TVTS diagnostic system, Grek and Johnson have replaced the MCP-SIT configuration described above with a Texas Instruments CCD in a tube with a large photocathode. There are several reasons for the changed system. First, the dynamic range of the CCD exceeds that of the microchannel plate by over an order of magnitude. In addition, the read-out noise is reduced over that of the SIT tube. The disadvantage of using a CCD tube is that the chip must be baked at a lower temperature in the tube. This makes it difficult to obtain good quantum efficiency. However, quantum efficiencies of 4%–6% have been achieved which is comparable to the 5%–10% reported for the Princeton MCP-SIT system. Presently, the background light is measured ±1 ms before the laser pulse and transferred to part of the lower 320–256 pixel segment of the CCD array. This is an improvement over the PLT system where the channels were imaged at the top and bottom of the photocathode where the sensitivity is poor. In addition, the PLT system only had one-third as many background channels as scattered light channels. However, it has the advantage that the plasma background is measured during the laser pulse. Since the PDX TVTS system also has separate background channels similar to the PLT system, an arrangement is being implemented where they are used to appropriately scale the background data obtained 1 ms before the laser pulse. An interesting use of the PDX optical fiber was recently reported where the authors reported on its use to obtain temporally and spatially resolved Hβ measurements.

In addition to the improvements in the television camera system described above, additional advances have occurred since the original work reported by Bretz et al. As discussed earlier, the original PLT system used thin glass fibers. These packed badly and also had many of the fibers misaligned at the end of bundle resulting in an optical efficiency of only 35%. These were subsequently replaced on the PLT system by large plastic fibers yielding improved optical efficiency. However, these suffer from poor transmission over long path lengths and also have very poor radiation resistance making them unsuitable for the next generation of fusion devices. Therefore, the PDX system has adopted a system similar to that anticipated for TFTR which will use an 8–10 m length of fibers in a neutron and gamma-ray environment. To satisfy this need, a fused silica fiber has been developed which is thin clad in silicon resin to allow close packing while also being resistant to TFTR lifetime radiation doses. Using this new fiber an optical efficiency of 65% is obtained.

The above-mentioned improvements made in the PDX system have resulted in an increase in sensitivity over the original PLT system. In addition, the laser pulse energy has been doubled from the PLT laser (≈2–4 J, 0.25 mrad divergence). The resultant increase in system performance has permitted the utilization of the full 56-point spatial resolution. In the PLT TVTS system the 56 spatial positions were averaged to yield only 28 data points across the plasma. A planned further factor of 2 increase in the PDX laser energy will lead to even better data quality.

2. Nd and excimer lasers

Thus far, all of the discussion has been concerned with ruby laser scattering systems. However, in recent years there has been considerable interest in other lasers including Nd-glass (1.06 μm), frequency-doubled Nd–glass, and excimer lasers. The primary motivation is easy to understand. As mentioned above, one wishes to follow the temporal evolution of the electron temperature distribution during the entire plasma pulseduration (i.e., ≥1 s). However, the solid-state ruby laser medium is only suitable for single-point measurements due to the increased divergence, etc., which occurs in the high repetition rate mode. A secondary motivation is the limited output energy (≈10 J) available from the ruby laser. The Nd–glass laser, although also a solid-state system, can operate in the burst mode (≈20–50 Hz) for time scales of ≥1 s which is sufficient to provide temporal data. However, perhaps of more long range interest are the recent advances in the area of high brightness gas
exci mers lasers in the blue-green portion of the spectrum. The self-cooling nature of the gas laser permits operation at the required high repetition rates.\textsuperscript{101} Both of the above-mentioned systems are briefly discussed below.

As mentioned above, the attractive features of Nd-glass laser (1.06 \( \mu \)m) are very high output energy and high repetition rate. However, there are some complications which must be understood in order to optimize the performance of such a system. First, the linewidth of the laser tends to be relatively broad. For example, a 100-J system may have a linewidth of the 25- to 50-\( \AA \) range.\textsuperscript{100} This can lead to serious problems when scattering in low-temperature plasmas. In this case, the electron temperature is extracted from a Voigt profile rather than the usual Gaussian profile.\textsuperscript{101} Second, detectors other than photomultiplier tubes must be employed. As discussed earlier, silicon avalanche photodetectors exhibit high quantum efficiency in the vicinity of 1 \( \mu \)m.

Recently, the first work describing a high repetition rate tokamak scattering system was reported by Rohr \textit{et al.}\textsuperscript{99} The system is presently operating on the ASDEX tokamak at Garching. The Nd laser produces \( \approx 1 \) J in a pulse of 20 ns and is capable of operating at a repetition rate of \( \approx 50-100 \) Hz for burst durations up to \( \approx 8 \) s. The beam divergence is \( \approx 1.5 \) mrad. Their detectors consist of large-sensitive-area (1 \( \times \) 7 mm) Si avalanche photodiodes (modified RCA C30950) with integrated preamplifiers. The devices exhibit an NEP = 1.6 \( \times \) 10\textsuperscript{-13} W Hz\textsuperscript{-1/2} and a quantum efficiency of 45\% at 1.06 \( \mu \)m with a bandwidth of 15 MHz and a gain \( >20 \). This results in a minimum detectable number of photons \( \approx 10^{3} \) within the 50-ns gate duration. The observed low stray light level together with a sharp cutoff light filter (\( T_{1-0.6} \mu \text{m} \approx 10^{-4} \)) result in a stray light level corresponding to an electron density of less than 10\textsuperscript{11} cm\textsuperscript{-3}. Density calibration is performed using anti-Stokes rotational Raman scattering from \( \text{H}_2 \). Presently, three signals are digitally recorded: the laser scattering signal, a delayed monitor pulse and finally, several microseconds later, the plasma background. This sequence is repeated every 20 ms (up to 400 times) for each spectral channel.

The advances in frequency doubling technology have also made it practical to consider doubled Nd–glass systems.\textsuperscript{100} Peterson \textit{et al.}\textsuperscript{107} have described a system using a 20–40-J Nd–glass laser and a 25-mm KDP doubling cell. A net conversion efficiency of up to 14.5\% resulted in this first system being utilized in scattering from a low-temperature (\( \approx 10 \) eV) theta pinch plasma (\( n_e = 10^{16} \text{cm}^{-3} \)).

Recently, as discussed earlier, there has also been considerable interest in gas excimer lasers in the blue–green.\textsuperscript{101} Both HgBr excimer lasers (5020 \( \AA \)) and resonantly Raman-shifted XeCl excimer lasers (4590 \( \AA \)) have produced multi-joule outputs with low divergence. Boody and McNeill\textsuperscript{101} have carefully investigated the design of a XeCl oscillator–amplifier system which employs a Pb-vapor Raman shifter. They anticipate pulse energies of \( \approx 3 \) J with 0.25-mrad divergence for a 1-s long burst (10–20 Hz pulses).

### C. Thermal ion temperature measurement

\( \text{CO}_2 \) or optically pumped FIR lasers are optimum for large-\( \alpha \) scattering measurements in magnetic fusion plasmas. If we further assume that we are scattering from thermal fluctuations, additional constraints are placed on the required laser power. For example, a laser pulse duration of \( \approx 1 \mu \text{s} \) results in a required power of the megawatt level (dependent on background, optical losses, and detector sensitivity) even when heterodyne detection techniques are employed.\textsuperscript{104–117} In addition, the sources must possess good spectral qualities (i.e., single mode operation) in order to accurately resolve the complicated scattered spectrum. Typical requirements are 1–10 MW power levels, \( \approx 1 \mu \text{s} \) pulse durations, and spectra widths (at the 10-dB points) \( \Delta f \leq 100 \) MHz. Calculations indicate that longer pulse durations significantly improve the final signal-to-noise levels\textsuperscript{108–117} by increasing the available integration time. Laser systems marginally approaching the above specifications at both CO\textsubscript{2} and FIR laser wavelengths have been developed.

Before proceeding to a discussion of the technology, it is appropriate to consider the limitations of the thermal scattering technique. First, a tacit assumption is made that in the frequency range of interest the plasma is thermal in nature.\textsuperscript{34} The presence of nonthermal fluctuations can lead to large scattering which will mask the thermal feature. The presence of impurities can also significantly distort the scattered spectrum. In addition, the shape of \( S(k) \) depends sensitively on \( T_e/T_i \). One must, therefore, accurately know \( T_e \) in order to unfold the scattered spectrum. The effect of impurities can be significant either when they are highly ionized (\( Z \) large) or when they constitute a significant fraction of the plasma.\textsuperscript{115} Here we should note that one can alternatively employ large \( \alpha \) scattering to determine the effective charge

\[
\bar{Z} = \frac{\sum \bar{Z}_j N_j}{\sum \bar{Z}_j N_j} \tag{7}
\]

in a plasma.\textsuperscript{115}

It should be stressed that in the case of thermal scattering, the postdetection of the signal is extremely important. Specifically, in order to resolve the ion feature, the scattered and down-converted signal must be fed into a multichannel spectral analyzer such as a filter bank and detector system. As is well known, the signal-to-noise ratio at the detector output is given by

\[
S = \frac{s}{1 + s\sqrt{1 + Br}}, \tag{8}
\]

where \( s \) is the signal-to-noise ratio at the mixer output, \( B \) is the bandwidth of each filter, and \( \tau \) is the detector integration time. One can, of course, express \( s = P_s/P_n \), where \( P_s \) is the collected scattered power and \( P_n \) is the total noise referred to the receiver input. One sees that for \( \tau \ll 1 \) the total signal-to-noise ratio is independent of scattered power and only depends on \( Br \). Therefore, to obtain large \( S \) one requires \( Br \gg 1 \). However, to obtain good spectral resolution one desires \( B \) as small as possible. Furthermore, \( \tau \) is limited to the duration of the laser pulse. A reasonable compromise is to take \( B = 50 \) MHz and \( \tau \approx 1 \mu \text{s} \) which results in \( S = 7 \).

A question arises as to the expected accuracy of the ion temperature measurement via Thomson scattering. Detailed numerical simulations were performed by Watterson \textit{et al.}\textsuperscript{114} which included the use of Monte-Carlo techniques in order to model the statistical properties of the power spec-

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trum. In addition, the authors treated the case of impurities. In this latter case, they found that by cutting off the low-frequency portion of the spectrum they were able to eliminate the systematic deviation of the mean \( \Delta T_i \). They predict that one can obtain ion temperatures with accuracies of \( \approx 10\%-15\% \).

In the following we will briefly survey the various CO\(_2\) and FIR thermal scattering systems.

1. CO\(_2\) systems

There have, to date, been several successful demonstration experiments employing CO\(_2\) lasers in the measurement of ion temperature \( (\approx 1-5\ eV) \) in high-density \( (n_e \approx 10^{15}\ cm^{-3}) \) arc plasmas.\(^{116-120}\) The system employed by Holzhauer\(^{116}\) consisted of a single-mode (longitudinal and transverse) hybrid CO\(_2\)-TEA laser\(^{121}\) which delivered 0.25 J in a 1-\(\mu\)s duration pulse. The homodyne detection system \( (\text{NEP} = 5 \times 10^{-19}\ W\ Hz^{-1}) \) consisted of a Ge-Hg photococonductor \( (\eta = 0.11) \) cooled to 20 K. A portion of the unscattered laser radiation \( (\approx 0.3\ W) \) serves as a local oscillator. Alignment and calibration is facilitated by the use of Brillouin scattering at the Bragg angle from acoustic waves launched in neutral gas (typically air at STP) via a piezoelectric transducer. Using matched Gaussian probe and LO beams, a mixing efficiency of \( >80\% \) is obtained. Using the measured \( T_e \approx 1.3\ eV \) and assuming \( T_i = T_e \) (calculated difference \(<5\%) \) yielded excellent agreement between the calculated and measured spectra. This experiment definitely demonstrates the viability of the diagnostic technique. A slightly different approach was taken by both Pasternak and Offenberger\(^{118}\) and Peebles and Herbst.\(^{117}\) In these studies, high-power \( (10-70\ MW) \) pulsed TEA lasers were employed. Spectral data were obtained via high-finesse Fabry-Perot interferometers which were scanned on a shot-to-shot basis. The interferometer output was detected using cooled photoconductors. Due to the low ion temperature, deconvolution was required to separate the instrument function from the scattered spectrum. Such a technique would be expected to work best in very-high-density devices \( (n_e > 10^{14}\ cm^{-3}) \), where the scattering angle is large enough that the stray light falling on the detector can be reduced to a sufficiently low level.

Recently, Kasperek and Holzhauer\(^{120}\) have extended the earlier studies of Holzhauer.\(^{116}\) As shown in Fig. 21, a hybrid CO\(_2\)-TEA laser is again employed. In these studies the technique was shown to be capable of resolving the change in the spectrum when other ion species were added to the hydrogen plasma (see Figs. 22 and 23).

![Fig. 21. Experimental arrangement for ion temperature determination in a thermal arc plasma via collective CO\(_2\) laser scattering (after Ref. 120).](image1)

![Fig. 22. The measured spectrum obtained in CO\(_2\) scattering from a pure hydrogen arc plasma yielding \( T_i = 1.6\ eV \) (after Ref. 120).](image2)

More recently, Taylor and Bretz have developed and installed a prototype infrared Thomson scattering system on the Princeton PDX tokamak which is to serve as a proof-of-principle experiment for a proposed TFTR ion temperature diagnostic.\(^{113}\) The probe laser consists of a 2.5-m-long hybrid CO\(_2\)-TEA laser which produces \( \approx 0.2\ J \) in a 2-3-\(\mu\)s duration pulse. The output is passed through a two-pass amplifier system resulting in \( \approx 2\ J \) of probe beam. The entire laser is feedback stabilized by diverting a small amount of the oscillator output onto a pyroelectric detector via a grating tuned to the 10P (20) line. An important feature to note is that contained within the hybrid oscillator structure is a separate cw CO\(_2\) laser local oscillator. This is attenuated to \( \approx 100\ mW \).

The scattered and local oscillator beams are combined in an Sb compensated 0.2\times0.2-mm Ge:Cu detector (supplied by Airborne Instruments Laboratories) cooled to 4 K. The measured NEP in the laboratory was \( <2 \times 10^{-19}\ W/\text{Hz} \) (100- to 1200-MHz IF bandwidth). The system as installed on PDX exhibits somewhat higher noise temperature \( (\approx 5 \times 10^{-19}\ W/\text{Hz}) \). After preamplification, the IF output signal from the detector is introduced into an eight-channel filter bank covering 100-1200 MHz. Each channel contains an appropriate bandpass filter, detectors, video amplifiers,
and gated integrators. The integrated signals are then multiplexed and may be displayed directly on an oscilloscope.

2. Far-infrared systems

During the past few years, there has been considerable interest in the development of far-infrared laser scattering systems for the thermal scattering measurement. Before proceeding to a survey of the far-infrared technology, it will be helpful to briefly discuss the problems peculiar to the far infrared. First, to obtain \( \alpha > 1 \) results in large-angle (\( \theta > 10^\circ \)) scattering in the far infrared. While this provides improved spatial resolution (compared to CO\(_2\)), it also increases the demands on machine access. Second, both refraction and diffraction problems are more serious than for CO\(_2\) scattering. In addition, the synchrotron radiation background is significantly higher in the far-infrared portion of the spectrum which increases the required laser power. The above constraints, together with the realities of present day source and detector technology, have led to the majority of the present efforts being directed toward the development of a 385-\(\mu m\) wavelength system.

The most intensive efforts on the development of a pulsed FIR scattering system have taken place at the MIT National Magnet Laboratory by Woskoboinikow et al. Therefore, we will concentrate primarily on a description of their apparatus which is presently installed on the Alcator C tokamak for a proof-of-principle test experiment.

Figure 24 shows a schematic of the entire scattering system. The probe laser is a \( \text{D}_2\text{O} \) laser operating in the vicinity of 385 \( \mu m \). This is a convenient wavelength as discussed earlier both because it satisfies the various scattering constraints and because the \( \text{D}_2\text{O} \) laser has yielded the highest output in the FIR. The \( \text{D}_2\text{O} \) laser is described in detail in Ref. 128. Basically, it consists of a high-power \( \text{D}_2\text{O} \) oscillator cavity pumped by a high-power, single-mode tunable (\( \approx 1 \) GHz) CO\(_2\) laser chain operating on the 9R (22) line. The CO\(_2\) system consists of a low-power grating tuned CO\(_2\) oscillator (\( \approx 100\text{-MJ} \) output) with a ZnSe étalon for longitudinal mode selection. The output is amplified to \( \approx 50 \) J (200–500 MW) in a multiple amplifier system containing an SF\(_6\) absorption cell to prevent oscillations on the 10-\(\mu m\) branch.

There are several reasons for employing a tunable CO\(_2\) laser pump. First, the 385-\(\mu m\) transition of \( \text{D}_2\text{O} \) is a stimulated Raman transition. Therefore, its frequency is tuned by the pump frequency. As we shall soon see, this is of great importance when one attempts to reduce the stray light background in the actual scattering experiment. Second, the \( \text{D}_2\text{O} \) linewidth is considerably reduced when pumped with a single-mode CO\(_2\) laser due to the Raman nature of the transition. Finally, the maximum output of the \( \text{D}_2\text{O} \) laser is fre-
frequency offset from the low-power absorption line center. The CO$_2$ output pumped a 4-m-long D$_2$O oscillator containing a grating (to eliminate the 359-$\mu$m output) with a Fox-Smith longitudinal mode selector. The system has produced $\approx 0.5$--0.9 MW with linewidths less than 100 MHz, pulse durations of $\approx 100$--200 ns and polarization of 90%.

The quasi-optical heterodyne receiver system (see Fig. 24) was developed by H. R. Fetterman and his colleagues at the MIT Lincoln Laboratory. A detailed description of the various components may be found in Ref. 67. Briefly, the scattered radiation is combined with the local oscillator beam via a frequency selective diplexer (see Ref. 133) and introduced into a quasi-optical Schottky diode corner cube mixer. The 380.6-$\mu$m local oscillator is a 25-mW deuterated formic acid (DCOOD) optically pumped far-infrared laser whose frequency was actively stabilized to the LO video maximum. With an 8.1--10.6 GHz, IF amplifier, the measured noise temperature of the system is 20 000 K DSB (NEP $= 2.7 \times 10^{-19}$ W Hz$^{-1}$). The amplified IF output is introduced into a 32-channel filter bank. The detected video output from each IF channel is then passed to a gated current integrator. Using the receiver portion of the apparatus, the synchrotron radiation background in the vicinity of 385 $\mu$m was measured in both Alcator A and C$^{135,136}$ and for $B_T < 90$ kG was found not to be a limiting factor.

Great care was exercised in the design of the external optics for the Alcator-C proof-of-principle experiment$^{123,135,136}$ together with the beam and viewing dumps. Nevertheless, it was found that the poor tokamak access resulted in unacceptable levels of stray light which reduced the dynamic range of their system. Their final solution to this problem was to employ a narrow band gas absorption cell as a filter at the front end of the receiver. Using a 6-m-long N$_2$O gas cell they measured 60-dB rejection. Here it should be noted that this approach absolutely required a tunable FIR laser as the N$_2$O absorption line is $\approx 360$ MHz away from line center CO$_2$ pumping of D$_2$O. This technique coupled with the addition of C$_2$H$_4$CN to the D$_2$O laser (to reduce the low level -60-dB D$_2$O laser linewidth) has resulted in a measured upper limit for nonthermal wide bandwidth scattering on Alcator C to $\approx 13$ dB above thermal. With a few minor improvements, it appears that this system will conclusively demonstrate the feasibility of this technique.

The MIT FIR scattering experiment described above was designed as a proof-of-principle experiment. If this proves successful, an actual system for a plasma device such as TFTR will require considerably longer pulse duration to accurately resolve the ion feature. Such a system is under investigation at the Princeton Plasma Physics Laboratory$^{24,131}$.

As a final note, recent work on a far-infrared ring resonator D$_2$O laser configuration has yielded encouraging results.$^{137}$ Specifically, essentially transform limited laser linewidths were achieved in the case of the ring resonator. The reason for the improvement over the standard linear resonator is easy to understand. The standing-wave electric field patterns existing in linear laser cavities lead to spatial hole burning and encourage the simultaneous oscillation of several longitudinal cavity modes even when the lasing transition is homogeneously broadened. Unidirectional, traveling-wave ring lasers eliminate these spatial inhomogeneity effects and thus offer improvements in mode quality and efficiency.

### D. Collective scattering from waves and fluctuations

As discussed earlier, when $\omega \gg 1$ one can probe waves and fluctuations in plasmas. The study of the power spectra of naturally occurring and induced plasma waves are important for a variety of reasons. First, low-frequency microinstabilities, such as drift waves and convective cells, are thought to be the main cause of the anomalous transport existing in present day tokamak plasmas. Second, rf heating at the lower hybrid and ion cyclotron frequencies are likely to develop into cost-effective alternatives to neutral beam heating. A study of the launching, propagation, mode conversion, and damping of such waves is of importance in understanding the heating mechanisms and improving efficiencies. Third, the external launching of waves in fusion plasmas and a measurement of their dispersion relations can lead to a spatially and temporally resolved determination of plasma parameters such as ion temperature.

Scattering from plasma waves has the advantage that large enhancements in the scattered power above thermal can be obtained, thus requiring lower power laser sources. The above is easily seen for the case of a coherent density fluctuation [i.e., $\delta n = \delta n \cos(kz - \omega t)$]. Under wave number matching conditions, the scattered power is given by

$$P_s = \frac{1}{4} P_o \omega^2 \lambda_0^2 (\delta n)^2 L_v,$$

where $r_e$ is the classical electron radius and $L_v$ is the length of the scattering volume. The lower power requirement has meant that cw CO$_2$ lasers of $\approx 10$--100 W can be utilized. The main requirement is good spectral quality and amplitude stability to enable low-frequency ($\approx 50$ kHz) waves to be studied. Since the scattering angles are typically $\ll 1^\circ$, good beam quality is also necessary especially in a homodyne system—where stray light can be troublesome. In the far infrared (FIR), cw power levels as low as 2 $\mu$W have proved sufficient to study many phenomena. This is due both to the fact that the scattered power from a coherent wave scales as $\lambda_0^2$ and to the recent advances in the technology of low noise FIR radiometers.$^{62-67,134}$ As will be discussed below, the FIR offers a wide range of wavelengths which are easily obtainable via optical pumping of molecular lasers. In addition, the larger scattering angles available in the FIR also permit multichannel scattering systems to be constructed.$^{138}$ This permits single-shot dispersion measurements to be obtained routinely. Finally, microwave (millimeter wave) scattering using extended interaction oscillators (EIO's) or carciitrons (O-type backward wave oscillators) continues to be an invaluable technique.$^{139-140}$ In the following: a brief survey of the present technology and recent results will be given for each of these techniques.

#### 1. Microwave scattering

Microwave scattering from laboratory plasmas developed into a useful diagnostic during the latter half of the 1960's.$^{150,151}$ However, it was not until the middle of the 1970's that this technique was successfully applied to toka-
An excellent review of some of this work is contained in Ref. 149. Before proceeding to a comprehensive brief survey of the results, it will prove helpful to review the available technology.

A scattering system requires one or two high-frequency sources depending upon whether it is configured in the homodyne or heterodyne mode. It is, therefore, appropriate to consider the various source options. The limit in reflex klystron technology is \( \approx 220 \text{ GHz} \). For example, the Varian VRY 213A provides \( \approx 10 \text{ mW} \) output in the 170–220-GHz region. The advantage with klystrons is the simplicity and low cost of the associated power supplies. On the negative side, however, are the long delivery time and low output power of these high-frequency tubes. In addition, further improvements appear rather doubtful. A related linear beam velocity modulated tube, called the extended infra red oscillator (EIO), offers watt level performance up to \( \approx 230–300 \text{ GHz} \) (Varian # VKY-2432E,F). Mechanical tolerances in the resonant cavity and segmented drift tube would appear to preclude significant increases in maximum operating frequency. The carcinotron or O-type backward wave oscillator\(^{155,156}\) appears to be the most viable candidate for use in millimeter wave scattering systems on the next generation of magnetic fusion devices. A variety of tubes are available on a commercial basis from Thomson, CSF. For example, the CO.10.1 device produces nearly a watt at \( \approx 300 \text{ GHz} \).

Recently, several TH4218C tubes have been supplied to UCLA for use in scattering systems on the UCLA Microtor and University of Texas TEXT tokamak devices. These tubes cover the 360-406-GHz region (\( \lambda_0 \approx 800 \mu m \)) and produce up to 100 mW. Under a development contract with the European Space Agency Laboratory (ESTEC), Thomson has developed the TH4211 tube which has provided power levels in excess of 50 mW at frequencies above 500 GHz.\(^{154,155}\) A problem with the carcinotron is that it is a voltage tunable device (\( \approx 10–20 \text{ MHz/V} \)). To provide a narrow line source for studies such as drift wave microturbulence one requires an extremely well regulated low-ripple high-voltage (\( \approx \text{mV} \) at 10–12 kV) power supply. This requirement coupled with the need to install sophisticated protective circuitry has rather limited the power supply market. Other tube-type sources such as gyrotrons\(^{156,157}\) produce considerably more power and have been suggested as scattering sources.\(^{158,159}\) However, thus far no experiments have been performed, therefore, leaving open such questions as linewidth, etc.

Solid-state technology will eventually provide suitable sources for millimeter wave scattering systems (local oscillator and probe beams). IMPATT diodes have provided outputs of \( \approx 25–50 \text{ mW} \) at frequencies as high as 240 GHz.\(^ {160}\) Furthermore, 2-mW output has been reported at 400 GHz using LN\(_2\) cooling.\(^ {161}\) Although IMPATT diodes are rather noisy compared to klystrons, the output can be narrowed using either injection locking (fundamental or subharmonic) or spectral filtering. A related device called the TUNNETT offers the promise of higher-frequency operation (up to \( \approx 800 \text{ GHz} \)) with lower noise albeit with lower output power (\( \approx 1 \text{ mW} \)).\(^ {162}\)

Let us now turn to a brief discussion of mixers. In the millimeter wave region, the majority of low-noise mixers employ a Schottky diode chip mounted in a fundamental mode rectangular waveguide. A Schottky barrier diode chip (typically GaAs) is mounted on the broad face of the waveguide and contacted with a whisker antenna. Tuning is accomplished with a movable backshort. Since the waveguide is designed to operate in the fundamental propagation mode, the dimensions become quite small at the desired operating frequencies. Nevertheless, during the past few years there has been dramatic progress in the development of low-noise waveguide mixers.\(^ {163–168}\)

With the above technology level in mind, it is appropriate to perform a simple calculation to provide a feeling for the minimum detectable density fluctuation level. For this purpose, we assume \( P_0 \approx 5 \text{ mW} \), \( \lambda_0 \approx 1 \text{ mm} \), and \( T_{\text{sys}} = 5000-\text{K} \text{ SSB} \). Then using Eq. 9 we find that

\[
P_t \approx \frac{1}{4} P_0 \tilde{n}^2 \pi L^2 \lambda_0 \approx 10^{-28} (\tilde{n}_{\text{em}})^2 \text{ W}.
\]

The major noise sources are the synchrotron radiation background and the receiver noise level. To calculate their magnitude, we require the receiver bandwidth. Here let us assume we have launched a coherent wave and assume a bandwidth of \( \approx 50 \text{ kHz} \). Thus, assuming \( T_e = 5 \text{ keV} \) and \( B_0 = 40 \text{ kG} \) and some chamber wall reflectivity (i.e., an effectively optically thick plasma), it is appropriate to assume a 5-keV blackbody in calculating the synchrotron radiation background. The resultant plasma and receiver noise background powers are \( 3.8 \times 10^{-11} \text{ W} \) and \( 3.4 \times 10^{-15} \text{ W} \), respectively. Without recourse to sophisticated signal processing techniques, we still calculate a minimum detectable \( \tilde{n}_{\text{min}} = 6 \times 10^6 \text{ cm}^{-3} \) or with \( n_0 = 10^{14} \text{ cm}^{-3} \) we have \( \tilde{n}/n_0 \approx 10^{-5} \). In the case of a broad spectrum such as appropriate for drift wave microturbulence (i.e., \( B \approx 10^6 \text{ Hz} \)), we still obtain \( \tilde{n}_{\text{min}} = 3 \times 10^9 \text{ cm}^{-3} \).

It is of interest to briefly examine a microwave scattering system. Figure 25 shows the 140-GHz transmitting and receiving antenna configuration employed by Mazzacato for the PLT microwave scattering studies.\(^ {169}\) An array of reflector antennas are installed within the vacuum vessel permitting the study of the wave number spectrum of low-frequency microturbulence at various spatial locations. The electromagnetic waves are launched with the electric field parallel to the toroidal magnetic field. The scattering angles range from \( \approx 2–60^\circ \) permitting the study of fluctuations with \( k_r \), ranging from 2-30 cm\(^{-1}\) (\( \lambda_t \approx 0.2–3 \text{ cm} \)) with a resolution (\( \Delta k_t \approx \pm 2 \text{ cm}^{-1} \)). Recently, the resolution has been improved to \( \pm 1 \text{ cm}^{-1} \). The scattering is primarily along the poloidal magnetic field and thus \( k_i \) is not measured. Since the publication of Ref. 169, the system has been configured in the heterodyne mode, thereby permitting the study of wave propagation directions.

2. Far-infrared scattering

Recent advances in millimeter wave source and detector technology have resulted in a spate of papers describing FIR scattering measurements.\(^ {138,170–179}\) In the following we will briefly describe some of the technological advances.

The heart of most of the above-mentioned systems is a low-noise mixer. In the preceding section, the conventional microwave fundamental mode mixer was described. Unfor-
fortunately, at the highest frequencies of interest, mechanical tolerances impose a severe impediment to fundamental waveguide mixer fabrication. Therefore, there has been considerable interest in so-called quasi-optical structures which serve to eliminate some of the fabrication problems as well as possessing low loss in the submillimeter region.

The first quasi-optical mixer configuration to demonstrate low-noise performance in the far-infrared region was the biconical mount. This design was originally conceived by Gustincic for a NASA-JPL atmospheric radiometry program and later extended to the FIR for radio astronomy and plasma diagnostic applications. As shown in Fig. 26, the diode chip and whisker are mounted at the apex of a biconical antenna.

An alternative approach is based on the familiar traveling long wire antenna. To produce a focused beam from this configuration Krautle et al. added a corner reflector. Figure 27 shows a realization of such a mixer by Fetterman et al. The structure is easily analyzed utilizing standard image techniques and can be thought of as a simple four-element phased array. This design has proven to work well and has demonstrated low-noise performance at frequencies in excess of 1 THz. Table I shows representative corner cube mixer performance obtained by Fetterman and his colleagues at the MIT Lincoln Laboratory. As can be seen, even at frequencies as high as 2.7 THz, adequate mixers exist for scattering measurements.

The other major constituents of an FIR scattering sys-

![Fig. 25. Experimental arrangement for millimeter wave collective scattering in the Princeton PLT tokamak.](image)

![Fig. 26. Schematic of the quasioptical biconical Schottky diode mixer (after Ref. 64).](image)

**TABLE I. Noise temperature measurements.**

<table>
<thead>
<tr>
<th>$\lambda$ (µm)</th>
<th>$T_s$ (Room temperature) (K)</th>
<th>$T_s$ (cooled to 41 K)</th>
</tr>
</thead>
<tbody>
<tr>
<td>432.6</td>
<td>5700</td>
<td>3800 K</td>
</tr>
<tr>
<td>432.6</td>
<td>4200</td>
<td>3800 K</td>
</tr>
<tr>
<td>184</td>
<td>19000</td>
<td>10.5 (dB)</td>
</tr>
<tr>
<td>119</td>
<td>32000</td>
<td>10.5 (dB)</td>
</tr>
</tbody>
</table>

Fusion plasma
tem are the local oscillator and probe sources. Thus far, most of the experiments have been performed in the homodyne mode so that only a single source is required. In addition, the systems have primarily employed far-infrared lasers. As discussed earlier in Sec. I, both optically pumped molecular lasers (C\textsuperscript{13}H\textsubscript{3}F, CH\textsubscript{3}I) and electrical discharge lasers (HCN) have been utilized. These typically produce \(\approx 10\)- to 100-mW output power levels which have been sufficient to perform scattering measurements with good signal-to-noise ratio. However, at wavelengths as short as 500 \(\mu\)m carcinotrons are also available.

Let us now briefly discuss some representative far-infrared scattering systems. The first far-infrared tokamak scattering experiment was performed on the UCLA Microtor device and is described in detail in Refs. 171 and 172. A schematic of the apparatus is shown in Fig. 28. The radiation source is an optically pumped FIR laser. In the experiments, two FIR laser lines were typically employed. The C\textsuperscript{13}H\textsubscript{3}F laser generated \(\approx 5\) mW of power at 1222 \(\mu\)m (245 GHz), while the CH\textsubscript{3}I laser produced \(\approx 30\) mW at 447 \(\mu\)m (671 GHz).

The output from the laser is reflected from an adjustable mirror onto a mesh beam splitter which directs part of the beam towards a quasi-optical biconical Schottky barrier diode mixer. This serves as the local oscillator for the homodyne detection system adopted in these measurements. The remainder of the beam is weakly focused \(\left(d_0 \approx 1.5\right)\) cm to a region located on a horizontal plasma diameter. The scattered radiation is collected by a movable mirror which permits a variation of scattering angle from 0°–20°. The entire laser and optical assembly can be translated along the major radius of the tokamak allowing spatial regions to be probed from the edge to the center of the plasma. The frequency shifted scattered beam is recombined with the local oscillator via the beam splitter producing an output signal from the mixer at the wave frequency. The receiver system is absolutely calibrated using known hot and cold blackbody loads together with a \(Y\)-factor analysis.

The Microtor scattering system was first used to study low-frequency microturbulence. The general features have been similar to those found in microwave\textsuperscript{159–160,164} and CO\textsubscript{2}\textsuperscript{16} scattering studies. For example, most of the fluctuation energy lay below 150 kHz with \(S(k,\omega) \propto \omega^{-n}\) above 100 kHz, where \(n\) varied between 1.5 and 3.

In addition to the low-frequency microturbulence studies, there has been considerable effort at UCLA devoted to wave measurements (via FIR scattering) in the ion cyclotron range of frequencies (ICRF). These experiments have been quite successful and represent the first direct observation in a tokamak of externally excited electrostatic modes in the ICRF. In this instance two FIR systems were in use at symmetric toroidal locations for the study of the toroidal plasma propagation.

In the Microtor tokamak, the toroidal field varies both spatially and temporally. Figure 29 illustrates this temporal variation together with typical FIR scattering signals. These data were obtained in a deuterium plasma possessing a small hydrogen concentration (\(\leq 10\%\)). The applied rf was \(\approx 20\) MHz which is approximately the second-harmonic frequency for deuterium. The observed signal changes are primarily caused by the time varying toroidal field. This sweeps \(\omega/\omega\text{\textsubscript{T}},\) and, therefore, modifies \(k_\omega\) in a manner determined by the wave dispersion. From the time history of the spatial dependence of the scattering signals the wave dispersion relations

---

**Fig. 28.** Experimental arrangement for cw FIR laser scattering on the UCLA Microtor tokamak (after Ref. 173). A portion of the output of the FIR laser is split off by the beam splitter and serves as a local oscillator for the quasioptical mixer in the homodyne scattering system. The remainder is focused onto the major horizontal plane in the tokamak and serves as the probe beam. The scattered and frequency shifted radiation is then combined with the LO beam in the mixer.

**Fig. 29.** Typical cw FIR scattering data from ICRF waves in the UCLA Microtor tokamak (after Ref. 173). As shown in the figure, the Microtor toroidal magnetic field varies temporally as well as spatially. At a fixed spatial position this sweeps \(\omega/\omega\text{\textsubscript{T}}\) in time and, therefore, modifies \(k_\omega\) in a manner determined by the wave dispersion. This gives rise to time-varying phase and electric field amplitude signals determined by wave number matching and the wave properties.
can be obtained. Typical experimental dispersion data are shown in Fig. 30 and compared with theoretical calculations.

More recently, there have been several reported scattering measurements at 337 μm employing HCN lasers. At Kyoto University, studies of low-frequency microturbulence in the WT-2 tokamak have been reported both in pure ohmically heated plasmas and during rf current drive. The system arrangement is somewhat similar to the Microtor system described earlier. The output of an 80-mW HCN laser passes through a Mylar beam splitter providing a local oscillator beam and probe beam. At these frequencies (890 GHz), material losses become significant so that a concave mirror is used for focusing rather than a lens as in the Microtor system. A quasi-optical corner cube mixer is used for signal downconversion. An interesting feature of the Kyoto work was the utilization of a plastic Bragg cell for calibration of the scattering apparatus. Acoustic waves are launched via a piezoelectric transducer (PZT) in a cell made of TPX plastic (4-methylpentene-1) which is transparent in both the visible and FIR regions. By performing Bragg scattering from the acoustic waves, the system optical collection efficiency and wave number resolution can be determined.

Unfortunately, because only a single detector channel was available in all of the FIR systems described above, dispersion relations were obtained by a shot-to-shot variation of the scattering collection angles. Therefore, the reliability of the frequency and wave number spectra was totally dependent upon the reproducibility of the plasma and the phenomena under investigation. In the study of random processes such as low-frequency microturbulence, the fact that there is only a single channel is an obvious restriction. In addition, diagnostic applications such as ion temperature determination via cw scattering from externally launched waves are dependent upon single-shot time-resolved wave dispersion data.

To satisfy the above-mentioned needs, a multimixer far-infrared scattering apparatus has been developed at UCLA. In this system, shown schematically in Fig. 31, scattered radiation from plasma density fluctuations is collected by six mirrors which are designed to accept signals corresponding to six discrete wave numbers. In order to accurately obtain dispersion relations, as well as the absolute level of the density fluctuations, considerable effort has been devoted to the system calibration. To achieve this goal, system parameters such as collection efficiency, wave number resolution, scattering volume, and detector sensitivities have been carefully measured. The calibration procedures have included acoustic cells, laboratory test wave scattering, as well as standard hot–cold load techniques.

An acoustic cell similar to that utilized by the Kyoto University group is employed for optical system calibration. Unfortunately, even if the PZT driving voltage is kept constant, the amplitude of the acoustic waves is a sensitive function of frequency. This variation arises since PZT transducers exhibit a series of resonances related to their thickness. Because of this strong frequency dependence, the amplitude of launched waves must be measured within the acoustic cell at each frequency. This is accomplished by means of small angle He–Ne laser scattering orthogonal to the FIR scattering plane.

The above-mentioned system is used on the Microtor tokamak to study ICRF heating, as well as low-frequency microturbulence. In addition, a larger six-channel system is also installed and operating on the TEXT National User's Facility tokamak administered by the University of Texas. Although the primary purpose of this system is to serve as one of a set of diagnostics in a coordinated program of study of low-frequency microturbulence, it has also been provided with the capability for ICRF studies. The optical system is motorized permitting both a vertical and horizontal scan of the scattering volume over ~80% of the entire plasma cross section. The initial scattering studies are being performed at 447 and 1222 μm using a homodyne laser scattering configuration. This will later be supplemented by an 800-μm carcinotron heterodyne arrangement so that wave propagation direction can be determined. Dielectric waveguide is employed for beam propagation because of the long path lengths involved.

E. CO₂ laser scattering

The viability of the cw CO₂ laser scattering technique was first demonstrated by Surko et al. in a laboratory test plasma. In this initial pilot experiment, scattering was successfully performed from driven electron Bernstein waves of known amplitude, frequency, phase, and wavelength as independently determined by Langmuir probes. Studies of ion acoustic turbulence in laboratory plasmas have also been reported. This was followed by the successful implementation of this technique in low-frequency microturbulence studies on the Princeton ATC tokamak. In addition, microw turbulence measurements and lower-hybrid wave propagation studies were performed on the MIT Alcator A and C.
Tokamaks.\textsuperscript{187-191} More recently, microturbulence studies have been performed on the Princeton PDX tokamak,\textsuperscript{192} as well as scattering from Alfvén waves in the Pretex tokamak at Texas.\textsuperscript{193} Excellent reviews of much of this work is contained in Refs. 96 and 194.

The short wavelength (10 µm) of the CO\textsubscript{2} laser leads to slightly different techniques and operating conditions than those discussed for the microwave and FIR scattering cases. First, for typical magnetic fusion plasma parameters, the collective regime (α > 1) is reached only for small scattering angles (θ < 1°). This leads both to advantages and disadvantages. The small scattering angle is quite attractive in large fusion plasmas where the access is severely limited. However, the small scattering angle results in a chord-averaged scattering with an attendant lack of spatial resolution. As we shall shortly discuss, this latter limitation can be eliminated with a crossed-beam correlation technique.\textsuperscript{187} In addition, it should be noted that the small scattering angles preclude the construction of multichannel systems.

Tokamak cw CO\textsubscript{2} scattering systems have employed lasers in the 35-300-W power level region.\textsuperscript{184-190} Here it is particularly important to employ a single-mode laser with good beam quality (TEM\textsubscript{00}). This restriction arises since it is quite easy for the unscattered light to saturate or even destroy the detector because of the small scattering angles involved. Figure 32 shows a schematic diagram of a typical small angle homodyne CO\textsubscript{2} scattering experiment. The detector is a cooled photoconductor. A frequently used detec-
tor is the liquid-helium-cooled Ge:Cu photoconductor counter doped with Sb.\(^{197}\) This detector possesses both good sensitivity over a broad IF bandwidth and a high saturation level. For example, Airborne Instruments Laboratories provided a 0.2×0.2-mm element together with a 4-K impedance matching network for the PDX CO\(_2\) thermal scattering measurement described in Sec. III C 2. The measured NEP was \(<2 \times 10^{-19}\) W Hz\(^{-1}\) over an IF bandwidth of 0.1–1.2 GHz. In addition, the same company provided the detector employed by Slusher et al.\(^{199-201}\) in their 2.45-GHz lower-hybrid wave studies in the Alcator A tokamak. This device also possessed a measured NEP of \(2 \times 10^{-19}\) W Hz at 2.45 GHz at an LO power level of 180 mW. An alternative detector is the HgCdTe photovoltaic device which has been utilized by Meyer and Mahn in the Wendelstein VIIA stellerator scattering system\(^{196}\) and studied by Bretz and Taylor\(^{113}\) for a possible TFTR CO\(_2\) scattering system. Here it should be noted that both of the above-mentioned detectors do suffer from degraded performance for IF frequencies in excess of \(\approx 4\) GHz. Therefore, although they are perfectly suitable for scattering from waves up to the lower hybrid range of frequencies, there may be difficulty in scattering from waves near the upper hybrid frequency.

Pulsed CO\(_2\) lasers have also been employed to study collective fluctuations.\(^{196-201}\) The first CO\(_2\) scattering experiment was performed by Kornherr et al. using 1-J TEA CO\(_2\) laser. The turbulence in a theta pinch plasma was studied using a cooled (4 K) Ge:Hg photoconductor as a video detector. Similar studies were made by Breiz and DeSilva\(^{199}\) on a fast theta pinch. Again, video detection of the scattering signal was performed with a Fabry–Perot interferometer used for frequency selection. Hybrid lasers consisting of a high-pressure TEA discharge and a low-pressure cw discharge have also been employed.\(^{202}\) Typical narrow line outputs of \(\approx 3–200\) kW (\(\approx 2-10\)-µs pulse duration) have permitted hybrid scattering measurements. Finally, pulsed low-pressure laser (\(\approx 50\) W, \(100\) µs) scattering measurements have also been performed. Recently, for example, Fahrbach et al.\(^{201}\) used such a laser together with homodyne detection to measure the lower-hybrid drift fluctuations in the boundary layer of a high-beta theta pinch plasma.

**F. Far-forward scattering**

In the preceding discussion, it was tacitly assumed that the scattering angle was larger than the divergence angle of the input probe beam. The opposite limit is also of interest for plasma diagnostics and has been referred to as far-forward scattering, ultraforward scattering or as homodyne scattering.\(^{96,203-207}\) A related technique is the use of “phase scintillations” of a probing electromagnetic wave as in radio astronomy.\(^{208}\)

Far-forward scattering depends upon the fact that an electromagnetic wave traversing a refractive medium such as a plasma undergoes phase and amplitude modifications which are related to the wavelength, position, amplitude, and frequency of phase fluctuations in the medium. The simplest case to treat is that of an incident coherent Gaussian beam.\(^{203}\) The geometry is illustrated in Fig. 33. The laser beam propagating in the \(x\) direction is diffracted by the plasma wave (propagating transverse to the beam in the \(x\) direction) which acts as a moving sinusoidal phase grating \(\phi (x) = \Delta \phi \sin (k x - \omega t)\). As a result of passage through the wave, the beam acquires an intensity component at the wave frequency and successively smaller components at harmonics of the wave frequency. The component at the wave frequency, i.e., the first-order diffracted beam, corresponds to the IF signal from the heterodyne detector. It can be shown\(^{207}\) that the amplitude \(I_1\) and phase \(\phi_1\) of this component are given by

\[
I_1 = \frac{P_0 \Delta \phi}{2 \pi w_0^2} e^{-(w_i^2 + u_i^2) - (w_i u_i \mp v_i^2)} \left[ e^{-2u_i v_i - 2 \cos \omega v_i t} \right]^{1/2},
\]

(10)

and

\[
I_1 \sin \phi_1 = \frac{P_0 \Delta \phi}{2 \pi w_0^2} e^{-(w_i^2 + u_i^2)(e^2 - e^{-2u_i v_i - 2 \cos \omega v_i t})},
\]

(11)

where \(P_0\) is the power of the incident beam,

\[
u = \frac{w_i}{1 + \rho^2},
\]

and

\[
\rho = \frac{w_{01}}{w_i} + \frac{k_0 w_{01}^2}{w_i} \frac{L_i - L}{L^2},
\]

where \(z_{i1} = k_0 w_{01}\) is the Rayleigh length and

\[
L = \left( \frac{L_1}{L_2} + \frac{1}{f} \right)^{-1}.
\]

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as shown by James et al.206,207 in a far-infrared laboratory experiment, even minute departures from the Gaussian beam assumption can lead to problems in interpreting the scattering data.

As a final note, one can also study density fluctuations by measuring the correlations between parallel beams. This technique has been employed by Jacobson209 who employed the seven parallel beams of his multichannel interferometer system to study turbulent fluctuations in the ZT-40 device.

IV. SYNCHROTRON RADIATION DIAGNOSTICS

A. Principles

It is well known that electrons orbiting about magnetic field lines radiate energy at the cyclotron frequency and harmonics. The measurement of the resultant electron cyclotron emission or synchrotron radiation has become an invaluable diagnostic especially in the determination of tokamak electron temperature. References 210–214 provide particularly excellent reviews of the technique. Therefore, in the following we will only briefly treat the physical principles and will primarily concentrate on the diagnostics instrumentation.

The determination of electron temperature profiles depends first on the fact that in nonrelativistic plasmas the electron cyclotron frequency is purely a function of the local magnetic field strength. Therefore, if the magnetic field varies monotonically (i.e., the $1/R$ variation of the tokamak toroidal field) then the cyclotron frequency is a unique function of distance along the line of sight. In addition, if we provide viewing optics which restrict the field of view perpendicular to the line of sight, a frequency selective receiver will detect emission from a small localized volume as indicated in Fig. 35. The second constraint is that the viewing system must not detect radiation which has been reflected from the vessel walls. This leads naturally to the requirement that the plasma be optically thick at the frequency of interest. In addition, this implies that the plasma radiates as a blackbody in this frequency range, and, therefore, that the...
plasma temperature is directly obtained using absolutely calibrated receivers.

The emission intensity from a thermal plasma can be written as

$$I(\omega) = I_{BB}(1 - e^{-\tau})/(1 - \xi e^{-\tau}),$$

(12)

where $I_{BB}$ is the blackbody emission [ = $\omega^2 K_{BB}(R)/8\pi^2 c^2$ for $\omega_0 < K_{BB}$], $\tau$ is the optical depth, and $\xi$ is the effective chamber reflectivity which is dependent upon the optical system. When $\tau \rightarrow \infty$ we see that the emission at $\omega$ approaches blackbody. For typical fusion tokamak parameters with $\omega_{pe} < \omega_{ce}$, the fundamental emission in the ordinary mode is optically thick as is the emission in the second-harmonic extraordinary mode. In addition, reported measurements on the PLT tokamak indicate that the fundamental extraordinary mode radiation is also blackbody.215 Since this is not possible for electromagnetic excitation, the authors attribute it to thermal excitation of electron Bernstein waves. Here it should be noted that these measurements were performed with the detection system on the inside of the torus since the $R$-wave cutoff is thereby avoided.

Even in the case of optically thin lines, the temperature profile can still be determined. For example, at higher harmonics the ratio of the intensities of the ordinary and extraordinary modes yields the temperature profile though $I_{m}^{\omega}(\omega)/I_{m}^{\omega}(\omega) \approx K_{T_e}(R)/m_0 c^2$. An additional method, described by Celata and Boyd,210 permits the determination of the temperature profile through the measurement of the ratio of the emission intensities at successive harmonics in the optically thin region. The temperature is then given by

$$K_{T_e}(R) \approx \frac{M_m I_{m+1}[\omega_{m+1}(R)]}{[I_{m}[\omega_{m}(R)] m_0 c^2]},$$

(13)

where

$$M_m = \frac{2m+1}{(m+1)^{2m+1}} \left(2m + \frac{1}{2m} + 3 \right)^{m+1}.$$  

More recently, Blau216 has performed detailed calculations and experimental comparisons related to the use of the optically thin third-harmonic emission for temperature measurements in the Doublet III tokamak. This technique permits an increase in the permissible maximum cut-off density for a given toroidal field strength. In the limit that $\tau < 1$ we have, after expanding the exponentials in Eq. 13 and using $\tau \propto nT_e^2$, that

$$K_{T_e}(R) \propto \left(\frac{1 - \xi}{n_e} I_{BB}\right)^{1/3}.$$  

(14)

One, therefore, requires a knowledge of $\xi$ together with the density profile. Nevertheless, preliminary measurements are encouraging.

It should be noted that more work must be done in system calibration. To date, many of the cyclotron emission systems simply provide relative temperature profiles and must be normalized to Thomson scattering data. However, since these systems can be absolutely calibrated, they in principle provide a “stand alone” diagnostic. Unfortunately, systematic deviations from Thomson scattering temperatures have been observed in PLT studies.217 Efforts are currently being made to explain this discrepancy.

Thus far, the discussion has been restricted to thermal plasmas in which the runaway electron population is negligible. In the case where this is not true, the emission spectrum can be dominated by the runaway component. Fortunately, however, the desired operating regime for tokamaks minimizes the runaway component.

In addition to electron temperature profiles, one may also in principle obtain the density profile from cyclotron emission data.218 This depends upon the fact that for emission at an optically thin harmonic

$$I(\omega) \approx \tau I_{BB}/(1 - \xi).$$  

(15)

For low density with $m < 5$, $\tau \approx n_e T_e^{-1}/(m\omega_e)^2$ which implies that $I(\omega) \propto n_e T_e^2/(1 - \xi)$. To illustrate the technique assume that the second harmonic is optically thick while the third harmonic is optically thin. Therefore,

$$I(2\omega_e) \propto T_e(\tau)$$

(16)

and the ratio of the third-harmonic specific intensity at time $t$ with respect to its value at some reference time $t_R$ is given by

$$I_3(t)/I_3(t_R) = \left[ n_e(t)/n_e(t_R) \right] \left[ T_e(t)/T_e(t_R) \right]^3.$$  

(17)

This yields

$$n_e(t) = n_e(t_R) I_3(t_R)/I_3(t).$$  

(18)

Although the validity of this technique has been demonstrated experimentally,218 to date cyclotron emission measurements have only been utilized for temperature profile determination.

Potentially, one can also utilize the polarization properties of the cyclotron emission to determine the poloidal magnetic field distribution in a tokamak. However, wall and window depolarizing reflections can seriously complicate this measurement. For example, the ordinary mode emission (relative to the extraordinary) has been found to be anomalously high. Polarization measurements have yielded values ranging from $\approx 0\%$ to $\approx 30\%$.219-224

B. Techniques

Currently, tokamak synchrotron radiation measurements are performed either with InSb video detectors or with conventional millimeter wave heterodyne receivers. Since the InSb detector is a broadband device, a frequency analyzing element must be inserted between the plasma and the detector. Operating systems include grating polychromators,225,226 fixed or scanning Fabry-Perot226,227 and scanning Michelson interferometers.221,226,228,229 The latter two systems are limited to observation of fluctuations with frequency $\approx 100$ kHz while the grating polychromator system permits the study of fluctuations up to $\approx 1$ MHz. In addition to the InSb systems, there has been considerable interest in conventional millimeter wave Schottky diode heterodyne receiver systems. These possess significantly wider bandwidth than the InSb systems. However, lack of suitable local oscillators and mixers has limited them to operation at frequencies $\approx 140$ GHz, whereas the InSb systems easily operate from $\approx 100-1000$ GHz. However, recent advances in mixer and source technology offer the promise of extension of low-noise heterodyne receiver systems to beyond 1 THz. In the following we will attempt to survey the current level of technology as well as future directions.
1. Heterodyne receivers

A typical millimeter wave fundamental mode \( f_{10} = f_{\text{signal}} + f_{10} \) heterodyne receiver system employs a Schottky barrier diode as the mixing element. The key requirements are a broadband mixer and a sweepable local oscillator.

Let us begin with the local oscillator. At frequencies above \( \approx 100 \) GHz, the only suitable broadband, sweepable local oscillator is the carcinotron described earlier. For the purposes of the present discussion we note that practical tube specifications are: center frequency range: 100–600 GHz; bandwidth: 20%; output: \( > 20 \) mW. One problem arises immediately. To cover the necessary frequency range for cyclotron emission measurements, one requires several tubes. However, the cost per tube is \( \approx 50,000 \) dollars. In addition, since this is a voltage tunable tube \((\approx 8–40 \text{ MHz/V})\) which operates at voltages up to \( \approx 10 \) kV, the power supply requires considerable regulation \((\leq 50 \text{ mV ripple})\) and is, therefore, relatively expensive \((\approx 50,000–40,000)\). An additional problem associated with carcinotrons is that they possess relatively large noise sidebands. The reported noise temperatures\(^{230-232}\) have varied significantly from \(1000–3000\) up to \(50,000\) K at IF frequencies in the 1.4–1.7-GHz region. These variations appear to be associated with the age and design of the tubes. A recent FIR receiver design study concludes that a noise temperature of \( \leq 5000\) K at an IF frequency of 1.5 GHz is realizable.\(^{233}\) In fixed tuned receiver applications diplexers can be easily constructed with noise rejection of \( \approx 20\) dB so that this is a negligible contribution. Unfortunately, this is not practical for swept cyclotron emission receivers. Thus, the local oscillator noise appears to be the limiting factor in determining the minimum detectable signal level (or equivalent temperature). A final problem with carcinotrons concerns the significant variation of the output power level with frequency. This necessitates the use of leveling systems to ensure relatively uniform mixer conversion loss across the desired operating band.

Let us now turn to a discussion of mixers. For the moment we will restrict our attention to fundamental mixers \((f_{\text{signal}} \approx f_{10})\). It should be noted here that the mixer must not only possess a sufficiently low-noise temperature, but it must also exhibit a sufficiently broad instantaneous bandwidth of at least 20% so that it will be compatible with the carcinotron local oscillator. As mentioned in Sec. III D 2, at these frequencies quasi-optical mixers such as the biconical mount\(^{64-66}\) and the corner cube mixer\(^{62,63,66}\) are employed as well as standard waveguide mixers.\(^{234}\) They should all perform equally satisfactorily below 600 GHz yielding system temperatures of \( \approx 4000\) K. However, it should be pointed out that this figure may rise in an actual measurement system since the diplexer must be broadband, thereby eliminating the Fabry–Perot and Michelson couplers commonly employed.

In any of the above-mentioned cases, one requires a high-frequency local oscillator of \( \approx 10 \) mW output power. However, as discussed earlier these tend to be quite expensive. It is, therefore, of interest to examine systems in which the local oscillator frequency is at a subharmonic of the signal frequency. One common approach is the harmonic mixer in which the diode generates harmonics of the local oscillator frequency which in turn mix with the signal. A simple rectangular waveguide mixer system can exhibit noise temperatures as low as \( 10,000 \) K in the 200–300-GHz region. Significantly better performance has been achieved by Zimmerman (SBS conversion loss of \( \approx 6\) dB at \( \approx 200\) GHz). More recently, Erickson and Fetterman have obtained a noise temperature at 671 GHz of 23,000 K DSB with a waveguide harmonic mixer.\(^{234}\) An alternative is the so-called subharmonic mixer in which the two diodes are mounted in reverse polarity spaced a small fraction of a waveguide apart. They are then switched at twice the local oscillator frequency and mix the signal at \( \approx 2f_{10}\). Carlson and Schneider\(^{235}\) have demonstrated performance at 230 GHz, equal to that of fundamental mixers. In combination with a fundamental mixer, the subharmonic mixer, therefore, doubles the operating range of each carcinotron.

It is of interest to explore other choices of mixer design which require much lower local oscillator power levels \((\leq 10\) \( \mu\)W). Such mixers are all cryogenic in nature. A review of the present status can be found in Refs. 236–240.

Assuming the ultimate success of cryogenic mixers, one can envision the use of frequency multipliers to better utilize each carcinotron. For the above mixers one can contemplate the use of frequency doublers and triplers so that each carcinotron serves as local oscillator to mixers at two or more frequencies. A description of frequency multipliers can be found in Ref. 236.

It is now of interest to examine actual systems. Figure 36 shows a schematic representation of the fast-scanning heterodyne receiver system planned for TFTR.\(^{24,241}\) This will provide temperature profiles every 5 ms by sweeping the frequency range of 75–210 GHz. The periscope chamber protects the silica windows from sputtering as well as allowing the insertion of a blackbody load into the field of view for absolute system calibration. The radiation collected by the conical horn lens-corrected antenna is divided in the three-port splitters. The 75- to 110-GHz and 110- to 170-GHz bands are covered by fundamental mixers while the 170- to

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**Fig. 36. Schematic of a scanning millimeter wave electron cyclotron emission detection system for the Princeton TFTR tokamak (after Ref. 241).**
210-GHz range employs a crossguide second-harmonic mixer. Figure 37 shows a schematic of the basic radiometer design, which is similar to the PLT system. The leveling system is required since the power output of the BWO typically varies by a 10:1 ratio across the band. The choice of IF bandwidth is a compromise between the requirement of high IF frequency (to avoid the BWO noise sidebands) and low frequency (to reduce the spatial separation between the two cyclotron layers observed in a double sideband measurement). The indicated IF frequency results in a spatial resolution of \( \approx 1 \text{ cm} \).

Thus far the discussion has centered around heterodyne measurements in thermal plasmas. However, the same techniques can be employed to study nonthermal plasmas (i.e., runaway distribution). In this case one often simply requires measurements at a few discrete frequencies rather than continuous coverage as discussed above. This allows for the use of other local oscillators such as optically pumped molecular lasers. Using a combination of carcinotron and laser local oscillators, measurements were made in the 7th and 30th harmonic range in the UCLA Microtor tokamak. More recently, measurements have been reported up to the 7th harmonic in the Alcator device. Nevertheless, it appears that the InSb bolometers to be described in the next section are more suitable for such measurements.

2. Fourier transform spectroscopy

Liquid-helium-cooled InSb crystals and germanium bolometers are excellent broadband detectors throughout the millimeter and submillimeter wave region of interest. However, cyclotron emission measurements require that these detectors be employed together with some form of frequency selective element. A commonly employed configuration utilizes a Michelson interferometer of the Martin-Puplett polarizing type. Such Fourier transform spectrometers are utilized in systems throughout the world.

Figure 38 shows a schematic of a typical Michelson interferometer arrangement for cyclotron emission studies.
The system envisioned for use on TFTR will produce 70 scans per second with an 11-ms scan duration and will cover the range from 90–1800 GHz (Δf = 3 GHz). Therefore, the entire temperature profile is obtained with a spatial resolution of 1–10 cm dependent upon central magnetic field and location in the plasma. Blackbody loads of known temperature are employed to provide absolute system calibration.

One limitation of the Fourier transform spectrometer arises due to the mechanical scanning of the mirror. This results in typical scan times of ≈ 10 ms. However, one sometimes wishes to improve the time resolution. One possible technique has been demonstrated by Bartlett et al.\textsuperscript{254} In their configuration, shown in Fig. 39, a transient pulsed plasma is placed in one path of the interferometer providing the scanning. A microwave interferometer provides a plasma density monitor. Comparisons with a standard Michelson established the validity of this technique. However, the drawback is that the system described in Ref. (250) was essentially single shot. Repetitive pulsing of the plasma is required in order to provide the necessary time history of the temperature profile.

A different approach is to employ a rotating mirror interferometer. Using such a system Campbell et al.\textsuperscript{251,252} obtained repetitive (100 Hz) scan times of 0.5 to 2 ms duration with corresponding spectral resolution of 6–3 GHz.

3. Grating polychromator

As discussed above, the Fourier transform spectrometer is a relatively slow instrument normally requiring ≈ 10 ms per scan. However, one often requires information concerning the temporal evolution of the temperature profile on a much faster time scale (i.e., ≈ 1 μs). In this case one wishes to employ multiple detectors with dispersive filters to select the proper wavelength region for each detector. For example, Boyd et al.\textsuperscript{219} employed a five-channel InSb detector (time response < 10 μs) system with fixed bandwidth (f / Δf ≈ 2) filters in each channel for measurements on the ATC tokamak. However, the grating polychromator\textsuperscript{225,226} appears to offer improved performance.

Figure 40 shows a schematic of the grating polychromator used on PLT.\textsuperscript{225,226} The system consists of a 0.5-m, f / 3 Czerny–Turner instrument with five output channels, each employing an InSb detector. For typical PLT parameters, the 1.625-mm pitch grating results in a channel spacing corresponding to 5 cm in major radius for a total coverage of 20 cm. As mentioned above, this system permits much improved time response compared to the Fourier transform spectrometer. However, there are several limitations which should be understood. First, the use of multiple detectors complicates the instrument calibration. Second, the limited frequency span (and, hence, spatial coverage) means that only a portion (≈ 25% on PLT) of the temperature profile is obtained.

More recently, Fisher, Boyd, and Cavallo have constructed a ten-channel system for second-harmonic measurements shown schematically in Fig. 41. The incident synchrotron radiation is collimated by a cylindrical TPX plastic lens and then dispersed by an aluminum reflection grating. During a single tokamak discharge a temperature profile covering a 26%–32% range in major radius is obtained. The detector elements are indium antimonide hot electron bolometers cooled to 4.2 K.

4. Scanning Fabry–Perot interferometer

An alternative approach is to employ a Fabry–Perot interferometer as the dispersive element.\textsuperscript{220,227} Further, one has the choice of operating in the fixed or the scanning mode. To eliminate ambiguities due to transmission in higher orders one should of course employ a low-pass filter in tandem with this instrument.
V. ION BEAM PROBE DIAGNOSTICS

During recent years, increasing use has been made of heavy-ion beam probes for the time- and space-resolved measurement of plasma potential in magnetic fusion devices. In addition, this technique also offers the possibility of providing similar information on plasma density, fluctuations, and electron temperature. This is extremely important as the knowledge of the amplitude and phase of such quantities will permit the detailed study of transport processes. For example, the radial particle loss flux \( j_r = \langle \frac{\partial n}{\partial t} \rangle \) is given by \( j_r = (mc/2\pi\beta_0) |\hat{n}| |\hat{\phi}_p| \sin \delta \) in the case of simple sinusoidal fluctuations where \( \delta \) is the \( \hat{n} - \hat{\phi}_p \) phase difference. Before discussing the diagnostic technology, it will prove helpful to briefly review the basic physical principles of this measurement.

Figure 43 shows a simple schematic of the measurement. A primary beam of singly ionized energetic heavy ions (typically Ti\(^+\), Li\(^+\), Na\(^+\), K\(^+\), Rb\(^+\), or Cs\(^+\)) is injected across the magnetic field into the plasma. If the Larmor orbit is large enough (\( v_i \) large enough), the primary beam will completely traverse the plasma. However, a small fraction of the primary ions will undergo ionizing collisions primarily producing doubly ionized ions. These ions carry away information about the plasma potential at their point of ionization, as well as possessing significantly different trajectories than the singly ionized primary ions. By measuring the energy difference between the secondaries entering the analyzer and the primaries leaving the ion gun, the space potential is determined to within a few volts. The spatial localization is obtained from the particle trajectories. The need for an energetic heavy ion beam is easily demonstrated. Relating the beam potential \( V_B \) to the perpendicular kinetic energy \( (mv^2) = eV_B \) and noting that the Larmor radius \( r_{LO} = v_i / \omega_o \) allows us to write \( V_B = e(r_{LO}/\beta_0)^2/m_e^2 \). Then requiring \( r_{LO} > a \) (plasma radius) results in the inequality

\[
V_B \geq e(a\beta_0)^2/m_e^2
\]

which is satisfied for heavy ions and/or very high ion beam energy.

The local plasma density is easily calculated from the ratio of the secondary to primary current. These currents are related through the effective ionization cross section by

\[
I_s/I_p = 2n_s/\sigma_{ab} = 2n_s \langle \sigma v \rangle / (2eV_B/m)^{1/2};
\]

where \( I_s \) is the length of the primary beam sampled by the detector. Therefore, if the temperature distribution is known, one can infer the local density.

By probing the same plasma volume with beams of different ion species, the electron temperature can be determined from the secondary ion currents if the ionization cross sections are known. Alternatively, one can employ the ratio of triply ionized to doubly ionized ions to infer the electron temperature using single ion species injection.

Figure 44 contains a simplified schematic of an ion beam probe system. One begins with an ion source. This is basically fabricated from zeolites (ion exchange resins)

![Diagram](image-url)

Fig. 43. Basic concept of heavy ion beam probing (after Ref. 253).

![Diagram](image-url)

Fig. 42. Schematic of a rapid scanning FIR Fabry–Perot interferometer for electron cyclotron emission measurements (after Ref. 227).

![Diagram](image-url)

Fig. 41. Schematic of 10-channel FIR grating polychromator for electron cyclotron emission measurements on TFTR (courtesy of D. Boyd).
which are heated to thermionic emission temperature. This is mounted in a Pierce-type gun. The extracted ions are accelerated and focused into a beam via a cylindrical lens. Systems with ion energies of up to 60 keV and beam currents of \( \approx 10^{-6} \) to \( \approx 10^{-4} \) A have been produced with focal diameters of \( \approx 0.5 \) cm.\(^{13}\) The electrostatic deflection plates (sweep plates) permit one to sweep the primary beam across the plasma. Therefore, using a single detector it is possible to map the entire plasma profile by sweeping the beam injection angle and energy.

The secondary ions are analyzed and detected in an electrostatic energy analyzer. Great care is taken in the design to ensure high resolution as the plasma potential is typically two to three orders of magnitude lower than the acceleration voltage. Therefore, one requires the measurement of energy differences to one part in \( 10^3 \) or \( 10^4 \). This is accomplished by using a modified parallel plate analyzer with a wide aperture and split plate detectors. The difference signal caused by the beam striking the two plates is amplified and fed back to the top analyzer plate which forces the beam to provide equal currents to the two plates. The feedback voltage is then directly proportional to the space potential. In operation, this feedback controlled analyzer has provided a sensitivity of \( 10^{-4} \) V.

Thus far, there has been no discussion of the time response of the diagnostic system. This is presently limited primarily by the high gain amplifier employed in the feedback loop. The 10- to 100-nA total secondary ion current is to be compared with the \( \approx 1 \)-nA sensitivity of the best available 1-MHz bandwidth amplifier. Therefore, the present systems are suitable for examining fluctuations in the MHD and drift wave frequency region, but not the ion cyclotron range of frequencies. A more basic limitation is imposed by the finite transit time of the ions through the analyzer (\( \approx 0.1-1 \) ns).

Let us now briefly examine some actual systems. Figure 45 shows a schematic of the system installed on the TMX Mirror device at the Lawrence Livermore Laboratory.\(^{257}\) The system was installed on the center cell of the tandem mirror device. The ion beam provides a spatial resolution of \( \approx 1 \) cm\(^{-3}\) and a sensitivity to potential of \( \approx 2 \) V with an absolute accuracy of \( \approx 25 \) V. The analog bandwidth of the

![Beam Probing System Diagram](image-url)
system was \( \approx 1 \) MHz. Using digital recorders, fluctuations to \( \approx 40 \) kHz (100 kHz sampling) could be followed for the entire discharge. Higher frequency fluctuations of up to \( \approx 1 \) MHz (2-MHz sampling) could be studied for periods of \( \approx 0.5 \) ms. The beam probe can be operated in several different manners. First, the observation point can be held constant yielding the time history of the potential at a fixed position. Alternatively, the observation point can be scanned (via the sweep plates) so that multiple radial profiles can be obtained during a single plasma discharge. Typical scan times of 100 \( \mu \)s to 10 ms yielded dozens to several diameter scans per shot.

As described earlier, one can measure ion temperature by using two different ion species probe beams assuming the effective ionization cross sections are known. The ratio of the secondary current signals is simply the ratio of the effective ionization cross sections at the appropriate electron temperature. Typical data obtained in a laboratory arc plasma are shown in Fig. 46 together with the inferred density and temperature profiles. As can be seen, the technique works well. The difficulty with this measurement is that at higher electron temperatures the ratio becomes relatively insensitive to \( T_e \).

It is also of interest to consider the scaling of future systems. We recall that the required beam energy scales as \( a^2 B^2 / m_i \). The Livermore tandem mirror device is being upgraded (TMX-Upgrade). The vacuum tank has been increased in size by a factor of 2 while the magnetic field has been increased from 2 to 3 kG. This necessitates a new ion source capable of injecting a Cs or Ti ion beam at energies up to 100 keV. \[260\] As part of the anomalous transport study program on the TEXT tokamak, a heavy ion beam probe system is under construction since it offers the possibility of directly measuring both \( \bar{n} \) and \( \phi \) as discussed earlier. The large size (limiter radius 28 cm) and magnetic field ( \( \approx 30 \) kG) results in a required beam energy of \( \approx 500 \) keV. Finally, a system with a beam energy of 6 MeV is required for the MFTF-B mirror device presently under construction.

In addition to a positive ion beam, other probes are possible including neutral beams (see Sec. VII) and negative ion beams. For example, Ishii et al. \[261\] suggest the use of a gold negative ion beam together with laser fluorescence detection of the resultant neutral beam. An additional possibility is the detection of the positive ion beam which is produced by double detachment.

VI. X-RAY DIAGNOSTICS

The high temperatures of present day magnetic fusion devices result in a considerable portion of the energy radiated in the x-ray region. This coupled with the short wavelength of the emission together with the transparency of the plasma has led to a wide variety of x-ray diagnostic techniques. In particular, x-ray measurements provide spatially resolved information on electron and ion temperatures, impurity content and transport, as well as MHD fluctuations and disruptions. The basic principles of x-ray diagnostics together with reviews of more recent measurement techniques are contained in Refs. 8, 9, 13, 16, 17 and 262–265. In the following we will only briefly describe the basic principles of the measurements and will primarily concentrate on the more recent developments.

The x-ray spectra in current fusion devices are comprised of the x-ray continuum together with line radiation from high-Z impurities. The former consists of free–free transitions (bremsstrahlung) and recombination radiation (free-bound transition). The continuum radiation energy

![Fig. 46. (a) Typical \( I_\phi / I_p \) data for Na\(^+\) and K\(^+\) beams. Note the different vertical scales for the two curves. (b) Electron temperature and density profiles deduced from the beam probe measurements (after Ref. 259).](image-url)
density is approximately given by
\[
\frac{dE_{\text{rad}}}{d(h\nu)} \propto Z_{\text{eff}} T_e^{-1/2} \exp\left(-\frac{h\nu}{kT_e}\right)
\]
\[
\times \left[g_{\text{ff}} + \sum_{\chi_n^m} \frac{2\chi_n^m \bar{g}_{\text{ff}}}{kT_e} \exp\left(\frac{\chi_n^m}{kT_e}\right)\right]
\]
where \(g_{\text{ff}}\) is the free-free Gaunt factor, \(\bar{g}_{\text{ff}}\) is the Gaunt factor for free-bound transitions into shell \(n\), and \(\chi_n^m\) is the ionization energy of the ion from the quantum level \(n\). As can be seen from the above, the electron temperature is obtained from the slope of the continuum. At the temperature of current fusion devices, light ions such as hydrogen and helium are fully stripped. Therefore, the line radiation typically consists of \(K-, L-,\) and \(M-\) shell radiation from high-\(Z\) impurities. Quantitative measurements of the impurity line emission yield important information concerning fusion plasma parameters. For example, the fractional abundance of high-\(Z\) components (i.e., Fe, Cr, Ni, etc.) can be deduced from the \(K\) lines. Then, by comparing the enhanced spectrum to a calculated hydrogen bremsstrahlung spectrum, the low-\(Z\) constituency can be obtained finally yielding \(Z_{\text{eff}}\).

High-\(Z\) impurity line radiation has also been successfully employed to determine ion temperature,\(^{263,266}\) as well as plasma rotation produced by neutral beam injection. This technique makes use of the fact that the linewidth (and shift) is a convolution of the natural linewidth (Lorentzian) together with a Gaussian due to the Doppler effect yielding a Voigt profile. Figure 47 shows a time sequence of Ti XXI \(K_{\alpha}\)-resonance lines obtained during neutral beam heating on the PDX tokamak. At \(t = 300\) ms, deuterium neutral beams were turned on for 150 ms. The resultant heating and plasma rotation are obvious from the broadening and shifts of the spectra. By fitting Voigt profiles to the data of Fig. 47 contained between the arrows, the time history of the ion temperature is obtained (see circles in Fig. 48). The triangles are measurements obtained from hydrogen charge exchange neutrals. From the shift in the line centers one can also obtain a time-dependent measurement of the Doppler shift. The associated toroidal plasma rotation is produced by the momentum transfer of the deuterium beams.

It is appropriate to now consider some of the instrument systems employed to obtain the data discussed above. First, we will discuss pulse-height analyzer systems\(^{16,262,267}\) which are used to obtain \(T_e\) and the broad features of the impurity concentration as well as departures of the distribution function from Maxwellian. Figure 49 shows a schematic of a multidetector instrument used on PLT.\(^{262}\) The use of multidetectors together with variable absorbers and apertures permits one to lessen the distortion of the high-energy portion of the spectrum due to pulse pile up. The low-energy portion of the spectrum is observed using thin foils to pass the low-energy x rays together with small apertures to reduce the count rate to an acceptable level. The high-energy portion of the spectrum is observed using thicker foils (to

![Fig. 47. Time sequence of Ti XXI \(K_{\alpha}\)-line profiles observed in PDX discharges during neutral beam heating during the period from 300-450 ms. The solid lines represent least-squares fit of Voigt functions. The arrows indicate the limits used for the fit (after Ref. 263).](image)

![Fig. 48. Ion temperature results obtained from the observed Ti XXI \(K_{\alpha}\)-line profiles shown in Fig. 47 (circles) and from measurements of charge exchange neutrals (triangles) as a function of time (after Ref. 263).](image)
block the low-energy x rays (and larger apertures so that the count rates are equalized. By overlapping the energy range of the channels the entire spectrum is obtained. The typical energy resolution is ≈ 150 eV. The detectors are LN_{2}-cooled lithium drifted silicon or Si(Li). The x-ray photons produce electron-hole pairs in the silicon. These are essentially charge pulses which are proportional to the photon energy. The pulse-height analyzer system sorts them according to energy and each energy bin (or channel) is counted. Calibration of the system is done using a biased filament together with various targets to provide the appropriate Ka line radiation. The detectors in such systems must be shielded with boron loaded polyethylene (for neutrons), lead (for x rays), and iron (for stray magnetic fields). A similar detector system is installed on PDX 262,267 as shown in Fig. 50. Here there are five separate movable detector tubes so that several regions in the plasma can be viewed simultaneously. Each tube contains three separate Si(Li) detectors in the PDX system. Typical data obtained with this system are shown in Fig. 51. The line spectra are clearly visible in Fig. 51 together with the continuum.

Detailed impurity line shape measurements, such as those shown in Fig. 47, require considerably better resolution than available with the pulse-height analyzer system. In addition, the pulse-height analyzer system does not permit one to examine the ultrasoft portion of the x-ray spectrum. To fulfill these requirements, one often employs crystal spectrometers where Bragg diffraction of the incident x rays provides the energy discrimination.5,16,262,264 Both plane and bent crystal spectrometers are utilized. Figure 52 shows a schematic of a rotating plane crystal spectrometer such as will be used on TFTR for impurity spectroscopy.264 The use of Soller slits with their large area increase the accepted radiant flux.9

![Figure 50. Schematic of the multidetector x-ray pulse-height analyzer (PHA) system installed on the Princeton PDX tokamak (after Ref. 267).](image-url)

Figure 53 shows a schematic of a curved crystal spectrometer262 such as the one used to obtain the PDX impurity line shapes shown in Fig. 47. The radius of the focal circle is one-half that of the bent crystal. The PDX spectrometer employed a quartz crystal cut in the 2023 plane which provided a resolution of \( \lambda / \Delta \lambda = 23000 \). A multiwire gas (typically 90% Ar, 10% methane) proportional counter is employed for x-ray detection. Two planes of cathode and anode wires are oriented perpendicular to each other. The incident x-ray photons produce fast primary photoelectrons which subsequently produce secondary electron-ion pairs. The resultant electron avalanche produces a charge pulse to the anode wires, thereby inducing an image charge on the cathode.
FIG. 51. Typical x-ray spectra obtained from simultaneous measurements along four chords in PDX in Fig. 50. The smooth curves represent polynomial fits to the continuum neglecting the C Kα, Ti Kα, and Ti Kβ impurity peaks (after Ref. 267).

wires. The cathode wires are capacitively coupled to a delay line and the difference in arrival times of the charge pulses at the two ends of the delay line is proportional to the position at which the photon struck the detector, thereby providing energy resolution. Typical time resolution of such systems is \( \approx 20-50 \) ms.

As mentioned earlier, x-ray measurements have been extensively employed in the study of MHD instabilities and disruptions. The success of these measurements has led to the development of x-ray imaging systems. Figure 54 illustrates the arrangement employed for imaging measurements on the PLT and PDX tokamaks. The plasma x-ray emission is viewed through a slot aperture via an array (see Fig. 55) of silicon surface barrier detectors or PIN diodes operated in the current mode. Since the plasma is optically thin, one obtains the chordal emission on each channel. The unfolding of these data to obtain the local fluctuations in emissivity will be described later. Since the currents generated by the detectors range from only several nanoamperes to tens of microamperes, one requires low noise, wide band amplifiers. The units employed for the PLT/ PDX measurements have noise levels of \( \approx 0.3 \) nA for a bandwidth of 500 kHz. The detectors are relatively broadband and can be used in the range of \( \approx 0.1-20 \) keV dependent upon choice of foil.

Fig. 52. Schematic of the rotating crystal x-ray spectrometer installed on the Princeton PLT tokamak together with a photograph of the crystal and detector assembly (after Ref. 262). The x rays emitted from impurity ions located in the shaded regions of the tokamak plasma enter the helium-filled chamber through a Be window and after Bragg reflection from the crystal impinge upon a position sensitive multiwire proportional counter.

Fig. 53. Schematic of the Johann x-ray crystal spectrometer used on the Princeton PLT and PDX tokamaks (after Ref. 262). The x rays emitted from impurity ions located in the shaded regions of the tokamak plasma enter the helium-filled chamber through a Be window and after Bragg reflection from the crystal impinge upon a position sensitive multiwire proportional counter.
FIG. 54. Schematic of the x-ray imaging system employing an array of silicon surface-barrier detectors as used on the Princeton PLT and PDX tokamaks (after Ref. 262).

Typical data from the imaging array is shown in Figs. 56 and 57. Figure 56 displays the chord-integrated x-ray emission observed with the $r = 0$ cm detector. As can be seen, a large increase in the signal occurs during neutral beam heating. An expanded time sequence is shown in Fig. 57 from several of the detectors showing the growth of the $m = 1$ oscillations.

It is appropriate now to discuss the techniques for interpreting and unfolding the chord-integrated emissivity measurements. A familiar technique from interferometry and spectroscopy is Abel inversion. However, in the standard approach one normally makes the assumption of circular symmetry which is certainly not valid for the helical perturbations found in tokamaks. Fortunately, the tokamak observations have shown that the oscillation frequency is independent of radial position. This indicates that the perturbation exhibits a rigid body rotation. Therefore, the authors of Refs. 268 and 269 Fourier analyze the chord intensity in time and interpret the components as $m$ modes of a Fourier expansion in the poloidal angle $\phi$. Detailed discussions of the reconstructions are contained in Refs. 268 and 269. Figure 58 shows a time sequence of the $M = 2$ perturbation during a minor disruption obtained using these techniques.

Reconstruction techniques other than the generalized Abel inversion have been pursued.283 Specifically, Chase et al.283 have discussed the use of computerized tomography techniques for x-ray image reconstruction. Since one is limited as to the number of viewing angles (and chords), a question arises as to the applicability of the convolve and back-project algorithm.283 Although even the most accessible of present day fusion devices do not allow for the optimal number of views, the authors show that suitable systems could provide considerable information. Specifically, they have calculated the response of various systems to hypothetical sources such as the disk-ellipse test pattern shown in Fig. 59(a). For example, Fig. 60 shows a proposed seven-view imaging system which they designed for use on the Alcator-C tokamak. This also required a new reconstruction algorithm since the views are not spaced uniformly around the plasma. The resultant reconstruction is shown in Fig. 59(b). Unfortunately, the system installed on the Alcator-A tokamak permitted much more limited viewing. Therefore, the

FIG. 55. Photograph of surface barrier detector array used for x-ray imaging the Princeton PLT tokamak plasma through a slot aperture (after Ref. 268).

FIG. 56. Central trace from Princeton PDX soft x-ray imaging system showing $m = 1$ oscillations during neutral beam heating (courtesy of K. W. Hill).

FIG. 57. Expanded time history of x-ray emissivity from several radial chords of PDX showing $m = 1$ oscillations during neutral beam heating (courtesy of K. W. Hill).
above-mentioned techniques could not be directly applied. These authors instead made use of the rigid rotation discussed earlier to justify a technique they refer to as rotational reconstruction. Basically, they use a stationary group of detectors and obtain views from different directions by equating them to views at different times. A reconstruction of a rotating $m = 2$ emissivity feature is shown in Fig. 61.

Recently, Granetz has discussed the possibility of an x-ray tomography diagnostic for the Alcator-C tokamak using multielement solid-state detector arrays. The proposed system would consist of 21 “fans”, each with eight chords, providing a total of 168 views. The calculated spatial resolution of the x-ray emissivity cross section is $3 \times 3$ cm.

VII. NEUTRAL PARTICLE DIAGNOSTICS

Neutral particles provide a convenient method of determining plasma properties in magnetic fusion plasmas. In particular, plasma ion temperatures are routinely obtained via the analysis of energetic charge exchange neutrals which escape the plasma volume. Measurements of the attenuation of neutral beam probes can provide line integral plasma densities.

The principles governing the use of charge exchange neutrals for the determination of plasma ion temperature are quite simple to understand. The technique first depends on the presence (within the plasma) of atomic neutrals of the ion species whose temperature is to be determined. For present devices, the atomic density is $\approx 10^7$ to $10^9$ cm$^{-3}$. The relatively large reaction rate for charge exchange ($\approx 10^{-7}$ cm$^{-3}$/s) results in energetic particle generation rates of $\approx 10^{14}$ to $10^{15}$ cm$^{-3}$ s$^{-1}$. By measuring the energy distribution of the es-
In the proposed seven-view x-ray imaging system for Alcator-C, one can determine the ion temperature. However, there are several problems. First, the escaping distribution of neutrals suffers significant attenuation passing through the plasma. Therefore, one primarily tends to obtain information on the cooler, outer region of the plasma. Exponential attenuation lengths are \( \sim 10 \text{ cm} \) so that there are significant difficulties in extending this technique to large devices such as TFTR \((a = 90 \text{ cm})\). Therefore, alternative ion temperature diagnostics such as laser scattering are currently under investigation (see Sec. III). An additional problem with this diagnostic is that the neutral atomic density peaks toward the outside of the plasma further weighting the low-energy portion of the spectrum.

The above-mentioned difficulties led to the use of so-called active neutral beam techniques. Here one injects a beam of the appropriate atomic species to artificially enhance the neutral density. By modulating the source and using synchronous detection, one can further enhance the signal levels.

Figure 62 shows a representative neutral analyzer system. The multichannel system was used on the Lawrence Livermore Laboratory mirror machine 2XIIIB in which the neutral beam injection greatly enhanced the charge exchange flux. The system is mounted on a trolley which facilitates radial scans of the plasma. The escaping neutrals enter the gas stripping cells where they are ionized at 0.1\%-10\% efficiency. The 10-kG magnet provides momentum analysis of the ions. This is then followed by 15 electrostatic energy analyzer and detector channels located over the interval from 18° to 108° of particle deflection. The combination, therefore, permits one to select neutrals of a particular mass and to time resolve 15 points on their energy spectrum in the range 0.5-40 keV. This is particularly important for mirror devices due to the essentially non-Maxwellian nature of the plasma.
ion distribution. Figure 63 shows a schematic of the 90° deflection parallel plate electrostatic analyzer and detectors used in the E11B neutral analyzer system (after Ref. 285).

The charge exchange diagnostic system for TFTR is considerably more involved due to the increased size, multi-ion species operation, and rf neutral beam heating scenarios. There will be an arrangement of 12 perpendicular analyzers for determination of the radial ion temperature profile and eight "tangential" analyzers to measure the slowing-down energy distribution of beam injected ions. The mass and energy of the particles are determined using an E11B analyzer. The basic arrangement is shown in Fig. 64, where H\(^+\), D\(^+\), and T\(^+\) are analyzed. More details of the channel electron multiplier are shown in Fig. 65. There are 75 energy channels for each species with an energy resolution which varies from 0.8% - 4% depending upon the anode pad location. With plate potentials of 0.6 - 6 kV, the ions can be resolved over an energy range of \(\pm 0.5 - 150\) keV. Since TFTR will produce large fluxes of 2.5- and 14-MeV neutrons, studies have been performed on the neutron response of the chevron microchannel plate. The measured response was 1.7 \(\times 10^8\) and 6.4 \(\times 10^{-3}\) counts/neutron, respectively, which is to be compared with the ion flux levels of \(10^6 - 10^8\) ions/cm\(^2\)/s which are detected with nearly 100% efficiency.

**VIII. NEUTRON DIAGNOSTICS**

As the temperature and density of mirror and tokamak plasmas continue to increase, the neutron yield arises correspondingly. It is therefore not surprising that there is increasingly more interest in neutron diagnostics. This is fortunate because, as noted in Sec. VI, present diagnostics such as charge exchange determination of ion temperature suffer limitations in these large plasmas.

Neutrons are produced by both D-D and D-T plasmas by the reactions

\[
D + D \rightarrow n(2.5\text{ Mev}) + 3\text{-He (0.8 MeV)}
\]

and

\[
D + T \rightarrow n(14.1\text{ MeV}) + 4\text{-He (3.5 MeV)}.
\]

Since tritium is also a product of the D-D reaction, there will also be a population of 14.1-MeV neutrons due to the resulting small D-T component. The neutron yield \(Y_n\) is given by

\[
Y_n = \frac{1}{\pi} \int dt n^2(r) \langle \sigma v \rangle,
\]

where \(\langle \sigma v \rangle\) is averaged over the particle distribution. Thus, if the density profile and neutron yield are known, one may determine the mean ion energy. Neutron production rates in present devices (i.e., PLT, TMX, etc.) range up to \(\pm 2 \times 10^{13}\) n/s. This is anticipated to increase up to \(10^{19}\) n/s in intense neutral beam heated TFTR discharges. These enhanced neutron yields have led to the feasibility of energy resolved measurements. For a Maxwellian distribution, \(\langle \sigma v \rangle \propto T^{-2/3} \exp(-\beta/T^{1/3})\). This shows how sensitive the neutron yield is to the presence of a small, high-energy tail and why one must take care in interpreting neutron measurements. Such a problem was in fact found in the neutron measurements on the Zeta device. The width of the neutron spectrum (in keV) for an isotropic, Maxwellian plasma is given by

\[
\Delta E_n \approx 82.5 \sqrt{T_i(\text{keV})} \quad \text{for D-D}
\]

and

\[
\Delta E_n \approx 180 \sqrt{T_i(\text{keV})} \quad \text{for D-T}.
\]
Thus, ion temperature can be directly determined from Doppler widths. The above relations also place constraints on the acceptable neutron spectrometer linewidths. The decay of the neutron yield can also be used together with the density decay rate to determine the energy confinement time $\tau_E$ through the relation:

$$\frac{1}{\tau_N} = \frac{2}{\tau_P} + \frac{3.5}{\tau_E}.$$  

(25)

There are several complicating processes which serve to confuse the above measurements. Scattered neutrons (from support structures, etc.) complicate flux measurements. The use of collimated detectors helps to minimize this error. Neutrons can also be produced by reactions other than D-T and D-D. For example, energetic electrons (such as runaways) can lead to neutron production via the direct electron dissociation of deuterium. In addition, the associated bremsstrahlung emission can lead to direct photodissociation of limiter materials resulting in an additional source of neutrons. The energetic photons can also lead to signals (or even saturation) in some detectors.

Despite the above caveats and difficulties, neutron diagnostics have proven to be extremely valuable. For example, the nuclear emission profile was measured using nuclear emulsions in the beam-heated PLT device by Strachan et al. In this technique, the neutron paths were mapped back into the plasma utilizing the recoil-proton tracks inside the emulsion. The resulting neutron emission profiles were strongly peaked on axis lending strength to the interpretation of the neutrons as thermonuclear in origin.

A variety of neutron spectrometers have been utilized to date including He$^3$ ionization chambers, liquid scintillators, and proton recoil telescopes. The important considerations are energy resolution, efficiency, and selectivity of the detector. Therefore, in designing a system one may wish to trade off efficiency for energy resolution for the moderate neutron yields associated with present fusion devices. A common choice is the NE213 liquid organic scintillator. Here the recoil protons cause scintillations in the liquid. These respond to both 2.5- and 14-MeV neutrons and have yield resolutions of 170 and 600 keV, respectively. An example of such an instrument which will be used on MFTF-B is shown in Fig. 66. Four separate detector and photomultiplier tubes fit inside the collimator and provide spatially resolved measurements along four independent chords through the plasma. The use of a flash digitizer with correction for differential nonlinearity results in a 7-MHz pulse-height analyzer so that neutron spectra can be obtained in 50–100 ms. By spatially mapping the Lawrence Livermore Rotating Target Neutron Source it was determined that the spectrometer possesses a resolution of $1'$ FWHM which corresponds to 10 cm at the MFTF-B plasma. A point to note from Fig. 66 is that since the NE213 detector also responds to hard x ray and $\gamma$ radiation one must employ pulse-shape discrimination to eliminate this source of error. In some tokamak operating conditions such detectors can even become saturated by the hard x-ray flux associated with runaway electrons. Therefore, it is of interest to consider more selective detectors. For example, Chrien and Strachen have developed a fast detector which employs ZnO(Ga) as the scintillating material and LiF enriched to 99.99% Li$^7$ as the detector for selective detection of 14-MeV neutrons in a background of 2.5-MeV neutrons, gammas, and hard x rays.

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**Fig. 65.** Geometry of the channel electron multiplier array (CEMA) used in Princeton TFFT neutral analyzer system (after Ref. 291). The three rectangular, semicontinuous active area strips are arranged so that one of them coincides with each of the mass columns for H$^+$, D$^+$, and T$^+$.  

**Fig. 66.** Schematic of time-space-energy resolved neutron diagnostic for the Lawrence Livermore National Laboratory MFTF-B mirror device (courtesy of D. Slaughter). The collimated enclosure consists of a large surrounding water tank with a long polyethylene tube providing a cylindrical aperture for the neutron pinhole camera. After passing through the collimator, the neutrons impinge upon the NE213 liquid organic scintillator detector which is mated to a selected photomultiplier tube. The electronics provide pulse-height analysis and discrimination against other events such as $\gamma$-ray interactions.
The full range of usefulness of neutron diagnostics can best be demonstrated by referring to Fig. 67 which contains a schematic of the neutron diagnostics system for TFTR. Activation foils can be pneumatically transported to stations located around the toroidal and poloidal circumferences of the torus, thereby providing information on the toroidal uniformity of the neutron production as well as the radial position of the fusing portion of the plasma. A nine-channel collimator system (see bottom of Fig. 67) provides time-resolved measurements of the radial profile of the neutron production. An interesting feature of this system is that it employs the concrete floor as the front shield. On the right-hand side of Fig. 67 is seen a neutron collimator–spectrometer which provides temporally and spatially resolved neutron spectra with high-energy resolution. The spatial resolution is ≈ 10 cm. Complete spectra will be obtained every 30 ms with the background gamma and scattered neutron contributions reduced to less than 1%. Finally, not shown in Fig. 67 is a paired U\(^{235}\) and U\(^{238}\) fission detector system which is employed to provide the time resolution required for interpretation of the neutron activation results.

**IX. LASER FLUORESCENCE SPECTROSCOPY**

The advent of tunable dye lasers has significantly enhanced spectroscopic plasma diagnostic applications. In particular, it has allowed spatially and temporally resolved measurements of impurity ion populations as well as neutral hydrogen/deuterium concentrations in fusion plasmas. These measurements are essential to an understanding of the power balance in both present and, more importantly, future fusion devices. The ignition temperature for a pure mass. These measurements are essential to an understanding of the source and mechanism of impurity introduction is essential to the eventual success of fusion.

**A. Impurity studies**

The processes responsible for the release of impurities in fusion plasmas include evaporation, sputtering, and arcing at both the limiter and vessel wall. Laser fluorescence experiments have permitted temporally and spatially resolved measurements of both the density and velocity distribution of such impurities. As described earlier, an understanding of the source and mechanism of impurity introduction is essential to the eventual success of fusion. Laser fluorescence has already made significant contributions towards this goal as well as monitoring the effectiveness of divertors and any deleterious effects caused by auxiliary heating.

Measurements in fusion plasmas have only been achieved recently and were preceded by a number of careful laboratory test experiments. The first measurements reported in fusion plasmas were made on the ISX-B device at Oak Ridge National Laboratory. Schweer et al. studied the density and velocity distribution of neutral Fe atoms in the plasma edge using a flash lamp pumped dye laser. They saturated the transition at 3020.64 Å and then observed the
decay of the excited state to a metastable level at 3820.4 Å. However, detection sensitivity limited their measurements to the current rise and afterglow periods and so prevented any detailed studies during the main discharge.

In contrast, Muller et al. performed measurements in ISX-B with sufficient overall sensitivity to observe both FeI and TiII during the main tokamak discharge.26 They used a Nd-YAG pumped dye laser which, coupled with efficient frequency upconversion techniques, provided a wavelength coverage of 2000–8000 Å. The experimental arrangement is illustrated schematically in Fig. 68. The photocell is used to monitor variations in the laser power and a hollow cathode lamp is adopted for wavelength calibration. The laser beam travels approximately 5 m to the ISX-B vacuum vessel and is then deflected onto a vertical path which passes through the 27-cm minor radius plasma approximately 1 cm inside the nominal plasma edge at the machine midplane. Calibration of the scattering system for FeI was achieved in two steps. First, an absolute calibration of number density was achieved in the laboratory using an ohmically heated effusive oven. This calibration was then transferred to the tokamak using Rayleigh scattering from N₂ in both the laboratory and tokamak environment. The same coaxial detection optics were used in both situations. Coaxial detection has the advantage of requiring only one port and possessing a high rejection against background emission. However, spatial optics were used in both situations. Coaxial detection has the advantage of requiring only one port and possessing a high rejection against background emission. However, spatial resolution is inferior to observation at 90°. The main conclusions from their data were deduced from a series of so-called plasma shift experiments where the plasma at a prescribed time is moved horizontally a few centimeters from its nominal center position. This movement caused the plasma to interact more closely with either the inner or outer limiter depending on the shift direction and so an assessment of the source of the dominant metal impurity (FeI) could be performed by variation of the limiter material. The outer limiter was either TiC or stainless (SS), whereas the inner limiter was always SS. The vacuum vessel was also SS. The shift experiments illustrated that the FeI impurity level remained virtually unchanged when the interaction between limiter and plasma was varied. In contrast, the TiII impurity level substantially increased when the plasma was shifted 1.5 cm toward the TiC limiter (see Fig. 69). These results firmly established that the dominant FeI impurities were due to interactions with the wall and not the limiters. Further detailed physics arguments strongly suggested that the generation mechanism for the FeI impurities was sputtering of the vacuum walls by charge exchange neutrals and not hydrogen ions or impurity ions.

Following these early measurements at ISX-B, a number of impurity studies have been performed in fusion plasmas in Europe and the USA. An excellent review of the current status may be found in Ref. 320. We shall conclude the discussion of laser fluorescence impurity studies by describing recent work of the Julich Group in Germany. They have performed experiments on the ASDEX tokamak to study the release, transport, and ionization of Ti atoms at the machine’s divertor plates.320–327

The experimental arrangement is illustrated in Fig. 70. Two separate laser systems have been utilized. The majority of the measurements were performed using a flash lamp pumped dye laser producing output powers of ~800 W in the UV with pulse lengths of ~0.5 μs. The bandwidth of the laser was ~0.08 Å but repetition rates were less than 1 Hz. In order to study the temporal development of the Ti fluxes (on a single shot), a high repetition rate (~100 Hz) excimer laser pumped dye laser was also used. This laser possessed

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**Fig. 68.** Schematic diagram of the ISX-B dye laser induced fluorescence experiment. The detected volume element is approximately 3 cm long and 3 mm in diameter. By translating the optics the detection volume can be scanned from –10 to + 30 cm relative to the machine midplate (after Ref. 326).

**Fig. 69.** (a) Measured FeI vs the plasma horizontal position shift at 105 ms in a D₂ ohmically heated discharge. (b) Measured FeI (plasma centered) and TiII (plasma centered) for case D and plasma out 1.5 cm for case E vs time (after Ref. 326).
powers in the UV of up to 40 kW with pulse lengths of 15 ns and a bandwidth of 0.06 Å. The fluorescence measurements were performed in the upper limiter divertor chamber of the ASDEX tokamak. The laser beam was incident directly on the divertor plate illuminating 1 cm².

The enhanced fluorescence was observed perpendicular to the laser beam and focused onto a movable slit which allowed spatial scans in a direction perpendicular to the plate with a resolution of 2 mm. Spatial scans parallel to the divertor plate were achieved by scanning the incident laser beam with a rotating mirror. Plasma radiation and stray laser light were reduced using a narrowband interference filter.

In order to measure the Ti density in the ground state the laser was tuned to 2941.9 Å (μ3F2→ν2F2). The fluorescence, however, was observed on a visible transition from the excited state at 4453 Å (ν3F2→b2F2).

Absolute density calibration was achieved by observing the fluorescence emission from evaporated Ti within the divertor. The density of the vaporized Ti was known from the measurement of Ti collected on a quartz microbalance. It should be noted that since evaporated and sputtered Ti atoms possess different distributions in their ground-state sublevels, a correction has to be applied. The total estimated error in the density measurement is approximately a factor of 2. Using the described calibration procedures, the average density of neutral titanium was determined to be ~2×10^8 cm⁻³ (assuming sputtering was dominant). In order to confirm that sputtering was the dominant mechanism, the velocity distribution (or wavelength spectrum) was measured again during evaporation and full tokamak discharges. The results are shown in Fig. 71. The upper profile spectral width is characteristic of the laser bandwidth. Note there is a slight shift to the blue caused by the small mean velocity in the direction of the divertor plates. The lower profile is more complicated and is interpreted as a convolution of two profiles shifted both to the red and to the blue. The red shift (the larger fluorescent intensity) indicates particles emitted from the divertor plate with a velocity of 3×10⁵ cm/s. The blue shift results from the resonance of laser light reflected (R~0.25) from the divertor plate with the same emitted particles. The reduction in fluorescent amplitude is consistent with the reflection coefficient of Ti. The measured mean velocity of 3×10⁵ cm/s lies within the same range measured for Ti in ion beam sputtering experiments confirming the assumption of sputtering and not evaporation as the main generation mechanism.

The spatial variation of Ti density perpendicular to the plate is illustrated in Fig. 72. The observed ionization length of 4.3 mm is a result of electron impact ionization. Using this measured value together with a known electron density allows the determination of electron temperature using ionization rate coefficients from the literature. The derived value of 10 eV allows corrections for ionization losses to be applied resulting in Ti densities of 6×10¹⁴ cm⁻³ and fluxes of ~2×10¹⁶ atoms/cm²s.
The time development of the Ti flux is illustrated in Fig. 73. The high repetition rate excimer laser pumped dye laser was used in these measurements. Operation with a 50-Hz repetition frequency indicated strong fluctuations of the fluorescence signals. The data points indicated in Fig. 73 are averaged over 60 ms and the solid curve is a further average over three neighboring points. The exact reasons for these temporal fluctuations are not fully understood but may be related to variations in the flow or position of the scrape off layer.

**B. Neutral hydrogen studies**

As discussed earlier, a knowledge of the concentrations and spatial distribution of neutral hydrogen in fusion devices is important in understanding the overall power balance. High-Z impurities generated by sputtering of high-energy charge exchanged H atoms can be a source of large radiation loss. The impurity studies on ISX-B, for example, have confirmed the above generation mechanism. In addition, the recycling of hydrogen near the vacuum vessel wall also needs the careful spatially and temporally resolved studies possible with the laser fluorescence scattering technique. Finally, a measurement of the velocity distribution of neutrals in the center of fusion devices may allow a spatially resolved determination of ion temperature through the charge exchange process. Such a measurement will be essential in machines such as TFTR and JET.

The above motivations have produced a number of proposals to perform scattering experiments from hydrogen. In this way the ground state is probed directly and such experiments have recently been performed as described below.

Historically the first-ever laser fluorescence scattering measurement in a fusion device was performed on the FT-1 tokamak in Russia by Razdobarin et al. The experimental arrangement is illustrated in Fig. 74. The flash lamp pumped dye laser (cresyl violet in ethanol) produced 10-μs pulses with an energy of 0.08 J. The laser bandwidth was ~8 Å and the power density was ≥ 600 W/cm² Å which was sufficient to saturate the entire H₃ (3→2) transition. In order to reduce the background H₄ noise emission, a two-channel compensating optical and detection system was employed. The collection optics viewed two adjacent regions of plasma only one of which was illuminated by the incident laser beam. A mirrored prism finally separated the signals emanating from these two regions onto two photomultiplier tubes. The resultant electrical signals were then subtracted and the result displayed on an oscilloscope. The fluctuations in H₄ background emission were therefore, substantially reduced resulting in an order of magnitude improvement in the fluorescent signal-to-noise ratio.

A rough idea of the magnitude of the fluorescent signal can be obtained by adopting a coronal plasma equilibrium model. The initial ratio of populations is then given by

\[
\frac{N_3}{N_2} = \frac{R_{13}}{R_{12}} \frac{A_{21}}{A_{31} + A_{32}}.
\]

where \( R_{13}/R_{12} \) is the ratio of transition rates for excitation by electron impact from the ground state and \( A_{21}/(A_{31} + A_{32}) \) is the ratio of probabilities of radiative transition from the second and third levels.

Using tabulated cross sections Eq. (26) gives \( N_3/N_2 \approx 0.85 \). If we now assume the incident laser beam saturates the transition, the ratio \( N_3^2/N_2^2 \) is given by the ratio of statistical weights \( g_3/g_2 \) which for hydrogen is 2.25. There-
fore, a reasonable (although not large) enhancement in the fluorescent emission can be expected. In order to determine the densities of hydrogen atoms from the fluorescent signal, two main assumptions were made. First, the laser power was assumed sufficient to saturate the excitation transition. Second, the laser pulse length was assumed long compared to the lifetimes of the excited transition levels.

Using these assumptions together with plasma models described in Refs. 322 and 323 the following coefficients can be calculated:

\[ \alpha = \frac{N_2}{\Delta N_2} \text{ and } \beta = \frac{\Delta N_2}{N_1 n_e} \]

The results are shown in Fig. 75. From the measured \( \Delta N_2, n_e \) the populations of atoms in both the ground and excited state can then be derived. The results are illustrated in Fig. 76 at different instants of time. Razdobarin et al. estimate the total error in their derived density profile to be 25%. However, it should be noted that in higher-density plasmas (\( > 10^{13} \text{ cm}^{-3} \)) the coefficients \( \alpha, \beta \) become very sensitive to the exact local value of density and so the total error can be substantially higher. The coefficient \( \alpha = \frac{\Delta N_2}{N_2} \) is also very sensitive to temperatures below 100 eV and so neutral densities derived may have large errors in the plasma edge region.

Recently an \( H_\alpha \) scattering experiment was performed on a reverse field pinch device at the Culham Laboratory.336 Experimentally similar techniques were used in the earlier Russian experiment with regard to background compensation. A collisional–radiative model developed at Imperial College324 was used to predict the relative populations of excited states prior to, during, and after laser irradiation. Using this model the increased population of \( N_3 \) can be calculated as a function of the ground state density for a particular \( T_e \) and \( n_e \). The calculated result for \( n_e = 10^{13} \text{ cm}^{-3} \) and \( T_e = 100 \text{ eV} \) is shown in Fig. 77. The linear dependence is a result of the close coupling of the lower excited states to the

![Fig. 75. (Upper) Temperature dependence of the coefficient \( \alpha = \frac{\Delta N_2}{N_2} \). (Lower) Temperature dependence of the coefficient \( \beta = \frac{\Delta N_2}{N_1 n_e} \). The hatched areas show the range of hydrogen atom densities \( N_2 = 10^{12} - 10^{14} \text{ cm}^{-3} \) and electron densities \( N_1 = 10^{13} \text{ cm}^{-3} \) curve (1), \( N_1 = 3 \times 10^{13} \text{ cm}^{-3} \) curve (2), \( N_1 = 6 \times 10^{13} \text{ cm}^{-3} \) curve (3), and \( N_1 = 10^{14} \text{ cm}^{-3} \) curve (4) (after Ref. 335).](#)

![Fig. 76. (a) Radial distribution of excited hydrogen atom density together with electron density. (b) Radial distribution of normal hydrogen atom density. Curve 1 corresponds to 4 ms after current onset and Curve 2 corresponds to the current maximum which occurs 15 ms after current onset (after Ref. 335).](#)
ground state. The scattered signal power was calibrated using a standard irradiance source whose color temperature was accurately known. This then allowed a determination of $\Delta N_I$, assuming knowledge of the scattering volume and collection solid angle.

The derived absolute value for neutral hydrogen concentration was $7 \times 10^9$ cm$^{-3}$. However, the measured value of fluorescence signal-to-background volume emission was found to exceed theoretical expectations by a factor of our. Explanations for this discrepancy included the possibility that collision rates between atomic levels used in the C-R model were incorrect. If this is in fact true, the errors existing in $H_a$ scattering experiments could be significant. This problem clearly needs further study if accurate neutral hydrogen density profiles are to be derived from fluorescence scattering studies.

The present status, therefore, appears to be that $L_a$ scattering is not feasible due to interpretive and signal-to-noise difficulties associated with insufficient laser power to saturate the transition. The $H_a$ fluorescence experiments illustrate that signal-to-noise problems can be overcome but that interpretive difficulties still remain. Further work is necessary to clearly establish the true accuracy of such measurement.

We shall now address the possibility of determining ion temperature (or more correctly neutral temperature) by measuring the fluorescent wavelength spectrum. This possibility is under investigation by Muller et al. on the Doublet III tokamak. The method adopted is not the usual one of employing a narrow bandwidth laser to scan the wavelength spectrum which for deuterium ($D_a$) is $\sim 8$ Å (FWHM). There are two difficulties with this approach. First, if a very narrow-band source is employed, the expected signal-to-noise ratio can be shown to be very small since the number of atoms excited is significantly reduced. Second, the use of a moderately broad source (e.g., $\sim 2$ Å) can result in a distorted measurement of the line profile due to saturation effects. The approach of Muller et al. is to use an extremely broad spectral output by setting the laser grating to zero order. In this situation, the whole $D_a$ line profile is pumped uniformly into the saturation regime. The fluorescence radiation spectrum would then be detected on a single shot with a set of five narrow-band interference filters chosen to optimize signal to noise and accuracy of the line profile determination. A feasibility experiment using one of these filters has already been performed. The $D_a$ fluorescent signal as a function of time is shown in Fig. 78. Signal-to-noise levels of 6:1 were obtained which, if all five filters were employed, would allow a determination of temperature within $\pm 15\%$. A conclusive measurement is expected during 1983–1984.

X. Magnetic field measurements

As discussed earlier in Sec. II, the determination of the local magnetic field distribution in fusion plasmas is of great importance. In addition to the Faraday rotation technique described in Sec. II, there are a large number of alternative methods which can potentially provide a measurement of the internal magnetic field. These include the use of heavy ion beam probes, the nonlinear generation of $2\omega_a$, emission at the upper hybrid layer, and the measurement of the polarization of the electron cyclotron emission. The interested reader is referred to the excellent review article by Peacock contained in Ref. 339 which describes these techniques as well as many others in great detail. In addition to local magnetic fields, one often requires measurements of global magnetic fields. In this section, we will restrict our attention to only three techniques: Zeeman polarimetry, the use of heavy ion beam probes, and the nonlinear generation of $2\omega_a$.

\[ [\text{Fig. 78. Increase in population; } \Delta N_I \text{ of the } N = 3 \text{ level with ground-state level density (after Ref. 336).}] \]
of the magneto-optic properties of fused silicon single-mode optical fibers, and laser light scattering from free electrons.

Figure 79 shows a schematic of the Zeeman polarimeter developed by West et al. for use on the TEXT device. The 100-keV Li neutral beam penetrates the large plasma and Zeeman splitting of the $2^3S_1\rightarrow 2^3P_0$ transition into three primary lines occurs. The polarization of the central $\sigma$ component lies along the direction of the magnetic field. McCor-mach et al. made use of this to measure the magnetic field in the Pulsator tokamak by filtering out the $\sigma$ components. Here they were aided by the fact that in a tokamak the poloidal field is much smaller in magnitude than the toroidal field. The magnitude of the magnetic field can also be determined from the separation between the $\pi$ and $\sigma$ components. However, the requirements on Li-beam velocity spread and divergence becomes more stringent compared to the polarization measurement. For example, West, et al. calculate that $\Delta V \leq 1$ and 100 V and $\theta \leq 10$ and 50 mrad are required, respectively, for the two measurements.

The system shown in Fig. 79 solves the problem associated with the low signal levels encountered with low beam currents and population densities. The output of a tunable cw dye laser propagates collinearly with the lithium beam and induces fluorescence of the Li beam, thereby considerably enhancing the signal levels. Single-shot current profiles will be obtained by measuring the phase angle between reference laser polarization direction and the laser-induced fluorescence signal at several positions along the tokamak radii.

The third technique mentioned earlier makes use of the modulation (at the electron cyclotron frequency and its harmonics) of the light scattered from uncorrelated electrons (i.e., small $\alpha$ scattering). For scattering directions sufficiently close to perpendicular to $B$, the scattered spectrum consists of a series of peaks separated by $\alpha_{ce}$ with an overall Gaussian envelope of width $2k\nu_e$. The modulation depth maximizes for $k_{sc}B$. The problem with this technique is that the number of scattered photons in each peak is extremely small rendering it difficult (if not impossible) to individually resolve the peaks. A particularly clever solution to the problem was suggested by Sheffield who proposed that a Fabry-Perot interferometer be employed with its free spectral range adjusted to coincide with $\alpha_{ce}$ so that the peaks are effectively superimposed upon each other. The experimental realization of this technique which was used to measure the current profile in the DITE tokamak is shown in Fig. 80. The poloidal magnetic field and $\theta$ profiles obtained with this system are shown in Figs. 81 and 82. The authors estimate that the system accuracy is sufficiently high to determine the field direction to within 0.15° and $\theta$ to within 5%.

Recently, Chandler and Jahoda have begun investigations into the use of single-mode low linear birefringence silica optical fibers for the measurement (nonlocal) of fusion plasma magnetic fields. This technique makes use of the Faraday rotation produced within the fiber by a longitudinal magnetic field. For a closed loop linking a current $I$, the rotation is given by the Verdet constant $V$ which for silica at 633 nm is $\approx 4.68 \times 10^{-6}$ rad A$^{-1}$. Thus far, Chandler et al. have demonstrated by the feasibility of the technique in sev-

![Fig. 79. Schematic diagram of the collinear beam Zeeman polarimeter for current-density profile measurement in the TEXT tokamak (after Ref. 392).](image-url)
eral laboratory test configurations. In addition, they have succeeded in sensing the 250-kA current of a dense Z pinch.

XI. EXTREME ULTRAVIOLET DIAGNOSTICS

In addition to the methods described in the preceding, both visible and vacuum ultraviolet spectroscopy are extremely valuable diagnostic techniques. However, because of the excellent references available on these topics we will only briefly discuss some recent developments in extreme ultraviolet (EUV) diagnostics. Specifically, we will describe the recent developments in time- and space-resolved extreme ultraviolet diagnostics.

Richards et al. have performed spatial EUV imaging by forming a spatial image using the slit of a monochromator. As they point out, this is a one-dimensional analog of a pin hole. Their detection system employed two microchannel plates (MCP's) with an inner aperture diameter of 18 mm. Above the first MCP is a biased wire mesh to prevent photoelectrons from escaping. The front MCP serves as a photocathode. The combination of two MCP's provides a

FIG. 81. The measured poloidal field $B_0(\rho)$ distribution in the DITE tokamak (after Ref. 350).

FIG. 82. Comparison of "q" distributions; $\bigcirc$ from poloidal field measurements, $--$ calculated from electron temperature profile assuming $j = T^{n/2}$ and constant $Z_{e,r}$ in the DITE tokamak (after Ref. 350).
gain of $\simeq 10^5$ which is high enough for pulse counting and low enough to prevent saturation of the second MCP. A ceramic header unit supports the 22 parallel anode strips which feed a set of pulse-amplifier–discriminators whose outputs are fed into counters and memory buffers. Finally, the buffers are strobed by a digital multiplexer to produce a string of 22 15-bit data words. Typically, this is repeated 400 times in a plasma discharge (100-μs time resolution) and the data Abel inverted to provide impurity profiles (i.e., volume emission rate profiles). The wavelength resolution is $\simeq 1$ Å and the wavelength coverage is $\simeq 300$–1700 Å for the normal incidence monochromator employed. Shorter wavelengths can be studied using a grazing incidence monochromator.

In addition to spatially resolved EUV spectra of impurity lines one requires time-resolved spectra of many lines in a single plasma discharge. To satisfy this requirement, Bell et al. have developed a multigrating instrument which provides the capability of both high-resolution operation ($\simeq 1$ Å) as well as lower dispersion ($\simeq 5$ Å) for survey studies. The detector consists of a 1024-element linear photodiode array coupled to an MCP image intensifier tube with a CuI photocathode. The integration time for each element can be as short as 3.7 ms so that many complete spectra can be obtained in a single plasma discharge.

XII. OTHER TECHNIQUES

The previous sections of this review have concentrated on those diagnostic measurements in which (in the opinion of the authors) there is the greatest instrumentation development activity. However, there are a number of techniques which are extremely useful which deserve mention. These include bolometry, neutral beam attenuation measurements, retarding grid energy analyzers, and Langmuir probes. In the following we briefly describe their use.

There are a variety of energy-loss mechanisms in magnetically confined fusion plasmas. These include cross field transport of charged particles (as well as end loss in opened devices such as mirrors), the loss of neutral atoms and electromagnetic radiation from both neutral atoms and charged particles (including impurities) with most of the radiation loss in the UV and x-ray region. To understand plasma confinement, it is important to obtain accurate measurements of the energy loss so that power balance calculations can be made. These measurements are commonly made using bolometers which are sensitive to all of the above-mentioned loss processes and simply respond to the absorbed energy. Basically, there are several ways in which bolometers are employed. First, by simply placing a number of uncollimated bolometers around the vacuum vessel wall the total power reaching the wall can be estimated. A second, more sophisticated measurement technique, utilizes collimated bolometer arrays. Here, localized radiation measurements can be made by inverting the information from a large number of chordal views using techniques such as those discussed in Secs. 1 and VI. For example, both the type and the amount of impurities present can be obtained via Abel inversion of the bolometry data.

The construction of bolometers is relatively simple. For example, the bolometers employed on PLT, PDX, and ATC consist of a thin ($\simeq 25$–38 μm) nickel or copper foil disk with a lightweight thermistor, referred to as a "thinner" by its manufacturer, mounted on the rear surface. The thermistor resistance is $\simeq 25$ kΩ and its dimensions are $0.5$ mm square with a thickness of $\simeq 0.38$ μm. However, this type of bolometer is not suitable for use on the next generation of fusion plasma devices such as TFTR because of the high anticipated temperatures and nuclear radiation levels. To surmount these difficulties, a new type of bolometer has been developed which utilizes a thin, die-cut, 3.5 μm-thick platinum grid which is sufficiently thick to even stop x rays of 9-keV energy. Although the temperature coefficient of resistance of platinum is much smaller than that of thermistor, this is nearly exactly compensated by its smaller thermal capacitance. To improve the measurement accuracy, the authors excite their Wheatstone-bridge configuration at 20 kHz and employ synchronous detection. At a bandwidth of 50 Hz the equivalent noise is $\simeq 100$ μW/cm² which is to be compared to a calculated power density at the detector of 2 mW/cm² for a 1-MW TFTR plasma.

Collimated neutral particle beams have found use in plasma density determination through measurements of the beam attenuation as it propagates along a chord through the plasma. By employing an array of beam probes the density profile can be inverted from the line averaged measurements using techniques similar to those discussed in the interferometry and x-ray diagnostic sections. In this case the fractional beam attenuation $A$ and transmission $T$ are related to the plasma density through the expression

$$1 - A = T = \exp(-\sigma_{\text{off}} \int n dl),$$

where $\sigma_{\text{off}}$ is the effective attenuation cross section (which depends on plasma parameters and a variety of interactions including charge exchange and beam ionization) and $\int n dl$ the chord integrated line density. As a representative example of a neutral beam attenuation diagnostic we will briefly describe the system used on the TMX device. Typically, deuterium atomic beams were produced with dimensions at the extractor grid of 7 × 35 cm with beam divergence angles (at the $e^{-1}$ points) of ± 0.5° in the narrow direction and ± 2.5° in the broad dimension. The detector arrays were then focused on one of these neutral beam sources.

The individual detector elements for the neutral beam attenuation measurements are simple to fabricate. Neutral particles enter a hole in the aluminum body and pass through two collimating holes (0.050-cm diameter) spaced 10.5 cm apart before impinging upon the detector emitter surface resulting in a secondary-electron current proportional to the incident neutral particle flux. Baffles located between the collimating holes prevent photons and particles which scatter from the walls from reaching the emitter surface. Signal levels on the various detectors ranged from 4.5 μA down to $\simeq 20$ nA which are to be compared with the 5- to 10-nA background noise (electronics, cabling, and EMI) and the signal-noise level due to beam variations (± 2% to 3% of signal level). Three arrays were used on TMX comprising a total of 54 detectors with two located in the end cells and
one in the central cell. The detector signals were digitally sampled at a 100-kHz rate 4096 times during the plasma discharge (typical duration of 25 ms). The resultant density profile was in agreement with a single data point from Thomson-scattering measurements.

Charged particles can be lost along the magnetic field lines in open-ended devices such as mirror machines. By using an electrostatic gridded energy analyzer, the energy distribution of the escaping particles may be determined. From these measurements the plug potential can be obtained as well as the average ion energy in the central cell. It is currently planned to supplement these end-loss measurements on TMX upgrade (TMX-U) by means of an $E/\beta$ energy analyzer patterned after that described in Sec. VII with the ions detected with a two-dimensional array of biased collector plates. Using such an end-loss ion spectrometer, it is anticipated that the thermal barrier potential can be determined as well as the potential in the transition region of the axisymmetric region.

In the case of TMX there were two end-loss analyzers positioned at each end wall. Typical current densities at the analyzer are 3 mA/cm$^2$. By setting the ion-repeller grid voltage to zero, the time history of the total end-loss current is obtained. Alternatively, by ramping the repeller voltage, the ion energy distribution is obtained. Note that since all of the ions will arrive at the wall with energy equal to or greater than the plug potential $\phi_p$, the analyzer current will be essentially constant for repeller voltage less than $\phi_p$. Typical equilibrium values of $\phi_p$ were 300–800 V with an estimated measurement accuracy of ±50–70 V.

Finally, Langmuir probes continue to find use in magnetically confined fusion plasmas. For example, Langmuir probes have been employed to determine the parameters of the halo (or edge plasma) in TMX-U. In addition, edge density fluctuations in tokamak plasma have been studied using Langmuir probe arrays. Details of Langmuir probe measurements may be found in Refs. 8 and 9.

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