Laser-aided plasma diagnostics

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ABSTRACT
Laser-aided diagnostics are widely applied in the field of high temperature plasma diagnostics for a large variety of measurements. Incoherent Thomson scattering is used for highly localized measurements of the electron temperature and density in the plasma. Coherent Thomson yields information on the fast ion population in the plasma and/or depending on the geometry and wavelength chosen electron density fluctuations. Interferometry and polarimetry are often combined in a single diagnostics set-up to measure the electron density and the component of the magnetic field parallel to the laser chord. Density fluctuations can be measured by means of phase contrast imaging, scattering and various other laser-aided techniques. In this paper a brief tutorial introduction in each of the techniques is given, followed by a description of some typical implementations on magnetic confinement devices and some examples of recent experimental results. For each of the techniques also the potential application to the ITER tokamak is discussed.
1.  INTRODUCTION

Active diagnostics that feature lasers as the probing source have a number of distinct merits: 1] the laser beam can be focused in the plasma, resulting in a good spatial resolution; 2] the measurements do not perturb the plasma because of the relatively small interaction cross sections; 3] lasers have a high spectral brightness; 4] both with pulsed- and cw-laser systems a good temporal resolution can be obtained; 5] the lasers (and in many applications also the detectors) can be positioned far from the plasma, where they can be more easily maintained.

Various types of laser-aided plasma diagnostics exist; all based on different physical interactions between the electromagnetic wave from the laser and the plasma. In general one can distinguish interaction based on: 1] absorption and/or emission, 2] refraction, 3] changes in the polarization ellipse and 4] scattering. An overview of the plasma parameters that can be measured with the various laser diagnostics is given in the parameter/technique matrix in Table 1.

Absorption and re-emission find their application in laser induced fluorescence (LIF). The wavelength of the laser in this application is tuned such that electrons in bound states of atoms or ions are pumped from energy level \( E_1 \) to energy level \( E_2 \). The excited electrons then spontaneously decay to a lower level \( E_3 \) (which may be equal to \( E_1 \)), while emitting a photon after some delay. The spontaneous emission is called fluorescence, and the intensity of the fluorescence signal can be used to determine various plasma parameters of the neutral or the partly ionized atoms in the plasma. Since
LIF is closely related to spectroscopy, it is described elsewhere in this issue.\(^1\) For the same reason it is also not included in Table 1.

A well-known method to measure the line-integrated density in a plasma is interferometry. This technique is based on measuring the phase shift that an electromagnetic wave experiences when traveling through a plasma. Birefringence in a magnetized plasma can lead to changes in the polarization ellipse of the electromagnetic beam. In principle there are two effects causing a change of polarization: the Faraday effect leading to a rotation of the polarization plane and the Cotton-Mouton effect leading to a change in the ellipticity of the probing beam. The first of these effects is often employed to measure the internal magnetic field in a hot plasma. More, recently the Cotton-Mouton effect is employed as a robust measurement of the electron density. Interferometry and polarimetry are mainly done with lasers in the far- and near-infrared, albeit that there are also systems employing microwave sources. These methods are described in Sect. 2.

Thomson Scattering (TS) refers to scattering of electromagnetic waves by free electrons. J.J. Thomson first described the theory of this process.\(^2\) Since then the scattering process has been studied in detail, including the separation between the electron and ion features by Salpeter,\(^3\) the description of collective effects\(^4\) and relativistic effects.\(^5\) The experimental application of TS as a diagnostic tool\(^6,7\) had to wait for the development of high power light sources, e.g. the (Q-switched) ruby laser in the early nineteenthies. Since then, various plasma parameters have been measured by means of this technique.\(^6,7\) The first demonstration of TS as a suitable diagnostic tool for hot plasmas was given by Peacock et al. in 1969 when they measured the electron
Further developments of TS as a diagnostic tool have led to present-day periodic-TS systems measuring the electron temperature ($T_e$) and density ($n_e$) of hot magnetically confined plasmas along the full plasma diameter, resolving up to ~100 spatial elements with time separations of ~10 µs - 10 ms.

Scattering of monochromatic electromagnetic radiation from hot electrons in a plasma, leads to a spectral broadening of the scattering spectrum due to the Doppler effect. Light scattering by an ensemble of electrons not only depends on the number of electrons but also on the interaction between them. To find the total scattered power, both the scattering cross-section of the single electrons and the phase relation between the scattered waves has to be taken into account. Correlated interactions between the plasma electrons only occur above a certain scale length, the so-called Debye length, $\lambda_D$.

\[ \lambda_D = \frac{v_e}{\omega_{pe}} \approx 7.4 \times 10^3 \sqrt{\frac{T_e}{n_e}} \text{ m}, \]  

(1)

where $v_e$ is the thermal velocity of the electrons, $\omega_{pe}$ the plasma frequency, and where $\lambda_D$ is in m, $T_e$ in eV and $n_e$ in m$^{-3}$.

Whether collective electron phenomena influence the scattering process depends on the ratio between $\lambda_D$ and the wavelength ($\lambda_0$) of the incident electromagnetic wave, which is expressed as the Salpeter (or scattering) parameter $\alpha$:

\[ \alpha = (k\lambda_D)^{-1} = \frac{\lambda_0}{4\pi\lambda_D \sin \frac{\theta}{2}}. \]  

(2)
where \( k \) refers to the differential scattering vector and \( \theta \) to the angle between the incident wave and scattered wave directions.

Depending on the value of \( \alpha \), three regions can be distinguished where the spectral distribution function of the scattered radiation shows quite different features. In the range \( \alpha << 1 \) the scattering spectrum results from scattering by individual electrons and is therefore determined uniquely by \( T_e \). This type of scattering is denominated as incoherent TS, or simply as TS.

In case \( \alpha \geq 1 \) the scattering contributions from different electrons and their surrounding shielding cloud will add up coherently since there is negligible phase difference between them. Under these conditions it is possible to deduce information about the ion velocity distribution because of the fact that ions are usually shielded by a cloud which consists of approximately 50% of electrons and 50% of (the absence of other) ions. It is the scattering by the electrons in this shielding cloud that can be observed experimentally. The first attempt to measure the ion temperature of a tokamak plasma using the collective effects of the TS process was made at Alcator-C in 1983.\(^9\) This so-called ion Thomson scattering or coherent scattering is experimentally very difficult to perform and has only been applied at a few confinement devices.\(^{10}\) Ion collective Thomson scattering is, however, still being actively explored, since it is one of the few diagnostic techniques that can be in principle used to study the fast ion distribution, and especially that of confined \( \alpha \)-particles in future fusion devices.\(^{11}\)

When \( \alpha >> 1 \) a third experimental possibility exists. In this case \( \lambda_0 \sin \theta / 2 \) is of the same order as the correlation length of the density fluctuations in the plasma (about the ion Larmor radius). As a result the scattered power is proportional to the square of the
density fluctuation level and collective motions of the electrons can be observed. This kind of scattering is also called collective scattering and is used to measure the electron density fluctuation spectrum in the plasma.\textsuperscript{12}

After a brief overview of the theory of incoherent TS in Sec. III.A, the typical experimental set-up for a TS system is presented in Sec. III.B. In Sec. III.C, a number of recent experimental applications of multi-position TS systems will be discussed. Section III.D is devoted to fast ion collective Thomson scattering with the emphasis on applications in the far-infrared and infrared. For the applications of fast ion collective Thomson scattering in the mm wave range one is referred to another paper in this issue.\textsuperscript{13}

Density fluctuations at spatial scales of the order of the ion Larmor radius and larger are ubiquitous in fusion plasmas. There is now little doubt that turbulent fluctuations arising from drift wave instabilities (Trapped Electron Modes TEM and Ion Temperature Gradient Modes ITG) are responsible for the lion’s share of heat and particle transport in the bulk of tokamak plasmas. The role of the short wavelength Electron Temperature Gradient Modes (ETG) is still uncertain. Density fluctuations also accompany other types of instabilities, including a host of MHD modes such as fast particle instabilities, Edge Localized Modes, magnetic islands and plasma waves driven for the purpose of radio frequency heating. Together these fluctuations span scale lengths in the radial and poloidal directions, ranging from the order of the plasma radius down to \( \sim 1 \text{ mm} \), fluctuation amplitudes \( \Delta n_e/n_e \) ranging from near 100\% at the plasma boundary to below \( 10^{-4} \), with associated frequencies from the kHz to the GHz range. This wide range has spurred very diverse diagnostics developments, many of which involve the use of electromagnetic wave probes, such as microwave scattering and reflectometry (see the
previous paper in this issue\textsuperscript{13}) and infrared laser scattering and imaging techniques. The latter techniques are presented in Sec. IV.

II. INTERFEROMETRY/POLARIMETRY/ELLIPSOMETRY

II.A. Theory of interferometry/polarimetry

The basic principles of interferometry and polarimetry were extensively presented in a number of review papers\textsuperscript{14,15,16,17} and textbooks.\textsuperscript{18,19} The full theoretical description of especially polarimetry is much more tedious and complicated than is described below. In this paper a choice has been made for a brief but transparent description of the basic theory. Starting point in the description of interferometry and polarimetry is the well-known Appleton-Hartree equation.\textsuperscript{20,13} This is the dispersion relation for electromagnetic waves propagating through a magnetized plasma. Since the waves are traveling close to speed of light $c$, the equation, which is an expression for the refractive index, $N$, is derived under the cold plasma assumption in which ion motion is neglected:

\begin{equation}
N^2 = 1 - \frac{X[1 - X]}{1 - X - \frac{1}{2} Y^2 \sin^2 \theta \pm \sqrt{\left(\frac{1}{2} Y^2 \sin^2 \theta\right)^2 + [1 - X]^2 Y^2 \cos^2 \theta}},
\end{equation}

where $X = (\omega_p/\omega)^2$ and $Y = \omega_c/\omega$, $\omega$ is the frequency of the propagating wave and $\theta$ the angle between the wave vector $k$ and the magnetic field $B$. The plasma frequency and the electron cyclotron frequency are given by:
\[ \omega_p = \sqrt{\frac{e^2 n_e}{\epsilon_0 m_e}} \quad \text{and} \quad \omega_c = \frac{eB}{m_e}, \]  

respectively, with \( e \) and \( m_e \) the charge and mass of the electron, \( n_e \) the electron density, \( \epsilon_0 \) the permittivity of vacuum.

For a linearly polarized probing wave traveling parallel to the magnetic field, \( B \) (i.e. \( \theta = 0 \)) through a uniform slab of plasma with thickness, \( L \), and under the assumption that the magnetic field is small but finite (i.e. \( \omega > \omega_c \) so only first-order terms in \( Y \) are retained), Eq (3) can be simplified to:

\[ N_z = 1 - \frac{X}{2} \pm \frac{XY}{2}. \]  

When entering a plasma, a linearly polarized electromagnetic wave is split into two counter-rotating circularly polarized waves, that propagate at different phase velocities corresponding to the + and – sign in Eq. (5) for the refractive index. After passage through the plasma slab the two waves will have experienced phase shift given by:

\[ \Delta \varphi = L \frac{\omega}{c} (N_z - 1). \]  

Superposition of the two circularly polarized waves will result again into a linearly polarized wave which has experienced a total phase shift of

\[ \Delta \varphi = (\Delta \varphi_+ + \Delta \varphi_-)/2 = -\frac{\omega}{2cn_e} n_e L, \]  

and whose plane of polarization has rotated by an angle

\[ \alpha = (\Delta \varphi_+ - \Delta \varphi_-)/2 = \frac{e}{2cm_e n_e} n_e BL. \]  

When the electron density becomes larger than the critical density \( n_c = \omega^2 m_e \epsilon_0 / e^2 \) the refractive index becomes imaginary and the wave cannot propagate but will be reflected.
In this regime one can diagnose the electron density profile by means of reflectometry, which is discussed elsewhere in this special issue. The rotation of the plane of polarization due to the parallel component of the magnetic field is known as the Faraday-effect.

Because the electron density and magnetic field component parallel to the probing wave are not constant along the beam path, Eqs (7) and (8) need to be integrated along the beam path. In most magnetic confinement devices the plasma is usually probed by beams that cross the plasma in a vertical plane, i.e. perpendicularly to the toroidal magnetic field (see Fig. 1). For this specific geometry Eqs (7) and (8) can be rewritten to

\[ \Delta \varphi = -\frac{\omega}{2cn_e} \int_{-z_0}^{z_0} n_e dz, \]  
\[ \alpha = \frac{e}{2cm_n e} \int_{-z_0}^{z_0} n_e B_\parallel dz, \]

where \( B_\parallel \) is the component of the magnetic field parallel to the probing beam. Eqs (9) and (10) have been derived under the assumption that the wavelengths of the probing beam is smaller than the typical scale lengths of the electron density and magnetic field.

To determine the local values of \( n_e \) and \( B_\parallel \) one must measure the full profiles of \( \Delta \varphi(R) \) and \( \alpha(R) \), which can be achieved by probing the plasma simultaneously along many different chords (and interpolating between adjacent chords). For a smooth cylindrically symmetric plasma it is possible to perform the transformation from line-averaged measurements to local quantities by Abel-inversion techniques. For plasmas which exhibit distortions from cylindrical symmetry one has to rely on numerical inversion procedures, function parametrization, or neural networks. In case the
plasma is probed from a number of independent directions, one can use tomographic reconstruction techniques to retrieve the electron density profile. The accuracy of inversion methods can be further improved by taking also data from other diagnostics like incoherent Thomson scattering into account. Because the computation time of many of the numerical analysis codes is rather limited, these techniques are well suited to be incorporated in feedback loops for plasma control.

II.B. Experimental constraints

There are various effects that could have a detrimental effect on the operation interferometry and/or polarimetry if not taken properly into account. From Eqs (9) and (10) (because of the dependence of ) it is evident that and . So the conclusion could be that the probing wavelength should be chosen to be as large as possible. However, at too long a wavelength refraction will cause the beam to deflect from its straight trajectory. As a result the beam may experience a different phase shift due to the changes path length in the plasma. Refraction gives rise to an upper limit for the wavelength to be employed in an interferometer system

with the maximum central density anticipated in the experiment and a constant which is typically in the order of 0.1 to avoid refraction. Usually, when designing an interferometer or polarimeter system extensive ray-tracing calculations are done for the envisaged plasma scenarios to properly estimate the effects of refraction.
Another instrumental effect comes from path length changes due to mechanical vibrations of for instance optical components. Vibration effects are most severe at smaller probing wavelengths, but can be largely overcome by applying a so-called two-color interferometer system.\(^{28}\) An alternative solution to overcome the effects of vibration is to employ a dispersion interferometer, in which the plasma is simultaneously probed by two beams, one of them having exactly the double frequency of the other.\(^{29}\) Other instrumental effects, which in general have a small effect on the choice of the optimum wavelength are source and detector noise.

Last but not least one should take the effect of the main toroidal field, \(B_\perp\), into account. This transverse field component makes the plasma, via the Cotton-Mouton effect linearly birefringent and causes the probing wave to become elliptically polarized. If a polarimeter system is not designed with care, the thus induced ellipticity can deteriorate the measurement of the Faraday rotation angle. De Marco and Segre\(^{30}\) have derived an upper limit for the probing wavelength below which the induced ellipticity can be safely neglected. Since the Cotton-Mouton effect is proportional to \(\lambda^3, B_\perp^2\) and \(n_e\), this reduces the upper wavelength considerably. In the past, many polarimeter systems were set up in such a way that the Cotton-Mouton effect could indeed be neglected. However, more recently it was realized\(^{31}\) that the effect can be also turned into an advantage. Since the electron density, \(n_e\), is the only unknown in the expression for the Cotton-Mouton effect, it is possible to determine this parameter by measuring the full polarization ellipse of the beam emerging from the plasma. The first successful measurement with ellipsometry (sometimes called full or complete polarimetry) was done on the W7-AS tokamak.\(^{32}\)
Most interferometer/polarimeter systems operate in the far-infrared wavelength region (100 – 500 µm). In this region and for most plasmas the phase shifts, and Faraday rotation angles are relatively large and easy to measure, whereas the deteriorating effects as refraction, ellipticity (if seen as a disturbance) and vibrations are tolerable or even negligible. Also systems have been developed in the near-infrared region (3-12 µm). Because the much higher sensitivity of these systems to vibrations, they are usually operate with two wavelengths simultaneously. In several magnetic confinement devices, the line-integrated density is measured by means of interferometry in the mm-wave length range.\textsuperscript{13}

**II.C. General implementation of interferometers and polarimeters**

Many different schemes have been developed for interferometry. The most common scheme is the Mach-Zehnder interferometer (see Fig. 2). Nowadays most interferometer/polarimeter systems employ heterodyne detection techniques,\textsuperscript{33} in which the probing beam and the reference beam have a frequency offset \(\Delta\omega\). This frequency offset (also referred to as beat frequency) can be easily introduced by either reflecting the reference beam off a rotating diffraction grating,\textsuperscript{14} or by using two separate frequency-locked sources.\textsuperscript{34} The temporal resolution that can be obtained in an interferometer system is directly related to the beat frequency between the probing and the reference beam. Mechanical modulation with a rotating grating usually limits the temporal resolution typically to several tens of kHz (values up to 100 kHz have been achieved in this way). By applying a dual-laser system, like a CO\textsubscript{2}-pumped far-infrared laser system
with two slightly detuned cavities one can easily generate beat frequencies in the MHz range\textsuperscript{35,36}

Apart from the Mach-Zehnder geometry sketched in Fig. 2, also the Michelson interferometer is often used in plasma interferometry. In this setup the probing beam is reflected after passage through the plasma by a mirror or a retro-reflector, such that it traverses the same plasma chord twice. Apart from the different optical geometries the basics of the Michelson interferometer are similar to those of the Mach-Zehnder system.

In several magnetic confinement devices a single-channel interferometer is used to probe the plasma along the magnetic axis to measure the line-averaged value of the electron density. This quantity is often used as input for real-time feedback on the electron density. If one wants to deduce the local density in the plasma (e.g. the density profile along the midplane) one must probe the plasma simultaneously along a number of chords, and apply (numerical) Abel-inversion techniques to obtain the local density values (see Sec. II.A.). In several interferometer (and polarimeter) set-ups this was done by splitting the probing beam into a number (typically 5-10) individual probing beams (see Fig. 3).\textsuperscript{37,38,39,40} The drawback of having many separate probing chords is that this results in a large number of optical components. The number of optical components can be largely reduced by probing the plasma with a slab-like or fan-like beam.\textsuperscript{16,26,41,42,43,44}

A basic feature of any interferometer system is that the phase shift (Eq. (7)) is seen by the detector as a number of interference fringes. Depending on the probing wavelength, $\lambda$, and the maximum density the number of fringes can increase from zero (before the start of the plasma discharge) to several tens or even several hundreds. The drawback of fringe counting techniques is that in case of fast density transients (e.g. due
to pellet injection, instabilities, disruptions) mistakes can be made in the number of counted fringes. Often this is evidenced by the fact that the fringe counter does not come back to zero at the end of the discharge. Various techniques have been developed to overcome the so-called fringe jumps. The most straightforward, but not fail proof, remedy is to increase the accuracy and sampling speed of the electronics. Another method that has been developed is to employ a beam that is scanned through the plasma, all the way from the very plasma edge to the plasma core.\textsuperscript{45} Although fringe jumps can still occur, the fringe counter is reset at every new scan, such that one doesn’t lose the information on the density for the rest of the discharge.

Instead of using interferometry one can also use polarimetry to measure the electron density, provided the plasma is probed tangentially to the toroidal magnetic field.\textsuperscript{46} So the component of the magnetic field parallel to the probing beam (see Eq. (10)) is the toroidal magnetic field which is known to high precision. So the only unknown in Eq. (10) is the electron density. A densitometer based on tangential polarimetry is completely insensitive to fringe jumps, since the Faraday rotation is usually well below 360 degrees. The density information is moreover weighed towards the point where the beam is tangent to the magnetic field. Another advantage is that even in shaped plasmas (i.e. non-circular elongated and/or triangular plasmas), the geometry of the experiment as seen from the top, is again circular which simplifies the Abel-inversion.\textsuperscript{47} Tangential polarimeters have been successfully applied at JT-60U,\textsuperscript{48} and LHD.\textsuperscript{49,50}
II.D. Combined implementation of interferometry/polarimetry

Hitherto we mainly discussed interferometry and polarimetry as related but independent diagnostics. However, as we will briefly discuss in this section, it is possible to combine the two diagnostics into a single system. For more detailed information one is referred to various review papers,\textsuperscript{16,17} and references therein.

A relatively simple technique to combine an interferometer and polarimeter system is to replace the beam combiner in front of the detector (see Fig. 2) by a polarizing beam splitter and to implement one additional detector per channel.\textsuperscript{14,15,37} Additionally, a rotatable half-wave plate and an additional polarizer need to be added to the system (see Fig. 5). The system is set up in such a way that one detector is mainly sensitive to the phase shift of the probing beam, whereas the other one is sensitive to the Faraday rotation. Although the Faraday rotation angle can be determined from the amplitude of the signal measured by the polarimeter detector, a much higher reliability can be achieved by determining the phase difference between the two detectors.

Many polarimetry systems and combined interferometry/polarimetry systems have been developed that are based on a modulation of the plane of polarization of the probing beam. Dodel and Kunz were the pioneers in this field. They developed a combined system in which the polarization plane of the probing beam is continuously rotating,\textsuperscript{51} and they also built a polarimeter system in which the polarization plane is modulated only by a few degrees.\textsuperscript{52} In both cases the Faraday rotation angle is deduced by means of a phase measurement. An alternative was developed for the MTX device.\textsuperscript{53} Here a beam with a continuously rotating elliptical polarization is used to probe the plasma. The charm of this system is that the interferometric phase shift as well as the
Faraday angle can be measured by a single detector per viewing chord. Also here the Faraday rotation angle is retrieved from a phase measurement. A disadvantage of this method is that there an interference between the polarimetry and the interferometry signals, which in practice means that one needs to strongly filter the interferometry signal. Another drawback, which is common for all polarization modulation and rotation schemes discussed thus far, is that mechanical methods are used for this purpose (e.g. a rotating ½ wave plate). This sets a limit on the maximum temporal resolution that can be obtained.

A method that doesn’t suffer from these drawbacks was developed for RTP. In this system, which is a generalization of the method developed by Dodel and Kunz, a triple-beam far-infrared laser system is employed with detuned cavities, such that three different frequencies are obtained. Two out of the three beams probe the plasma in orthogonal polarization states (see Fig. 6); the third one is used as reference beam. Since no mechanical modulation methods are used, it is principle possible to probe both the density and the poloidal magnetic field in the plasma with high temporal resolution (for the polarimeter up to 100 kHz). The same method has been adopted for NSTX, albeit that here a tangential setup is employed. The latter has been to study the effects of para- and diamagnetism on the toroidal magnetic field. In a spherical tokamak these are considerably larger than in a high-aspect ratio tokamak. The triple-beam system is also used at the MST reversed field pinch. With this system magnetic field measurements are routinely done with a typical time resolution of 10 µs (see Fig. 7). The current density profile can be followed during a sawtooth cycle, and is observed to peak during the ramp phase and promptly broaden at the crash. Even microscopic magnetic field fluctuations in
MST can be measured with the fast polarimeter, with the note that the typical fluctuations
levels in MST are approximately two orders of magnitude higher than in a tokamak.

As was mentioned in Sec. II.B it is possible to determine the electron density also
via the Cotton-Mouton effect. Since this measurement is, in contrast to interferometers,
not sensitive to fringe jumps there have lately been developments to combine the
diagnosis of the Cotton-Mouton and the Faraday effects into a single system. Two
methods to realize this have been proposed by Segre and Zanza.\textsuperscript{57} One of these methods
has been successfully tested in laboratory.\textsuperscript{58} At JET an approximate method has been
developed in order to extract both effects from the measurements.\textsuperscript{59} Care should be taken
in the analysis since the Faraday effect and the Cotton-Mouton effect are not completely
independent.

Despite the fact that interferometry and polarimetry are already being used on
magnetic confinement devices for 30 - 40 years, there are still continuously new
developments that aim towards higher accuracy, better time and spatial resolution and
including new measurement options.

\section*{II.E. Application to burning plasma experiments}

Interferometry and polarimetry are ideally suited for application in a burning plasma
experiment.\textsuperscript{60} For the ITER tokamak two systems are envisaged. One of them is a
poloidal polarimeter system operating at a wavelength of 118 \(\mu\text{m}\),\textsuperscript{61} the other is a
tangential polarimeter system operating at 10.6 \(\mu\text{m}\).\textsuperscript{62} The systems have completely
different aims. The objective of the poloidal polarimeter system is to measure the current
density profile. The system is set up as a Michelson interferometer with ø 37 mm retro-reflectors indented about 25 cm in the blanket modules at the high field side of ITER, and has a fan of nine viewing chords via one of the equatorial port plugs (see Fig. 8). A recent study has strongly hinted that erosion and deposition effects on the retro-reflectors indented in the blanket modules in the high field side will not have a significant effect on the operation of the poloidal polarimeter system. The electron density profile can be measured along the same nine chords. Because of the relative long wavelength of 118 µm, this can be done via the Cotton-Mouton effect (offering robustness (no fringe jumps, not sensitive to vibrations), a good time resolution and accuracy) or via normal interferometry (high time resolution and accuracy but vulnerability to fringe jumps and vibrations). An automated alignment system is proposed to keep the beams directed onto the retro-reflectors. This is done by means of a scanning mirror, which will automatically scan the beam over an area of 10 x 10 cm² in case the signal level drops below a certain threshold value.

The tangential polarimeter system is set up as interferometer/polarimeter combination with five lines of sight in the equatorial plane. Because of the much longer path length through the plasma a small wavelength is proposed to cope with refraction. Two wavelengths (10.6 and 9.27 µm) need to be applied in collinear beams to compensate for vibrations. A mirror in the entrance port is used to spatially separate the five beams. The beams are reflected back along the same beam path and leave the tokamak via the same port as the entrance beam. As for the poloidal polarimeter a real-time alignment control system needs to be employed to keep the beams focused onto the reflectors throughout the whole ITER discharge. Because the required optical tolerance of
the retro-reflectors is much higher than that of the poloidal system, the retro-reflectors need to be positioned behind relative long ducts. Still deposition on the retro-reflector could have a deleterious effect on the operation of the interferometer/polarimeter,\textsuperscript{63} and much further research should be devoted to finding the optimum set up.

III. THOMSON SCATTERING

III.A. Thomson Scattering Theory

The mathematical framework of Thomson-scattering theory is rather complicated and has been extensively described elsewhere.\textsuperscript{3,18,64,65} Therefore, we shall restrict ourselves to discussing the basic ingredients.

First consider a single electron of mass $m$ and charge $e$, located at position $r$ with respect to a coordinate system. The position vector $\mathbf{R}$ denotes the location of the detector (Fig. 9). The amplitude vector for the electron position $|r|$ is assumed to be much smaller than $|\mathbf{R}|$. The incident wave with amplitude $E_0$, propagation vector $\mathbf{k}_0$ and angular frequency $\omega_0$, can be written as:

\begin{equation}
E = E_0 \cos(\mathbf{k}_0 \cdot r - \omega_0 t).
\end{equation}

Under the action of this field, the electron experiences a net acceleration:

\begin{equation}
\ddot{r}_e = \frac{e}{m} E_0 \cos(\mathbf{k}_0 \cdot r - \omega_0 t).
\end{equation}

In the far field zone ($|r| >> \lambda$) for $dr_e/dt << c$, the radiation emitted from an accelerating charge is given by:
\[ E_s(R, t') = \frac{e}{4\pi \varepsilon_0 c^2 R} \left( \hat{r}_s \times \hat{k}_s \right) \times \hat{k}_s, \]  
\[ (14) \]

where \( \hat{k}_s \) is a unit vector in the direction of the scattered field direction and \( t' \) is the retarded time:

\[ t' = t - \frac{k_s \cdot r}{c|\hat{k}_s|}. \]  
\[ (15) \]

Combination of Eqs. (14) and (15) yields the field of the scattered radiation:

\[ E_s(R, t) = \frac{e^2 (E_\theta \times \hat{k}_s \times \hat{k}_s)}{4 \pi \varepsilon_0 mc^2 R} \cos(k_\theta \cdot r - \omega_0 t'). \]  
\[ (16) \]

Substituting the value of \( t' \) and using the relation \(|E_\theta \times \hat{k}_s \times \hat{k}_s| = E_\theta \sin \phi\), where \( \phi \) is the angle between \( E_\theta \) and \( k_s \) (see Fig. 9), Eq. (16) becomes:

\[ E_s(R, t) = \frac{r_0^2 E_\theta \sin \phi}{R} \cos(k \cdot r - \omega_0 t'), \]  
\[ (17) \]

with \( k = k_s - k_\theta \) the differential scattering vector (see Fig. 9) and where \( r_0 \) refers to the classical electron radius:

\[ r_0 = \frac{e^2}{4 \pi \varepsilon_0 mc^2} = 2.82 \times 10^{-15} [\text{m}]. \]  
\[ (18) \]

The differential scattering cross section is obtained by dividing the power scattered by an electron in a solid angle \( d\Omega \) (in the direction of \( k_s \)) by the incident power per unit area:

\[ \frac{d\sigma_T}{d\Omega} = \frac{\frac{1}{2} e^2 \varepsilon_0 |E_\theta|^2 R^2}{\frac{1}{2} e^2 \varepsilon_0 |E_\theta|^2} = r_0^2 \sin^2 \phi, \]  
\[ (19) \]

Some interesting properties of the scattering process can be understood from Eqs. (16) – (19). The re-radiation has a maximum in a plane perpendicular to \( E_\theta \) (i.e., \( \phi = 90^0 \)), and as
a result, scattering of randomly polarized light leads to polarization of the scattered light. 

The small value of the scattering cross section of $8 \times 10^{-30} \text{ m}^{-2}$ results in a very low scattering yield. From Eq. (18) it follows that the power scattered by ions is much smaller than that of electrons by a factor of $(m/M_{\text{ion}})^2$ and is therefore negligible. Since the cosine term of the scattered field contains the term $k \cdot r$ only the components of $r$ and thus of $v_{\text{electron}}$ parallel to $k$ are observed in the spectrum of the scattered light.

Hitherto only the scattering of an electromagnetic wave by an individual electron has been considered. However, in a plasma many electrons contribute to the scattering process. If the contributions of the individual electrons add coherently at the point of observation, the scattered power is $N_e^2$ times larger than that from a single electron. Here, $N_e$ refers to the total number of electrons in the observation volume. In case that there are no correlations between the waves scattered by individual electrons (incoherent TS), the total scattered power will be proportional to $N_e$.

Without going in detail we present the result of electromagnetic theory, that describes the scattered power $P_s$ as follows:

$$P_s = P_0 \frac{d \sigma_T}{d \Omega} \sin^2 \phi n_e \Delta L \Omega S(k, \omega),$$  \hspace{1cm} (20)

with $P_0$ the incident power, $n_e$, the electron density, and $\Delta L$, the length of the scattering volume. The dynamic form factor $S(k, \omega)$ describes the frequency shifts resulting from electron motion and when $\alpha > 1$ also the effect of correlations between electrons. The contribution from electrons in each velocity interval must be integrated over the electron velocity distribution, $f(v)$, to determine the net contribution of the scattered light to the spectrum in each frequency interval:
in which the $\delta$-function takes care of the fact that each velocity leads to a Doppler-shifted, frequency

$$\omega_s(v) = \omega_0 + k \cdot v .$$

Therefore, the theoretical scattering spectrum will have the same profile as $f(v)$. When $f(v)$ along $k$ is a Maxwellian distribution

$$f(v_k) = \frac{1}{a \sqrt{\pi}} \exp\left[-\left(\frac{v_k}{a}\right)^2\right] ,$$

with thermal velocity: $a = \sqrt{\frac{2 k_B T_e}{m_e}}$, the dynamic form factor becomes:

$$S(\lambda_s) = \frac{1}{\Delta \lambda_e \sqrt{\pi}} \exp\left[-\left(\frac{\lambda_s - \lambda_0}{\Delta \lambda_e}\right)^2\right] ,$$

in which $k_B$ is the Boltzmann constant, $\lambda_0$ the wavelength of the incident radiation and $\lambda_s$ that of the scattered radiation. Eq. (24) predicts that the TS spectrum will have a Gaussian shape with a 1/e width:

$$\Delta \lambda_e = 2 \lambda_0 \frac{a}{c} \sin \frac{\theta}{2} = 2 \lambda_0 \frac{\theta}{c} \sin \frac{\theta}{2} \sqrt{\frac{2 k_B T_e}{m_e}} ,$$

which is

$$\Delta \lambda_e (nm) = 1.94 \sqrt{T_e (eV)}$$

for $\lambda_0 = 694.3$ nm (ruby laser) and $\theta = 90^\circ$. The width of the scattering spectrum is proportional to the square root of $T_e$, e.g. for $T_e = 1$ keV one finds $\Delta \lambda_e \sim 61$ nm.

Assuming a Maxwellian velocity distribution, TS is a powerful diagnostic to determine $T_e$ from the shape of the scattering spectrum, according to Eq. (25), and $n_e$
from the total scattered power, (in fact from the integral of the scattering spectrum) using Eq. (20).

Hitherto we have implicitly assumed that the thermal velocity is small compared to the speed of light. When this is not the case, relativistic effects have to be taken into account, leading to two effects. First, the classical Thomson cross section is slightly reduced due to the mass defect by a factor $1/\gamma^2$, e.g. for $v/c = 0.1$ ($T_e \sim 2.6$ keV) $1/\gamma^2 = 0.99$. Second, and more important, the “head-light” effect leads to radiation preferentially emitted in the forward direction, as seen by an observer at rest. The power scattered by an electron beam with $v/c = 0.1$ increases 36% and decreases 26% when moving towards or away from the observer. This “head-light” effect results in a so-called blue shift of the scattering spectrum. Analytical descriptions of the relativistic effects are given in literature.$^{5,67,68}$

In Fig. 10 some theoretical relativistic scattering spectra are shown for different values of the electron temperature. The calculation was performed for $\lambda_0 = 694.3$ nm and a scattering angle of 90°. It is clear from Fig. 10 that due to the relativistic effects the spectra tend to shift towards the blue. The shift is larger for increasing $T_e$. The total scattering yield increases only weakly with increasing $T_e$; e.g. for 10 keV the total scattered power is 12% larger than for a cold plasma.

III.B. General implementation of a Thomson scattering system

Thomson-scattering systems typically feature laser sources operating in the visible and near infrared spectral region and usually employ detection systems observing the light
scattered over ~90°. For a plasma with $T_e$~1 keV and $n_e$~$5 \times 10^{19}$ m$^{-3}$, the Debye length is $\lambda_D$~33 µm. Therefore, $\alpha$ is in the range $(2.5 - 5) \times 10^{-3}$, which means that the condition for incoherent adding of the powers scattered by the individual electrons is well met.

Due to the small value of $d\sigma_T/d\Omega \sim 8\times10^{-30}$ m$^2$, the scattering yield is very low. For typical experimental values $n_e$~$5\times10^{19}$ m$^{-3}$, $\Omega = 5\times10^{-3}$ sr, $\Delta L = 5\times10^{-3}$ m and a transmission of the optical system, $\tau \sim 0.2$ the ratio between scattered and incident power $P_s/P_0 \sim 2\times10^{-15}$. Consequently, light sources with high power are required to get sufficiently high scattering yield. Therefore, present TS systems use Q-switched lasers, which produce output powers of $> 100$ MW in ~15 ns.

In Fig. 11 the typical layout of a TS system is shown. In the remainder of this section the basic parts and features of such a TS diagnostic will be reviewed.

**Laser**

Most present TS experiments employ Q-switched ruby or Nd:YAG lasers as source. The ruby laser operating at 694.3 nm produces outputs up to 25 J in 15 ns ($P_0 = 1.67$ GW). However, their repetition rate is usually rather low: $\leq 5$ Hz (1 J/pulse). When more than several pulses per minute are required, an intra-cavity ruby laser can produce a burst of high energy pulses ($\sim 15$J/pulse, $\Delta \tau \approx 1$ µs) with a rep. rate of $\sim 10$ kHz. Ruby lasers are usually employed in systems where a good spatial resolution is preferred above a high time resolution. The most frequently used system for periodic TS measurements is based on the application of Nd:YAG lasers operating at 1064 nm, with outputs of $\sim 1$ J, 15 ns and a repetition rate of 20 - 50 Hz. Combining a set of lasers the repetition rate can be increased (see Sec. III.C.1).
The beam divergence of both type of lasers is about 0.3 - 1.0 mrad. The polarization of the laser beam is chosen perpendicular to the scattering plane. The high laser powers require special precautions for the used optics. Laser beam diameters should be kept large enough such that for 15 ns pulses the energy density is below the damage threshold of about 5 J/cm$^2$. Transmitting surfaces need to be coated and tilted with respect to the beam propagation to prevent back-reflected light entering the laser system again. Furthermore, curved transmission optics should have concave entrance surfaces, to prevent focusing of the back-reflected beams (which might lead locally to very high power densities).

Although other types of pulsed lasers (e.g. frequency-doubled iodine lasers, Ti:Sapphire and Alexandrite lasers) have been proposed, the authors are not aware of any real applications of these sources to Thomson scattering for high-temperature plasmas.

**Stray light reduction**

The laser beam enters and leaves the plasma vessel, through vacuum windows. Passing these windows - especially the entrance one - generates stray light, which can reach the collection optics. Without precautions this stray light level can be six to eight orders of magnitude larger than the TS light. Reduction of the vessel stray light can be achieved by tilting the windows (placing them under the Brewster angle is very effective), by positioning them relatively far from the plasma, by using baffles in both entrance and exit ducts, and by mounting a viewing dump on the vessel wall opposite to the collection optics (see Fig. 11). A very effective light trap (reduction up to 100 times) can be made from a stack of knife-edge blades. Carbon tiles on the inner wall of the plasma vessel can
give a reduction by a factor of \(~20\). Finally, the laser beam is dumped on a piece of absorbing glass placed under the Brewster angle.

**Collection and relay optics**

Scattered light is collected after passing a vacuum window and subsequently relayed to a spectrometer. Because of the low scattering yield the transmission of the collection and relay optics should be of course as high as possible. In devices with hot plasmas a shutter is required between the plasma and the window to reduce deposition of all kind of materials on the inner window surface, during the times the diagnostic is not in use.

Various kinds of optics are used to collect the scattered light: Cooke triplets, Bouwer katadioptric systems, achromatic doublets, etc. These systems are used along with other lenses or mirrors to guide the TS light to the spectrometer. Basically there are two possible ways to guide the scattered light from the plasma to the detection system: via flexible fibers and via conventional optics (lenses and mirrors).

The main advantage of fibers above conventional optics is that the linear etendue of the source can be matched to that of the detector, albeit at the cost of a reduced spectral resolution. For this purpose the fiber array is rearranged such that the slit height is reduced and the slit width increased, for example by a factor of 2. As a result, the usable solid angle of the collection lens increases by a factor of 4. However, for TS diagnostics at small-sized plasma devices (\(2a \leq 300 \text{ mm}\)) where the detection system can be positioned relatively close to the plasma (\(\leq 10 \text{ m}\)) conventional optics gives a better transmission (up to a factor of 3) than a fiber optics.\(^{70}\) For comparable systems, e.g. the multi-position TS systems of JFT-2M\(^{71}\) and RTP,\(^{72}\) the overall transmission is larger for
systems using conventional relay optics (RTP: 25%) than for those using fiber optics (JFT-2M: 7%). The major contributions to the losses in fiber optics systems are the core-cladding ratio, the packing fraction, the attenuation, input- and output reflection losses and an increase of the exit cone.

For fiber-optics arrays transmissions of about 55% and even higher have been reported.\textsuperscript{73,74} Despite the lower transmission, fiber-optics systems have to be preferred when the scattered light needs to be relayed over longer distances (e.g. to get outside the biological shield of the plasma device). To bridge long distances ($\geq$ 10m) with conventional optics would require many large sized lenses and mirrors, resulting in a low transmission as well.

\textit{Spectral analysis}

Mainly two different techniques to disperse the scattered light are in use for TS systems: filter and grating spectrometers. In filter spectrometers (Fig. 12) the scattered light is separated into different wavelength bands by means of a cascade of interference filters.\textsuperscript{75,76,77} The number of separate wavelength channels in these systems is usually rather limited (3 - 8 channels), and therefore the interpretation of the data relies on the assumption of a Maxwellian electron velocity distribution in the plasma. In grating spectrometers a grating is used for dispersing the scattered light (see Fig. 14). Both mechanically and holographically ruled gratings are used for this purpose. In this case, the number of independent spectral channels can be quite large: up to 80 for the TVTS system on TEXTOR.\textsuperscript{74} In case of good photon statistics, this enables to determine the shape of the scattering spectrum and possible deviations from a Maxwellian distribution.
To prevent that the residual of the vessel stray light disturbs the TS spectrum, the laser wavelength should be carefully filtered out after dispersion has taken place. This can be done by blocking the laser wavelength, by reflecting light at this wavelength away from the detector or by focusing it onto a special detector. Both filter and grating spectrometers have typical stray light rejection ratios of $10^{-4}$-$10^{-5}$ for the spectral channels adjacent to the laser line, decreasing to $\leq 10^{-5}$ at the edge of the spectral range.

Detection and data acquisition

Two different types of detection systems can be distinguished: time-resolving single- or multi-element detectors, and signal-integrating multi-element detectors. The first category includes: Avalanche Photo Diodes (APD) employing the high quantum efficiency of Si between 500 and 1000 nm, Photo-Multiplier Tubes (PMT), photo-diode arrays, and multi-anode PMTs. These systems enable time resolutions of the order of the laser pulse duration of 15 ns. As a result plasma light can easily be sampled just before or after the laser pulse. TS systems using periodic Nd:YAG lasers mostly apply APD for detection of scattered and plasma light. The signals of PD detectors are recorded with charge integrating Analogue-to-Digital Convertors (ADCs) or by means of fast transient recorders.\textsuperscript{78}

Time-integrating detectors have a lower time resolution because and are called TV systems because the detection principles are similar to those of a television camera able to receive a two-dimensional image. Vidicon, CCD (Charge Coupled Device), CMOS (Complementary Metal Oxide Semiconductor ) and also streak cameras belong to this category.\textsuperscript{79} These cameras have large numbers of image pixels, e.g. $10^6$ – $10^7$. The low read-out time can vary for different types of cameras. For a 16-bit CCD camera the
read-out time can be ~ 1 second, while ultra-fast CMOS cameras sample with frame rates of $10^4$ images/sec at a 12-bit dynamic range (see Sec. III C). The scattered light of the short laser pulse is captured with a gated image intensifier coupled to the TV-like recording system using a lens system. Data detected with TV systems is usually stored in internal memories and after termination of the plasma discharge send to a computer for analysis.

Both PD and TV detectors employ different kinds of photo-cathode materials to improve the photon-electron conversion process. PMT and image intensifiers are equipped with GaAsP, super S25, extended S20 cathodes to reach high conversion efficiencies in the visible and near infrared. The signal-noise-ratio (S/N) of a detector directly depends on this conversion efficiency:

$$ S / N = \sqrt{N_{pe}} = \sqrt{N_{ph} \eta_{\text{conversion}}} , $$

(26)

where $N_{pe}$ denotes the number of photo electrons generated by the incoming photons ($N_{ph}$) and $\eta_{\text{conversion}}$ the efficiency of the conversion from photons to electrons. However, the S/N of the complete detector will be lower because of the noise added by the amplification and read-out processes. More useful for evaluation of a detector is the effective detector efficiency, which includes the noise factor: $\eta_{\text{det}} = \eta_{\text{conversion}} / \text{noise factor}$. The noise factor refers to noise increase in the detector caused by the amplification process.

The S/N of a complete detection system is determined by the statistical noise, the dark current of the detector and background signals, as plasma light and stray light.
Plasma light due to Bremsstrahlung and line radiation easily can be corrected for when PD detectors are used. Using TV systems in combination with fiber optics for light relay offers the ability to sample plasma light just next to the laser beam and guide this to the same detector for simultaneous recording. The contribution of plasma light can be kept negligibly low when the laser energy is high (> 10 J) and the sampling interval is almost as short as the laser pulse (40 ns, for a pulse with 15 ns FWHM). Grating spectrometers combined with TV systems result in a large number of spectral channels, which enables line radiation to be suppressed. Whether correction of the vessel stray light is necessary depends on how effectively the stray light can be suppressed.

**Calibration**

Calibration is an important issue for TS systems. The width and shape of the spectral channels of a filter spectrometer can be calibrated by means of a tunable light source (e.g. a white light source in combination with a monochromator). The wavelength calibration of a grating spectrometer is done with the use of spectral lamps (He, Ne, Ar). The relative sensitivity of the spectral channels is calibrated using a tungsten filament lamp with known emitted power as function of the wavelength, in combination with an integrating sphere for uniform spectral radiance. The absolute sensitivity of the complete system can be calibrated by filling the vacuum vessel with nitrogen or hydrogen at relatively high densities and performing Rayleigh or Raman scattering. The cross sections for these scattering processes are well known. Since the calibration geometry is identical to that of the actual experiment it is a simple matter to relate the calibration measurements to the scattered signal from the hot electrons.
III.C. Some Elementary Application Examples

During the first 10 – 20 years of applications at magnetic confinement devices, most TS systems were basically diagnostics yielding the values of $T_e$ and $n_e$ at one local position in the plasma during a single instant in time. Many of the present-day TS systems aim to determine $T_e$ and $n_e$ along the full plasma diameter and at several time instances during the plasma discharge. Basically, three main TS systems are in use nowadays: (1) periodic photo-diode (PD) systems with a relatively high temporal resolution based on repetitive lasers (usually Nd:YAG) combined with a separate filter spectrometer (see Sec. III.B) for each spatial point. Often APD’s are applied for detection; (2) LIDAR TS systems in operation at JET and proposed for ITER, using the time-of-flight of a short laser pulse to obtain spatial resolution along the laser chord, and (3) single, dual or multi-pulse TVTS systems with a high spatial resolution featuring high power ruby lasers combined with intensified CCD or CMOS cameras. Below we will briefly describe typical examples of each of these three systems: the DIII-D repetitive TS system, the JET LIDAR TS system and the TEXTOR high-resolution TVTS system.

III.C.1. Repetitive systems

One of the most advanced PD systems is in use at the DIII-D tokamak. Quite similar systems are operated at ASDEX-UG, TCV and Alcator C-Mod. The DIII-D system features eight Nd:YAG lasers operating at 20 Hz. The laser beams are packed closely together in a $2 \times 4$ matrix. The maximum sampling rate that can be obtained is 160 Hz using 1 J, 15 ns pulses. The lasers can be triggered independently to generate
arbitrary pulse trains at 20 Hz repetition rate, with a minimum time separation of 100 µs. Scattered light from 40 spatial elements is collected by a set of viewing lenses and guided by fiber optics to filter spectrometers (Fig. 12), each with eight wavelength channels. The stray light rejection ratio amounts to $10^{-4}$-$10^{-5}$. Avalanche Photo Diodes (APD) employing the high quantum efficiency of Si (50 - 85%) between 600 and 1060 nm are used for detection of the dispersed light. The noise factor is typically 2 – 3. The signals of APD detectors are recorded by charge integrating Analogue-to-Digital Converters and subsequently stored in computer memories for further analysis.

**III.C.2. LIDAR systems**

A special case of a PD system is the so-called LIDAR. LIDAR is the acronym for LIght Detection And Ranging and is based on the time-of-flight measurement of a short light pulse. In a LIDAR system the light scattered over 180° (back-scattered) is collected. The spatial resolution is directly related to the length of the laser pulse and to the speed of the detection system. The laser pulse should be as short as possible to have a reasonable spatial resolution, and the power should be as high as possible to have high enough a scattering yield. Because of the very high laser power one should be careful to stay below the damage thresholds of the various optical components. The way to do this is to expand the beam to a large diameter in the optical transmission system. In the JET LIDAR system (see Fig. 13), which is the only operating LIDAR system world-wide, light from a ruby laser with a pulse length of 0.3 ns and with a corresponding physical length of 90 mm enters the plasma through a vacuum window. Back-scattered light is collected, led to a filter spectrometer with six channels, and subsequently detected with fast Micro-Channel Plate detectors. By analyzing these signals as a function of time both $T_e$ and $n_e$
can be determined as a function of the spatial coordinate along the line of sight, since light from different positions arrives at different times at the detector. Values of $T_e$ can be determined with an accuracy of 8% at $n_e = 5 \times 10^{19} \text{ m}^{-3}$.

III.C.3. TV systems

In the TEXTOR TVTS system the beam of a double pulse Q-switched ruby laser (694.3 nm; $2 \times 10^8$ J, 30 ns FWHM) is guided along a vertical path through the plasma. Stray light originating from the in- and output windows is brought to a minimum by means of baffles, a viewing dump of carbon tiles, and a notch filter in the spectrometer, thus suppressing the vessel stray light level by eight orders of magnitude. A shutter placed between the plasma and the detection window protects the window against deposition of all kind of materials originating from the plasma edge. Scattered light from a laser chord of 900 mm is collected by means of a multi-element lens and guided to the spectrometer outside the biological shield via a coherent fiber bundle.

The grating spectrum is of the Littrow type (Fig. 14), and enables to use a relative large etendue without vignetting. The output of the fiber bundle is projected onto the entrance slit by lens (2). Pupil imaging is performed by a doublet field lens (3) and a spherical mirror (7). After dispersion the spectrum is imaged at the two-part mirror (7) and finally by a camera objective (8) onto the cathode of an image intensifier (9). Subsequently, the intensified light of the image is recorded by means of two intensified CCD cameras (13). The intensifiers of the CCD cameras are used as switches to separately capture the first and the second laser pulse.
The heart of the detection system is formed by the high quantum efficiency (33%) GaAsP cathode of the image intensifier. Including the noise of the multiplication process the tube efficiency is ~ 18%. Due to the large number of detector elements the TVTS system of TEXTOR enables a determination of $T_e$ and $n_e$ profiles along the full plasma diameter with 120 spatial points of 7.5 mm length. At an electron density $n_e = 2 \times 10^{19}$ m$^{-3}$ and 10 J laser energy variations in the density can be determined with an accuracy equal to the statistical error of ~ 3 %, while the observational errors on $T_e$ are ~ 6 % in the range of 25 eV - 3 keV. In Fig. 15 an example of a double pulse measurement is given, showing evidence for MHD islands.$^{85}$

Recent work in the field of TV Thomson scattering systems is focused on overcoming the drawback that one can only measure one or two snapshots, albeit with high spatial resolution. The basic idea explored is that the plasma itself is contained in the laser cavity.$^{69}$ This allows generating 5 – 10 ms long bursts of 15 J laser pulses, with a repetition rate of 10 kHz inside the bursts. For TEXTOR it is expected that three of such bursts can be generated during one discharge. For the detection of the scattering spectra at these high sampling rates novel > 10 kHz CMOS cameras are employed. Two of these cameras are mounted in a similar set-up as sketched in Fig. 6. One camera monitors the Thomson scattering light, while the other one monitors the plasma light. The latter measurement is very important for this high-rep rate system because the individual laser pulses inside the bursts are typically 1 µs long, such that also the plasma light is integrated over this much longer time (compared to 40 ns for the double-pulse system).
III.D. Ion Collective Thomson Scattering (in the THz range and at CO$_2$ wavelength)

As was mentioned in Sec. III.A. the dynamic form factor $S(k, \omega)$ defined in Eq. (21) describes the frequency shifts resulting from the electron motion as well as from the effect of collective behaviour of the electrons. Because the large ion to electron mass ratio the Salpeter approximation may be used, which decomposes $S(k, \omega)$ into separate electron and ion contributions:

$$S(k, \omega) = C_e \Gamma_e(v_{Te}) + \left( \frac{\alpha^2}{1 + \alpha^2} \right) C_i \Gamma_i(v_{Ti}),$$

where the subscripts e and i refer to the electrons and ions, respectively, the C’s are multiplication constants and the $\Gamma$’s are the scattering form factors. In a simple two-particle plasma consisting of electrons and ions, the spectrum contains a broad feature with a width roughly given by $kv_{Te}$ and a narrow superimposed peak with a width $kv_{Ti}$. Schematically, one obtains a spectrum as depicted in Fig. 16. It is obvious from Eq. (27) that when $\alpha \ll 1$, the ion term can be completely neglected, thus yielding the formulae for incoherent Thomson scattering that have been already discussed in Sec. II.

In principle the local ion temperature in the scattering volume can be determined from the width of the narrow peak. This, however, cannot be done in the same straightforward way as the electron temperature can be deduced from the incoherent TS spectrum. The presence of impurities and geometric effects can severely complicate the analysis. Moreover, the interpretation of the spectra depends very sensitively on the electron temperature, which must be accurately measured by means of another diagnostic.
First ion TS measurements were reported in the mid sixties from relatively dense ($n_e \sim 10^{21} \text{ m}^{-3}$) and cool ($T_e \leq 5 \text{ eV}$) plasmas using ruby lasers at scattering angles of $\sim 1^\circ$ or larger. These included experiments on hydrogen arc plasmas, $^{86}$ thetatron, $^{87}$ and theta pinches. $^{88,89}$ The first attempt to measure the ion temperature of a tokamak plasma using the collective effects of the TS process was made at Alcator-C in 1983 by Woskoboinikov. $^9$ Successful ion TS measurements were performed by the TCA group, using a high-power D$_2$O-laser at 385 $\mu$m and a heterodyne receiver system. $^{10}$ Typical accuracies that could be obtained in the ion temperature determination were about $\pm 25\%$ at a density of $5 \times 10^{19} \text{ m}^{-3}$.

Presently, several ion TS diagnostics are in use or in development. $^{13,90}$ The emphasis of these diagnostics is put on the measurement of the velocity distribution of confined $\alpha$-particles and fast ions. The ion CTS systems can be divided into two categories. Firstly, the systems featuring high-power CO$_2$ lasers (10 $\mu$m), which necessarily use small scattering angles to achieve large enough values for the Salpeter parameter $\alpha$. The main advantages of these systems are the well developed laser technology and the wavelength that is far from background plasma emission and refractive effects. This approach was initially proposed by Hutchinson et al. $^{91}$ and a proof-of-principle result was obtained on the ATF stellarator. $^{92}$ The system is presently being commissioned on the JT-60 tokamak. $^{90}$ To meet the condition for collective scattering, a scattering angle of about $0.5^\circ$ is required. So far the results have not been conclusive due to electrical noise from the laser, the large stray light levels and due to higher order laser modes outside the notch filter. Secondly, the systems featuring high-power gyrotrons that operate in the millimeter range as source. The fast-ion collective TS
system at TEXTOR is the only CTS system that hitherto has reached the status of a routine diagnostic. Since this system works in the microwave spectral region, it is described – along with other microwave systems, elsewhere in this special issue.\textsuperscript{13}

### III.E. Application to Burning Plasma Experiments

Incoherent TS is ideally suited for application in a burning plasma experiment.\textsuperscript{60} On ITER a LIDAR Thomson Scattering system will provide the basic equatorial core measurements.\textsuperscript{93} A major benefit of a LIDAR system over a conventional system is the single port access and the automatic alignment between the laser beam and the detection system. Inside the port plug the same optical elements are used for both the laser beam and for the collected scattered light. As a result no remote alignment of these elements is required. The first optical mirror is deeply recessed into the port plug, thus minimising the degradation caused by plasma deposition and erosion. Nevertheless, as with any first mirror in ITER, careful consideration should be given to the selection of the mirror material, the detailed design of the mirror (cooling or heating) and the mirror geometry.\textsuperscript{94}

Using two different lasers, high spatial/low temporal and high temporal/low spatial resolutions measurements may be combined. One laser ($\tau = 3$ ps) could provide about 7 cm resolution ($a/30$) at moderate time rate (10 Hz), while the second laser ($\tau = 1$ ns) would allow for a faster measurement rate (100 Hz) with a more coarse spatial resolution (20 cm). Ruby, Nd:YAG and Ti-sapphire lasers are considered as candidate laser sources. To achieve a high spatial resolution, detectors are required with response time better than 300 ps and with a large spectral bandwidth ($T_e$ on a BPX could reach
values above 30-40 keV, leading to a very broad scattering spectrum). Such detectors, based on GaAs photocathodes, have recently been developed for JET.\textsuperscript{95,96} In the case of ITER the very wide TS spectrum (at $T_e(0)$ up to 40 keV) will reduce the accuracy of $T_e$ and $n_e$ at the highest temperatures but not affect the spatial resolution.

Significant effort has been focused on developing TS for the difficult to access divertor area.\textsuperscript{97} Resolutions of 5 cm along the leg and 3 mm across the divertor leg seem to be achievable on ITER. The remaining limitations are the limited spatial coverage, the low time resolution (20 Hz) and the lower $T_e$ limit.

LIDAR systems, operating with a laser line tilted to the flux surfaces, are also considered to provide measurements adjacent to the X point with a 0.5 cm effective radial resolution at 20 Hz. A similar technique (LIDAR or conventional TS measurements with the laser line tilted to the flux surfaces for an enhanced effective radial resolution) will be also used for edge $T_e$ and $n_e$ measurements in the main plasma. The technique has been extensively tested in JET where edge profiles are measured with 1 cm resolution at a low frequency ($\approx 1$ Hz) and reduced radial extension ($\approx a/10$).\textsuperscript{98}

Since the optimum wavelength for a fast ion collective Thomson scattering system on ITER is the microwave region, we refer to the previous paper in this issue for more information on this specific implementation.\textsuperscript{13}
IV. LASER DIAGNOSTICS FOR PLASMA DENSITY FLUCTUATIONS

IV.A. Propagation in inhomogeneous media

For probe frequencies $\omega_0 >> \omega_{pe}, \omega_{ce}$ and fluctuations with scale lengths far exceeding the Debye length $\lambda_D$, the plasma may be considered to be a continuous, albeit inhomogeneous, refractive medium with refractive index $N=(1-\omega_{pe}^2/\omega^2)^{1/2} \approx 1 - \omega_{pe}^2/2\omega^2$.

The propagation of a plane wave probe beam of wave vector $k_0$ is governed by an equation of the form

$$\left( \nabla^2 + k_0^2 \right) E = -k_0^2 \left( N^2 - 1 \right) E .$$  \hspace{1cm} (28)

Formally its solution may be expressed in the form of an integral equation, for which the well known Born approximation assumes that the incident field is only weakly affected by the refractive inhomogeneities. For small scattering angles and large distances from the scattering medium, the scattered or diffracted field, $E_d = E - E_0$, where $E_0$ is an incident plane wave, can be written in the Born approximation as

$$E_d (r,t) \equiv \frac{k_0^2}{4\pi |r|} \int_{V'} \left[ N^2 (r',t) - 1 \right] E_0 (r',t) \exp(ikr') dV',$$  \hspace{1cm} (29)

where $k = k_d - k_0$, $k_0$ is the incident wave vector and $k_d = k_0 |r|$ is the wave vector of the light scattered to the location $r$ satisfying $|r| >> |r'|$. Eq. (29) shows that in the far field the diffracted wave into the direction $k_d$ is proportional to the Fourier transform of $N^2 - 1$ for the spatial frequency $k$. Eq. (29) underlies the interpretation of wave scattering measurements.
For strong interactions there is in general no simple approximation for Eq. (28). One particular one, corresponding to forward scattering, is however noteworthy, since it corresponds to the approximation of geometrical optics. It may be obtained by ignoring the vector nature of the field amplitude, using the Ansatz $E(r) = \exp(\rho + i\phi(r))$, where $\rho$ and $\phi$ are real functions (Rytov transformation). The real and imaginary parts of the wave equation can be written as

$$\left(\nabla \phi\right)^2 = k_0^2 N^2 + \nabla^2 \rho + (\nabla \rho)^2, \quad \text{and}$$

$$\nabla^2 \phi = -2\nabla \phi \nabla \rho. \quad (31)$$

For media which are homogenous at a scale greater than the probe wavelength, the second and third terms at the right-hand side of Eq. (30) can be neglected, yielding the eikonal equation of geometrical optics. This equation defines rays which at every point in space are parallel to $\nabla \phi$ and perpendicular to the wavefronts defined by $\phi = \text{constant}$. The solution of the eikonal equation $^{100}$ is obtained as

$$\phi(r) - \phi(r_0) = k_0 \int_{r_0}^{r} N(l) dl, \quad (32)$$

where the line integral is evaluated along a ray.

The field amplitude is obtained by using Eq. (31):

$$\rho(r) - \rho(r_0) = -\frac{1}{2k_0} \int_{r_0}^{r} \frac{\nabla^2 \phi}{N} dl, \quad (33)$$

and back-transforming:

$$|E(r)| = \exp \left[ -\frac{1}{4k_0^2} \int_{r_0}^{r} \frac{\nabla^2 \phi}{N} dl \right] |E(r_0)|. \quad (34)$$
In most practical applications it is desirable that the amplitude be approximately constant and rays be nearly straight. This can be seen to be the case when \( L \equiv |\mathbf{r} - \mathbf{r}_0| << 1/|\nabla N/N| \).

A comparison of Eq. (34) with Eq. (29) shows that diffraction phenomena are beyond the scope of geometrical optics. Paradoxically however, many experimental situations can be described using combinations of both the diffraction theory and geometrical optics. For this reason it is needed to investigate the conditions for which geometrical optics may be applied in media where diffraction is known to occur. In general one is interested in small refractive index changes expressed as \( N = 1 + N^* \), where \( |N^*| \) can be arbitrarily small. The corresponding amplitude changes, estimated from Eqs (30) and (31), are of order \( \rho \sim k^2 L^2 |N^*|/4 \), where \( k \) is the wavenumber of the refractive perturbation of interest. Considering the terms neglected in the derivation of Eq. (32), the most stringent condition is obtained by requiring \( |\nabla^2 \rho| = k^2 \rho << k_0^2 \Delta N^2 \approx 2 k_0^2 |N^*| \), or, introducing the wavelengths \( \lambda \) and \( \Lambda \) of the incident radiation and of the refractive perturbation \( L << \Lambda^2/\lambda \) (near field condition).

The condition is pictured on Fig. 17, showing a plane wave incident onto a refractive slab of thickness \( L \), with transverse perturbations of scale \( \Lambda \). The above inequality is equivalent to stating that the diffracted waves, scattered at angles \( \pm \lambda/\Lambda \), must have diverged by much less than a distance equal to \( \Lambda \) from the incident wave on exiting the slab.

Fig. 18 illustrates the wavenumber matching condition \( k_{\pm 1} = k_0 \pm k \). Since only low frequency perturbations, \( \omega_{\pm 1} = \omega_0 \pm \omega \equiv \omega_0 \), are considered here, it needs to be required that \( |k_{\pm 1}| \equiv |k_0| \), leading to the Bragg relation, \( \sin(\theta/2) = k/2k_0 \). Permissible scattered wave vectors lie on the circle in the figure. It may now be noted that for an interaction volume
of finite depth $L$, $k$ is only defined within $\Delta k_z \approx 1/L$. From this one may distinguish two limiting regimes depending on whether or not the Bragg relation can be simultaneously satisfied for both orders $+1$ and $-1$. As seen from Fig. 18 the regime of simultaneous diffraction in both orders corresponds to $\Delta k_z >> k^2/k_0$ which is equivalent to the above near field condition. This regime is also referred to as Raman-Nath diffraction\(^{100}\) and refractive media satisfying the near field condition are called ‘thin’ phase objects. Hence over short enough distances geometrical optics may be used to describe wave propagation even if, over larger distances, a wave optics treatment is required.

Measuring the diffracted waves emerging from the plasma allows for some of the information on the refractive perturbations to be retrieved. It is up to the experimenter to devise a diagnostic that provides the desired information, depending on the nature of the fluctuations to be studied, within the often stringent access constraints of fusion research devices. Depending on circumstances, it may be advantageous to measure the diffracted light components in the far field (scattering methods), yielding information on the fluctuation wavenumber spectrum, or in the near field, as is the case for imaging diagnostics, which provide spatially localised information. If the amplitude and phase of all of the diffracted light could be measured, both approaches would be equivalent.

IV.B. Collective scattering diagnostics

Scattering diagnostics have been used extensively to study drift wave turbulence in high temperature plasmas,\(^{101,102,103,104,105,106,107,108,109,110,111,112}\) and to a lesser extent externally excited plasma waves such as Ion Bernstein Waves,\(^{114}\) Lower Hybrid Waves\(^{115,116}\) and Alfvén Waves.\(^{117}\) A comprehensive review of small angle scattering has
been given by Basse et al.\textsuperscript{118} (See also ref. 13, this issue). A typical arrangement is shown in Fig. 19. It involves a probing beam, typically a coherent Gaussian beam from a laser or a microwave source. Since scattered field amplitudes are extremely small, heterodyne or homodyne detection is required and the scattered radiation is made to interfere with a local oscillator provided by a second source, often at a slightly different frequency (heterodyne detection), or by splitting off a fraction of the source radiation (homodyne detection). The intersection of the two beams defines the scattering volume. The local oscillator can cross the probe beam or it can be introduced using a beam splitter after the scattered radiation has left the plasma (dashed lines in Fig. 19). The first case has the advantage of a reduced sensitivity to mechanical vibrations, but the wavenumber sensitivity drops for \( k < 1/w \), where \( w \) is the beam width and is strictly zero for \( k = 0 \). (See ref.\textsuperscript{119}, appendix B). The resulting power density on the detector can be written (up to a factor \( \varepsilon_0 \)) as

\[
I = \left| E_d + E_{lo} \right|^2 = \left| E_d \right|^2 + \left| E_{lo} \right|^2 + E_d E_{lo}^* + E_d^* E_{lo},
\]

(35)

where \( E_{lo} \) designates the local oscillator field and * denotes a complex conjugate. The detector size is chosen as to intercept all of the local oscillator and scattered radiation, corresponding to a power:

\[
P = \int (E_d E_{lo}^* + E_d^* E_{lo}) dS.
\]

(36)

Equation 36 is a scalar product and means that the detected power will be proportional to the projection (in the functional sense) of the scattered field onto the local oscillator field. This scalar product is conserved during propagation and may be evaluated anywhere along the local oscillator beam. This explains the high selectivity of heterodyne methods. Radiation not originating from within the portion of space defined by the local oscillator
beam, or not propagating in the direction \( k_{0i} \) defined by the local oscillator, does not contribute to the signal. For probe and l.o. fields with Gaussian profiles, the wavenumber resolution is given by \( \Delta k_y = \Delta k_x = 2/w \) and \( \Delta k_z = k/(k_0w) \).\(^\text{120}\) The Gaussian beam half-width \( w \) of the incident and l.o. beams is defined at the 1/e points of the beam intensity and similarly \( \Delta k_{x,y} \) are defined at the 1/e points of the power spectrum of the l.o. beam. The length of the interaction volume along the direction \( z \) is given by \( l = 2w/k_0 \). For a refractive perturbation of the form of a plane wave, \( \tilde{N}^* = \tilde{N}^*_0 \cos(kr - \omega t) \), satisfying the Bragg relation for a single order of diffraction, the scattered power from within the observation volume can be expressed as \( P_s = \tilde{N}^*_0 k_0^2 l^2 P_i / 4 \), where \( P_i \) is the probe beam power.

In traditional scattering arrangements the probe beam is nearly perpendicular to the magnetic field lines. The choice of probing wavelength then is a compromise between the conflicting desires for spatial resolution and for wavenumber resolution, as well as the necessity for measurements over a sufficiently wide range of \( k \). Adopting short wavelengths leads to small scattering angles, allowing for a wide range of wavenumbers to be collected using the same collection optics at the expense of spatial resolution along the probe beam direction. For many plasma waves of interest, scattering angles in the mid and far-infrared are too small to achieve any spatial resolution along the direction of propagation. Far infrared lasers\(^\text{106,107,108,110,111,117}\), and CO\(_2\) lasers emitting in the mid-infrared\(^\text{104,109,112,113,115,116}\) are the most popular sources. Consequently detection schemes may involve sub-millimeter diode mixers or liquid nitrogen cooled infrared photoamperic diodes or photoconductive detectors.
The first laser measurements of drift wave fluctuations were obtained in the small ATC tokamak by CO₂ laser scattering (λ=10.6µm)¹⁰² using a single photoconductive detector, scanning the wavenumber spectrum and beam position on a shot-to-shot basis. Further detail was obtained by a radially and vertically scanned multichannel scattering diagnostic on the TEXT tokamak.¹⁰⁶,¹⁰⁷,¹¹⁰ It used two coherent far infrared lasers with λ=1.222 mm, which were optically pumped by the same CO₂ laser. The probe and local oscillator lasers were actively feedback-tuned in order to achieve heterodyne detection with a stable intermediate frequency of 1 MHz, permitting a Doppler measurement of the direction of propagation of low frequency drift waves. The group also developed a scattering system allowing simultaneous measurements, at the same location, of the wavenumber spectra in perpendicular directions for addressing the issue of the poloidal isotropy of the turbulent spectrum (Fig. 20).¹¹¹

The main drawback of using a CO₂ laser is the lack of spatial resolution along the beam path. Under some assumptions, information on the spatial localisation can be retrieved using dual beam cross correlation.¹⁰³,¹¹² Spatial localisation is also possible by making use of the well founded assumption that $k_\parallel<<k_\perp$, i.e. turbulent structures are highly elongated along the magnetic field direction.¹⁰⁹,¹¹³ This implies that the scattered wave vector lies in a plane perpendicular to the local magnetic field lines and the spatial origin can be determined by relating the direction of the scattering angle to the local magnetic pitch, $B_0/B_T$ (Fig. 21). For a beam passing through the magnetic axis, the spatial resolution along the beam direction so obtained depends on the wavenumber resolution,
\[
\frac{\Delta z}{a} = \frac{\Delta k}{k} \frac{qR}{a} \left(1 - \frac{rdq}{qdr}\right)^{-1} = \frac{\Delta R}{w} \frac{q}{a} \pi \left(1 - \frac{rdq}{qdr}\right)^{-1}
\]
to first order. For the system on Tore Supra \(\Delta z \sim a/10\) for \(k=26\text{ cm}^{-1}\) in the core and \(\Delta z \sim a/3\) for \(k=6\text{ cm}^{-1}\) near the edge. This approach benefits from large beam widths \(w\) and should provide even better localisation in stellarators, reversed field configurations and low aspect ratio devices. Fig. 22 shows an example of a spectrum obtained on the Tore Supra device using this method.\(^{113}\)

As recently proposed\(^ {121}\) for a tokamak geometry, a substantial enhancement of spatial resolution can be obtained in a scattering arrangement where the probe beam is nearly parallel to the magnetic field lines. For small enough probe wavelength, as achieved using a CO\(_2\) laser, the scattered radiation originates from the immediate vicinity of the tangency region.

### IV.C Imaging and optical filtering methods for weakly refractive media

Imaging offers attractive alternatives to the more traditional scattering techniques in the case of perturbations with spatial scales such that the plasma can be considered to be a thin phase object. If a fluctuating wavefield has a characteristic spatial structure, a real space representation – an image – is likely to be more directly interpretable than its spectrum. Low amplitude fluctuations causing small phase shifts (\(\Delta \phi \ll 1\)) can for instance be measured using a homodyne Mach-Zehnder interferometer (see Fig. 2), tuned to fractional fringe observation, i.e. the probe and reference beams are adjusted as to have a phase difference of \(+\pi/2\) or \(-\pi/2\). Stabilising such an interferometer to a small fraction of a radian in the noisy environment of a fusion device, can however be a challenge.\(^ {122}\) A heterodyne version, as operated on the DIII-D tokamak\(^ {123}\) [M A Van Zeeland et al 2005]
Plasma Phys. Control. Fusion 47 L31-L40 'Alfvén eigenmode observations on DIII-D via two-colour CO2 interferometry'], with an intermediate frequency higher than the highest frequency of interest, may be less demanding in mechanical stability. Optical filtering methods do not require an external reference beam and are more tolerant with respect to mechanical perturbations, yet they can still retrieve phase information appropriate for fluctuations with large enough wavenumbers. The most accomplished of these is the phase contrast method, for which Zernike was awarded the 1935 Nobel prize.\textsuperscript{124} One group\textsuperscript{125} used it to study turbulent perturbations in high density plasmas, using a pulsed ruby laser and photographic film. The phase contrast imaging (PCI) method has since been applied to study density fluctuations in tokamaks and stellarators using wide probe beams produced by CO\textsubscript{2} lasers and multi-element detectors. Topics investigated include plasma turbulence,\textsuperscript{126,127,128,129,130,131,132,133} MHD fluctuations such as ELMS,\textsuperscript{134} quasi-coherent edge modes,\textsuperscript{135} and Alfvén Cascade modes,\textsuperscript{136} externally driven radio-frequency waves such as kinetic Alfvén waves,\textsuperscript{137,138} ion Bernstein and ion cyclotron waves.\textsuperscript{139,140}

\textit{IV.C.1 Detection of small phase shifts.}

Let’s consider the case of a plane wave beam with an amplitude profile described by $B(x)$. Its interaction with a thin phase object may described by a multiplication with the eikonal phase factor Eq. (32)

$$B'(x,t) = \exp(i\phi(x,t)))B(x) \equiv (1 + i\phi(x,t))B(x) \text{ for } \phi<<1.$$  \hfill (37)

For simplicity the time dependence and the vector notation are dropped for the position $x$.

In order to obtain a detectable change in intensity (power density) a wavefield $iB(x)$ may be added that can constructively interfere with the diffracted field $i\phi(x)$:
PCI and similar methods use no external reference (Fig. 23). The diffracted wavefield is separated from the undiffracted in the focus of a lens (L1 in Fig. 23) and phase-shifted by \(-\pi/2\) (or \(+\pi/2\)) with respect to it. This leads to:

\[
B''(x) = B(x) + i\phi(x)B(x) \quad \text{and} \quad I(x) = |B''(x)|^2 = 2|B|^2(x)(1 + \phi(x)). \tag{38}
\]

Other techniques can be understood similarly, although the phase shifts obtained are not necessarily the ideal \(\pm \pi/2\) for all wavenumbers in the spectrum. Clearly no such ‘internal reference’ instrument can measure absolute phase changes like an interferometer.

\textbf{IV.C.2. The phase contrast technique}

The phase object in \(\Sigma\) is illuminated with a suitably expanded beam of parallel light from a source such as a laser (Fig. 23). The spatial Fourier spectrum of the transmitted wavefield is mapped onto the focal plane of the lens L1, \(k \rightarrow kAf_1/2\pi\).\textsuperscript{141} The difference in thickness of the phase plate P is such as to phase-shift diffracted light with \(|k| > k_c\) by \(-\pi/2\). Lens L2 produces a Fourier back-transform of this filtered spectrum (i.e. an image of \(\Sigma\)) in \(\Sigma\). In the following analysis of the transfer properties a magnification equal to unity shall be assumed.

The transmitted beam in plane \(\Sigma\) may be written as:

\[
B'(x) = \{1 + i\phi(x)\}B(x), \tag{40}
\]
where $B(x)$ is the incident beam. In the Fourier plane the wave amplitude is

$$
\tilde{B}'(k) = \left\{ \delta(k) + i\tilde{\phi}(k) \right\} \otimes \tilde{B}(k),
$$

(41)

where $\sim$ designates the Fourier transform (FT) and $\otimes$ denotes a convolution. The effect of the phase plate is described by (Fig. 24):

$$
\tilde{B}''(k) = \left[ 1 - (1-i)\tilde{C}(k) \right] \tilde{B}'(k),
$$

(42)

where $\tilde{C}(k) = 1$ for $|k| \leq k_c$ and $\tilde{C}(k) = 0$ for $|k| > k_c$. The wave amplitude distribution in the image plane is obtained by an FT back-transform:

$$
B''(x) = R(x) + D(x),
$$

(43)

where

$$
R(x) = B(x) + (i-1)B(x) \otimes C(x)
$$

$$
D(x) = i\tilde{\phi}(x)B(x) - (1+i)\left\{ \phi(x)B(x) \otimes C(x) \right\}
$$

and $C(x)$ is the inverse FT of $\tilde{C}(k)$. The interference terms of $R$ and $D$ are linear in $\phi(x)$:

$$
\Delta I(x) = R(x) \ast D(x) + c.c. = 2B(x) \left[ B(x) \otimes C(x) \right] \left\{ \phi(x) - C(x) \otimes \left[ B(x)\phi(x) \right] \right\}.
$$

(44)

$\Delta I(x)$ is proportional to the difference between $\phi(x)$ and a weighted average of $\phi$ over the neighborhood of $x$. The instrumental impulse response is obtained by setting $\phi(x) = \delta(x-y)$:

$$
h_i(x, y) = 2B(x) \left[ B(x) \otimes C(x) \right] \left\{ \delta(x-y) - \frac{B(y)C(x-y)}{B(x) \otimes C(x)} \right\}.
$$

(45)

In the shift-invariant approximation (near $y=0$) this further simplifies to

$$
h_i(x, y) = h(u) \equiv 2I_0 \left\{ \delta(u) - \frac{B(-u)C(u)}{B(v)C(v)dv} \right\}.
$$

(46)

We have also used $B(x) \left[ B(x) \otimes C(x) \right] \equiv \left| B(x) \right|^2 = I_0(x)$. The FT of $h(u)$ provides the instrumental transfer function (wavenumber response):
\[ H_i(k) \equiv 2I_0 \left\{ 1 - \frac{\widetilde{B}(-k) \otimes \widetilde{C}(k)}{[B(-l) \otimes C(l)]_{l=0}} \right\}. \]  

(47)

The transfer function has a soft cutoff at \( k = k_c \) and becomes wavenumber independent for higher wavenumbers. If \( k_c \) is chosen as to match the diffraction-limited spot size corresponding to \( B(x) \) the instrument is sensitive to perturbations with wavelength shorter than the beam width in the object plane. Its transfer function then approaches that of a theoretical ideal internal reference interferometer, which depends only on the beam intensity profile. (See ref. 119, appendix A). For \( k > k_c \) the response is four times larger than for an interferometer making use of the same laser source (Fig. 25). In order to overcome detector saturation limits, sensitivity can be enhanced by reducing the undiffracted (local oscillator) power in the image plane using a phase plate with a partly transmitting central area of amplitude transmittance, \( \gamma \leq 1 \), to obtain \( \Delta I / I_0 \equiv 2\phi/\gamma \) for \( k > k_c \).

In a real system shift invariance is limited. For off-axis imaging with a cutoff near the diffraction limit, the impulse response becomes asymmetrical and consequently, for \( k < 2\pi/d \), where \( d \) is the beam diameter, the phase of \( \widetilde{H}(k) \) deviates from zero by up to \( \pm \pi/2 \) as \( k \to 0 \). A correction for \( k < 2\pi/d \) may be achieved by inverting the system transfer matrix, obtained by digitizing the position dependent impulse response \( h(x,y) \).

**IV.C.3. The ‘knife edge’ and phase blade methods**

It is instructive to compare phase contrast to its venerable ancestor, Foucault’s knife edge method, when used for weakly refractive media. In this arrangement all of the
diffracted light with $k > k_c$ is stopped, while all the light with $k < k_c$, including most of the undiffracted light, is allowed to pass and form an image in $\Sigma'$. The effect of the knife edge can be written as

$$\tilde{B}''(k) = B'(k)\left[1 - \tilde{E}(k)\right],$$

(48)

where $\tilde{E}(k) = 1$ for $k > k_c$ and $\tilde{E}(k) = 0$ for $k < k_c$.

Following the same formalism as for phase contrast (see ref. 147, page 4.10), one obtains

$$H_i(k) = i\lambda\left\{\frac{\tilde{B}(k) \otimes [\tilde{E}(k) - \tilde{E}(-k)]}{B(0)}\right\}.$$  (49)

An improved version of Foucault’s knife edge method consists in replacing the knife edge by a transparent blade producing a phase shift equal to $\pi$. This case is treated by setting $\tilde{E}(k) = 2$ (instead of 1) for $k > k_c$ in Eq. (21), resulting in a transfer function that is twice that of the knife edge (Fig. 25c). $H_i$ is imaginary and odd, meaning that all wavenumber components are $\pi/2$ phase shifted, which leads to a scrambling of complex images. A variation of this method (inaccurately called phase contrast) was tested on the CDX device.\textsuperscript{143}

**IV.C.4. Near field and defocusing effects**

A simple way of obtaining intensity variations as a result of phase variations is to position the detector a distance $z$ behind the phase object or by defocusing the object in an imaging arrangement. Propagation of the transmitted wavefield $B'$ at small angles over a distance $z$ is described by:

$$B''(x) = B'(x) \otimes h_z(x) = \left\{1 + i\phi(x)\right\}B(x) \otimes h_z(x)$$
where \( h_z(x) \equiv \frac{k_0}{2\pi iz} \exp\left\{ ik_0 z \left( 1 + \frac{x^2}{2z^2} \right) \right\} \) is the impulse response of free space propagation in the Fresnel approximation.\(^\text{141}\) The interference terms in the expression for the intensity are:

\[
\Delta I(x) = -i \{ B(x) \otimes h_z(x) \} \otimes (B^*(x)) \otimes h_z^*(x) \} + c.c.
\]

The near field may be defined as the region close to the phase object where \( B(x) \otimes h_z(x) \equiv \exp(ik_0 z)B(x) \). Following the same formalism as for phase contrast, one obtains the instrumental transfer function for propagation over a distance \( z \):\(^\text{119}\)

\[
H_i(k) \equiv -2I_0 \sin\left( \frac{zk^2}{2k_0} \right). \quad (50)
\]

This method is known as ‘amplitude scintillation’.\(^\text{144}\) The very oscillatory character of \( H_i \), as depicted on Fig. 25d, makes out-of-focus measurements difficult to interpret. Also, part of the information in the transmitted wavefield is not revealed in the intensity modulations. The effect of defocusing can be included in the formal treatment of phase contrast and interferometry. For phase contrast, the in-focus transfer function must be multiplied by the factor \( \cos\{ zk^2 / 2k_0 \} \), while for the homodyne interferometer, because of interference with both the reference and the transmitted beams, the correction factor is \( \sqrt{2\sin(\pi/4 + zk^2 / 2k_0)} \). For a heterodyne interferometer with beat frequency \( \omega_b \), the modulation phase \( \exp(-i\omega_bl) \) is shifted by \( \exp(izk^2 / 2k_0) \), scrambling complex images. Requiring these corrections to be small is equivalent to the near field condition.
IV.C.5. Line integration through a fluctuating medium

Imaging methods suffer the drawback that the information obtained is not local in the direction of propagation. Nonetheless, it is often possible to infer estimates of local fluctuation amplitudes from line-integrated measurements. Using a random walk argument, one can show that for homogeneous turbulence there is a relationship between the average local fluctuation variance and the variance of the line integrated fluctuations:

\[
\langle \overline{N}(x,t)^2 \rangle = l_z L \langle |N(x,z,t)|^2 \rangle , \tag{51}
\]

where \( \overline{N}(x,t) = \int N(x,z,t)dz = \phi / k_0 \),

\( L \) is the length of the probe beam path in the random medium and

\[
l_z = \frac{\int \langle N(x,z,t)N(x,z+\Delta z,t)dz \Delta z \rangle }{\langle |N(x,z,t)|^2 \rangle} \text{ is the coherence length along the direction of propagation. The length } (l_z L)^{1/2} \text{ is called the effective integration length.}
\]

More generally, the autocorrelation function of \( \overline{N} \) is a line-integral of the autocorrelation function of \( N \). This implies that the wavenumber and frequency spectra of \( \overline{N} \) are identical to those for \( N \), taken for \( k_z=0 \).\textsuperscript{126}

A second situation involves a quasi-monochromatic refractive wave in spherical or cylindrical symmetry\textsuperscript{146}. Although in this case an Abel inversion is in principle possible, the number of detectors in the image plane may not be sufficient. A useful approximation is given by\textsuperscript{147}

\[
\int G(r) \exp(i2\pi r / \Lambda - \omega t)dz = \sqrt{\Lambda x} G(x) \exp(i2\pi x / \Lambda + \pi / 4 - \omega t) , \tag{52}
\]

where \( r=(x^2+z^2)^{1/2} \) and \( G(r) \) a slowly varying envelope function.
IV.C.6. Sensitivity and calibration

Since all imaging methods can be compared on the basis of their transfer functions, we only need to treat the case of phase contrast, assuming photoamperic detectors which produce a photocurrent given by \( j = A e \Phi \eta \Phi \) where \( \Phi = I / \hbar \omega \) is the photon flux onto the detector of area \( A \) and quantum efficiency \( \eta \). In most circumstances, when the background intensity \( (I_0 / \gamma^2 \) for phase contrast) is high enough, the main limitation for photoamperic detectors is due to shot noise\(^\text{145}\)

\[
\langle j^2 \rangle = 2e\Delta f \times j = 2e^2 A \eta \Delta f / \Phi ,
\]

where \( \Delta f = \Delta \omega / 2\pi \) is the bandwidth under consideration.

The mean signal power for phase contrast is given by

\[
\langle j^2 \rangle = 4A^2 e^2 \eta^2 \gamma^2 \langle \Phi^2 \rangle / \Phi_0^2 ,
\]

with \( \Phi \equiv \Phi_0 = I_0 / \gamma^2 \). The signal-to-noise power ratio is then

\[
S / N = \frac{2A \eta l_0 \langle \phi^2 \rangle}{h \omega \Delta f} .
\]

At high intensities it may be advantageous to attenuate the direct light (e.g using a BaF\(_2\) substrate \( (\gamma=0.2) \) for producing a phase mirror\(^\text{146}\)) in order to avoid saturating the detectors. Although the wavelength dependence appears to favour the far-infrared, the mid-infrared offers far higher laser powers (tens of Watts), detectors with quantum efficiencies in the range 20-70\% and satisfy the near field condition for larger \( k \). If a CO\(_2\) laser \( (\lambda=10.6\mu m) \) and HgCdTe photoamperic diodes \( (A \eta l_0=1mW) \) are used, the sensitivity limit \( (S/N=1) \) corresponds to \( \langle \phi^2 \rangle^{1/2} \sim 3 \times 10^{-6} \) radians or \( \int n_{d\ell} - 10^{14} \) electrons/m\(^2\)
for $\Delta f = 1$MHz. For fluctuations with effective integration lengths of a few cm, this corresponds to $\Delta n_e/n_e \sim 10^{-4}$ for typical magnetic confinement devices.

An absolute calibration as a function of $k$ can be obtained using sound or ultrasound waves, together with a calibrated sonometer to measure the pressure and hence density oscillations, relating the local refractive index variation to the phase fluctuations, using the effective integration length $\sqrt{R\Lambda}$, (Eq. (52)) where $R$ is the distance to the sound source and $\Lambda$ its wavelength.\textsuperscript{146,148} Somewhat counter-intuitively, the impulse response can be measured by scanning a narrow opaque object through the beam to produce nearly the same response as a narrow phase object.\textsuperscript{146,147}

IV.C.7. Fluctuation measurements using laser imaging in magnetic fusion devices

Fig. 26 shows as an example, the phase contrast apparatus used on the TCA tokamak ($R=61$cm, $a=18$cm, $B_T=1.5$T).\textsuperscript{146} An 8-Watt waveguide CO$_2$ laser and beam expansion optics were mounted on the rear side of the optical ‘breadboard’ (not shown). The beam was expanded to $23 \times 5$cm and refocussed with parabolic mirrors used slightly off-axis. The beam was relayed to, and back from, the plasma by two sets of three flat mirrors, $45 \times 7$cm in size, which were mounted in rigid boxes at right angles to each other in a corner-cube arrangement providing immunity against rigid-body movements. Without such passive vibration proofing, active feedback stabilization, as used by Coda and Porkolab,\textsuperscript{148} may be necessary in order to keep the focal spot in the groove of the phase mirror.

Elongated NaCl windows ($23 \times 3.6$ cm clear aperture, 3cm thick, later replaced by ZnS) provided access to more than half of the plasma cross section and wavenumbers
down to $k_c=0.3\text{rad/cm}$. This long wavelength capability allowed the identification of a spectral maximum at $k\approx 1.2\text{ rad/cm}$ ($k\rho_s\approx 0.3$) for drift wave turbulence, well below the capabilities of scattering diagnostics using narrower beams. Unfortunately such generous access has not been possible on later implementations of the technique.

As an example, Fig. 27 shows amplitude and phase profiles of mode converted radio-frequency waves on the C-Mod and TCA tokamaks with a wide range of radial wavelengths. Both systems were equipped with multi-element photoconductive HgCdTe detectors. In the case of C-Mod the laser was acousto-optically modulated in order to downshift the plasma-wave induced intensity modulations at the RF frequency (50MHz) to a beat frequency near 1MHz, which was within the bandwidth capabilities of the 32 element detector and acquisition. Fig. 27a shows a localized, damped ion cyclotron wave generated at the D-$^3\text{He}$ hybrid layer. The convention in the figure is such that a phase with negative slope corresponds to propagation towards the low field side. The system on TCA relied on analog synchronous detection at the RF frequency (2 MHz) and mode conversion was observed at the shear Alfvén resonance layer (Fig. 27b and c).

As in the case of scattering diagnostics, spatial resolution can be improved by making use of the fact that turbulent structures are highly elongated along field lines ($k_||<<k_\perp$). For tangential imaging, as apparent from Eq. (52), the density fluctuations which contribute to the signal are strongly weighted towards the region where the beam is tangent to the flux surfaces (and preferably parallel to the magnetic field lines). This arrangement also provides a measure of isotropy ($k_\theta$ versus $k_\varphi$). Another method, first implemented on LHD, relies on a vertical beam passing near the magnetic axis and makes use of the variation of the magnetic pitch angle along the beam. LHD is ideally
suited for this technique because the pitch angle varies nearly linearly along the beam path by some ±40°. The image plane features an 8 by 6 element photoconductive array, which allows for 2D wavenumber spectra to be computed. By slicing these spectra according to pitch angle before Fourier backtransforming, the local density fluctuations along the incident beam can be reconstructed with a resolution equivalent to ~1/10 of the minor radius. A snapshot of turbulence in LHD, resolved into the lower and upper halves of the plasma, is shown in Fig. 28.132

IV.C.8. Optical filtering for strongly refractive media

The above optical filtering methods for weakly refractive media have striking instrumental similarities with optical methods used for short lived high density plasmas such as z-pinches,149,150 plasma focî,151,152 and laser plasmas,153,154 and generally use pulsed visible or VUV lasers and photographic film of CCD’s as recording material. Unlike phase contrast, these methods have been developed primarily for refractive media producing large phase shifts (ψ>>2π), even at visible wavelengths and are interpreted in the frame of geometrical optics.155,156,157 These methods usually emphasize regions characterised by refractive light deflection by gradients, as well as localised sources of scattered light. A classic example is the knife-edge Schlieren method invented by Foucault for the inspection of optical surfaces.142 In the geometrical optics description, refractive index gradients produce changes in the direction of light propagation (‘refraction’) which are given to first order by

\[
\alpha(x, y) = k_0 \nabla_\perp \phi(x, y),
\]

(57)
where $\phi$ is the eikonal wave phase and $\alpha$ is the angular deflection. The effect of the knife edge is to stop light with propagation angles below a threshold given by its position in the focus of L1, from illuminating the image. As a result only those regions of the phase object which produce phase gradients transverse to the knife edge exceeding a value depending on the position of the edge will be bright. Numerous variants exist depending on the nature of the filtering and illumination\textsuperscript{157}. If the illumination is produced by an extended source, or the knife edge is not exactly in the focus of L1, the intensity in the image can be made to be proportional to $\alpha$ over a certain range. An example from a Rayleigh-Taylor unstable plasma is seen in Fig. 29. Since the Schlieren methods provide simultaneous measurements of transverse wavenumber and position they are, just as scattering methods, limited by the uncertainty principle to $\Delta(\nabla_{x,y}\phi)\Delta x, y \geq 2\pi$.\textsuperscript{158} Out-of-focus imaging, in this context referred to as shadowgraphy (no filter), produces intensity modulations proportional to the second transverse derivative of $\phi$ (see ref. 152 for examples). From Eq. (34) to first order:

$$\frac{\Delta I(x, y)}{I_0} = -\frac{z\nabla^2_{x,y} \phi(x, y)}{Nk_0}.$$  \hfill (57)

Comparison with Eq. (50) shows that this is equivalent to out-of-focus imaging of weakly refractive objects only in the $k \to 0$ limit.

**IV.D. Summary and outlook**

Improvements in detector and laser technology have led to plasma physics applications and refinements of generic optical techniques such as scattering and imaging, which have
established the ubiquity of drift wave turbulence in magnetic confinement devices and allowed measurements of many other types of instabilities and waves. These measurements can be directly compared to theoretical predictions of fluctuating wavefields, thereby validating (or, as may be, invalidating) our understanding of these fundamental, albeit complex processes. A drawback compared to particle probe diagnostics, such as the Heavy Ion Beam Probe (see another paper in this issue\textsuperscript{159}) is however that they do not measure the fluctuating electrical fields and their phase relation to the density fluctuations and therefore cannot provide a direct measure of turbulent fluxes. It should be noted that the measurement of plasma waves with known dispersion properties can itself provide information on plasma parameters such as ion temperature (from ion Bernstein waves\textsuperscript{114}) or safety factor and mass density (from Alfvén waves\textsuperscript{137}), the latter being linked to the D/T density ratio in a burning plasma.

Wide probe beams and high quality optical components are required for optical filtering techniques such as PCI, if very long wavelength fluctuations are to be studied. Wide beams are also required for achieving high spatial resolution by making use of the pitch angle variation in scattering arrangements\textsuperscript{109} and PCI\textsuperscript{131}. A tangential scattering diagnostic can achieve good spatial resolution with very manageable beam widths, as proposed\textsuperscript{121} for the JET device and similarly good resolution may be expected from a tangential imaging diagnostic\textsuperscript{143}. The examples of DIII-D\textsuperscript{110,148}, Tore Supra\textsuperscript{109} and LHD\textsuperscript{131} show that laser scattering and imaging techniques can be successfully applied also to large devices. Scattering and imaging diagnostics can make use of the same laser source, detectors and beam delivery optics and therefore could be reconfigured into each other, depending on experimental requirements. They could be applied to burning plasma
experiments relatively easily for modest size beam cross sections, which could be relayed by a succession of mirrors, allowing the vacuum windows and detectors to be protected from neutrons. In addition to the wide range of applications mentioned in this article they could contribute to diagnosing turbulence in conditions of intense alpha particle heating, as well as fast particle instabilities.

V. CONCLUSION

Nearly 40 years after the first TS measurements in fusion devices, the field of laser plasma diagnostics has diversified into a broad range of applications and is still developing. Although many of the methods that were described in this paper are already applied for several decades, there is still a continuous stream of innovative ideas that extend the diagnostic possibilities. In this paper we have tried to explain the basic principles of the various laser-aided diagnostics and to report on the current developments in the field. Because of length restrictions to this paper it was unfortunately not possible to give an exhaustive overview of the complete field, neither was it possible to include references to all work that has been published. The keen reader therefore is encouraged to discover the rich literature this field is endowed with, which extends well beyond the reference list below.
ACKNOWLEDGMENTS

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REFERENCES


G. VAYAKIS, “Generic issues for a BPX,” This issue.


H. PARK et al., ”Tokamak ion temperature determination via CW Far-Infrared laser scattering from externally excited Ion Bernstein Waves” Nucl. Fusion 25, 1399 (1985).


J.A. Snipes et al., ”Active and fast particle driven Alfven eigenmodes in Alcator C-Mod”, Phys. Plasmas 12, 056102 (2005)


Table 1: Parameter/technique matrix for laser-aided plasma diagnostics

<table>
<thead>
<tr>
<th>Parameter/technique matrix</th>
<th>Electron density (profile)</th>
<th>Poloidal magnetic field/current density profile</th>
<th>Electron temperature (profile)</th>
<th>Electron density fluctuations</th>
<th>Ion temperature</th>
<th>Fast ion population</th>
<th>Described in Section</th>
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<td>Interferometry</td>
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<td>Polarimetry</td>
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<td></td>
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<td></td>
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<tr>
<td>Ellipsometry</td>
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<td></td>
<td></td>
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<tr>
<td>Incoherent Thomson scattering</td>
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<td>☀</td>
<td>★</td>
<td></td>
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<tr>
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<td></td>
<td></td>
<td></td>
<td>★</td>
<td>☀</td>
<td>★</td>
<td>III.D (fast ions)</td>
</tr>
<tr>
<td>Phase Contrast Imaging</td>
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<td></td>
<td></td>
<td>★</td>
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</table>

★ Main, ☀ Backup + Supplementary, ☀ Supplementary
Fig. 1. Standard geometry for interferometry and polarimetry measurements on a tokamak plasma.

Fig. 2. Schematic representation of a single-channel heterodyne Mach-Zehnder interferometer.
Fig. 3. Schematic drawing of the nine-channel interferometer / polarimeter on TEXTOR. The line-integrated electron density and the Faraday rotation angles are measured along nine vertical and one horizontal lines of sight.

Fig. 4. Typical geometry for measuring the electron density via polarimetry. Figure taken from ref. [50]
Fig. 5. Optical setup of one probing beam of the combined interferometer/polarimeter at TEXTOR. The polarization of the waves is indicated by bold double-headed arrows. Figure taken from ref. [37].
Fig. 6. Schematic of the triple-beam polarimeter setup after [54]. Two laser beams with angular frequencies $\omega_1$ and $\omega_2$ are transmitted through the plasma in orthogonal circular polarization states. Co-alignment of the two beams is achieved without power loss using a polarizer/quarter wave plate combination. A third laser at frequency $\omega_3$ acts as local oscillator for heterodyne detection of the interferometric phase shift.

Fig. 7. (a) Plasma current and average toroidal magnetic field. (b) Faraday rotation signals for standard Ohmic MST plasma. Sawtooth crashes are marked by spikes in the toroidal gap voltage, $V_{tg}$. The vacuum vessel center corresponds to $x = 0$. (Figure taken from [56])
Fig. 8. Sketch of the poloidal polarimetry system (at 118.8 µm) showing rays and locations of retro-reflectors. The vertical chords that are indicated are only proposed.
Fig. 9  Scattering geometry and definition of parameters.

Fig. 10. Theoretical relativistic scattering spectra for a plasma diagnosed with a ruby laser (694.3 nm) under a scattering angle of 90°.

Fig. 11. Typical layout of an incoherent Thomson-scattering system.
Fig. 12. Layout of a filter spectrometer as in use at DIII-D (taken from Ref. 75).

Fig. 13. Layout of the JET LIDAR system (taken from Ref. 65)
Fig. 14. Littrow grating spectrometer of the TEXTOR high resolution TS system.

1a fiber array (bulk TS)  1b fiber array (esge TS)
2 relay lens  3 field lens doublet  4 entrance slit  5 Littrow triplet  6 grating
7 two-part mirror  8 camera objective  9 image intensifier  10 coupling lens  11 beam splitter  12 fast CMOS camera

Fig. 15. First ten of 18 $T_e$ and $n_e$ profiles, recorded every 200 $\mu$s, smoothed over 4 spatial points. The corresponding laser energy ranged from 12 to 8 J. The profiles for $T_e$ and $n_e$ are plotted above each other with an equidistance of 0.3 keV and $1 \times 10^{19} \text{m}^{-3}$, respectively.
Fig. 16. Schematic illustration of the different contributions to the scattering form factor (not to scale).
Fig. 17. Diffraction by a slab with refractive perturbations: illustrating the near field condition.

Fig. 18. The wavenumber matching condition

Fig. 19. Typical scattering geometry for heterodyne detection. B incident beam, LO local oscillator, V scattering volume, S beam splitter, D detector.
Fig. 20. Bi-directional multi-channel scattering diagnostic on the TEXT-upgrade tokamak.
Fig. 21. Pitch angle localisation technique for a typical tokamak geometry. The y axis designates the direction of the probe beam, Ψ is the local angle of the scattered wave vector with the magnetic axis.\textsuperscript{113}

Fig. 22. Wavenumber spectrum of drift wave turbulence for 0.7 < r/a < 1 in Tore Supra obtained by CO2 laser scattering at different levels of ion cyclotron heating power.\textsuperscript{113}

Fig. 23. Optical setup for phase contrast. Σ object plane, Σ' image plane, L1, L2 lenses, P phase plate, D diffracted wave component, UD undiffracted wave component.
Fig. 24. Fourier plane quantities and transfer function for phase contrast.

Fig. 25. Impulse responses (top) and transfer functions (bottom) of methods for measuring small phase shifts.

a), a’) Interferometry
b), b’) Phase Contrast
c), c’) Knife edge (solid line) and phase blade (broken line)
d), d’) Near field and defocused imaging (‘amplitude scintillation’)
Fig. 26. Optical arrangement for phase contrast on the TCA tokamak. PL optical table, $\Sigma$ object plane at the plasma midplane, F vacuum window, CV vacuum chamber, BT toroidal field coil, GA major axis, P parabolic mirror ($f=190.5\,\text{cm}$), M flat mirror, MP phase mirror, MI imaging mirror ($f=27\,\text{cm}$). $\Sigma'$ first image plane.

Fig. 27. Mode converted RF waves measured using phase contrast imaging. a) Ion cyclotron wave in C-Mod. (modified from ref146). b) Kinetic Alfvén wave in TCA. c) Standing kinetic Alfvén with two nodes of oscillation.
Fig. 28. Dynamics of plasma density fluctuations in LHD. a) Raw phase contrast image with 1µs time resolution. b) Reconstructed upper (z>0) and (c) lower halves of plasma. White dashed and dotted arrows indicate the direction of the magnetic field at ρ=0.9 in upper and lower parts respectively. Consecutive pictures at a rate of 1 MHz reveal propagation into electron diamagnetic direction (magenta and cyan-blue arrows). d) k-spectra versus normalized radius in upper and lower (e) halves. From ref. [132].