INSTITUT FÜR PLASMAPHYSIK

GARCHING BEI MÜNCHEN

THE LASER AS A TOOL FOR PLASMA DIAGNOSTICS

Wilhelm H. Kegel

IPP/6/9 Oktober 1963

Presented at the Fifth Annual Meeting of the Division of Plasma Physics of the American Physical Society at San Diego (Nov. 6-9, 1963)

Die nachstehende Arbeit wurde im Rahmen des Vertrages zwischen dem Institut für Plasmaphysik GmbH und der Europäischen Atomgemeinschaft über die Zusammenarbeit auf dem Gebiete der Plasmaphysik durchgeführt.

The development of the laser has created a very intense and extremely monochromatic light source which is becoming an important tool in plasma physics. There are several experiments that are improved or are even made possible by the use of lasers. I do not intend to give a complete survey on all possible applications of lasers in plasma diagnostics. I just want to discuss some of them, especially those with which we are involved in the laboratories of the Max-Planck-Institut für Physik und Astrophysik and the Institut für Plasmaphysik at Munich.

1. Determination of Local Plasma Parameters by Light Scattering Experiment.

a) Theoretical Aspects.

We consider a volume V containing n_0 electrons (the mean electron density then is $n = n_0/V$) and $N_0 = n_0/Z$ positive ions of atomic charge Z . If a light beam passes through this volume and we ask for the light scattered, it is sufficient to consider only the light scattered by the electrons, because of the high mass ratio of ions to electrons. The intensity I_s of the radiation with the frequency Ω_2 and wavenumber $\vec{K}_2$ that is scattered into the solid angle do is given by

$$d I_s(\vec{K}_2, \Omega_2) = \sigma_e I_0(\vec{K}_1, \Omega_1) \langle |n(\vec{k}, \omega)|^2 \rangle do \qquad (1)$$

where $\vec{K}_2 = \vec{K}_1 + \vec{k}$ and $\Omega_2 = \Omega_1 + \omega$ (2a)

or $\vec{K}_2 = \vec{K}_1 - \vec{k}$ and $\Omega_2 = \Omega_1 - \omega$ (2b)

I_0 is the intensity of the incident radiation, σ_e the scattering crossection for a single electron, and $n(\vec{k}, \omega)$ the Fourier transform of the electron density, the brackets denote the ensemble average.

If we assume the incident light to be perfectly monochromatic, then the spectral distribution of the scattered light I_s represents the spectral distribution of the density fluctuations. By making some special assumptions, e.g. assuming thermal equilibrium, one can calculate the quantity $n(\vec{k}, \omega)$. If one is able to derive the corresponding data from light scattering experiments, one can decide by comparison whether these assumptions are fulfilled by the plasma under investigation or not. Principially one can detect all micro instabilities, as far as they lead to density fluctuations, with wavenumbers

$$0 \leq |\vec{k}| \leq 2 |\vec{K}_1| \tag{3}$$

by measuring the spectra of light scattered into different angles.

The problem of calculating the quantity $n(\vec{k}, \omega)$ for a plasma in thermal equilibrium has been treated during the last few years by many authors. I shall only mention the paper by E. E. SALPETER (1960), from which the formulas for the following discussion are taken.

In the case of classical Thomson scattering one assumes that the electrons, by which the light is scattered, are statistically independent. Then $\langle |n(\vec{k}, \omega)|^2 \rangle$ is proportional to n_0, the total number of electrons in the scattering volume. Whether this is a good approximation or not, depends on the parameter

$$\alpha = \frac{1}{|\vec{k}| D} \tag{4}$$

where

$$D = \sqrt{kT / 4\pi n e^2} \tag{5}$$

is the Debye length, n being the mean electron density and T the electron temperature. For $\alpha \ll 1$ the assumption made before is fulfilled. For $\alpha \gg 1$ collective effects must be taken into account, because the electrons cannot be considered as free particles over distances large compared to the Debye length. This has two consequences. On the one hand the total intensity of the light that is scattered into any given solid angle is reduced by a factor μ which is given by

$$\mu = \frac{1 + Z\alpha^2}{1 + (1+Z)\alpha^2} \tag{6}$$

The function $\mu(\alpha)$ varies between 1 and 0.5 and is shown in Fig. 1 for the case $Z = 1$. The essential decrease occurs in the region $0.2 < \alpha < 1.5$. This region corresponds to a variation in the electron density by a factor of 50. On the other hand there is also a very pronounced effect on the spectral distribution. The characteristic shape changes in the same region of α. For $\alpha \ll 1$ we have the classical case of a Gaussian profile, the width of which corresponds to the thermal velocities of the electrons. For $\alpha \gg 1$ the spectrum consists of a central line, the width of which is given by the thermal velocities of the ions, and two peaks symmetrically located about the central line at distances given approximately by

$$\omega_0 = \pm \sqrt{\omega_p^2 + \frac{3kT}{m} k^2} \tag{7}$$

which is the well known dispersion relation for plasma oscillations.

More precisely the spectral distribution is given by

$$I_s(\vec{K}_2, \Omega_2) \sim \langle |n(\vec{k}, \omega)|^2 \rangle \sim \frac{\overbrace{|1-G_i|^2 F_e(-\frac{\omega}{|k|})}^{I} + \overbrace{Z|G_e|^2 F_i(-\frac{\omega}{|k|})}^{II}}{|1-G_e-G_i|^2 |\vec{k}|} \tag{8}$$

where F_e and F_i are the Maxwellian distribution functions for the electrons and ions and

$$G_e(\omega) = -\alpha^2 [1 - f(x) + i\sqrt{\pi}\, x\, e^{-x^2}]$$

with $\quad x = \omega/\omega_e \quad ; \quad \omega_e = \sqrt{2 k^2 kT/m} \tag{9a}$

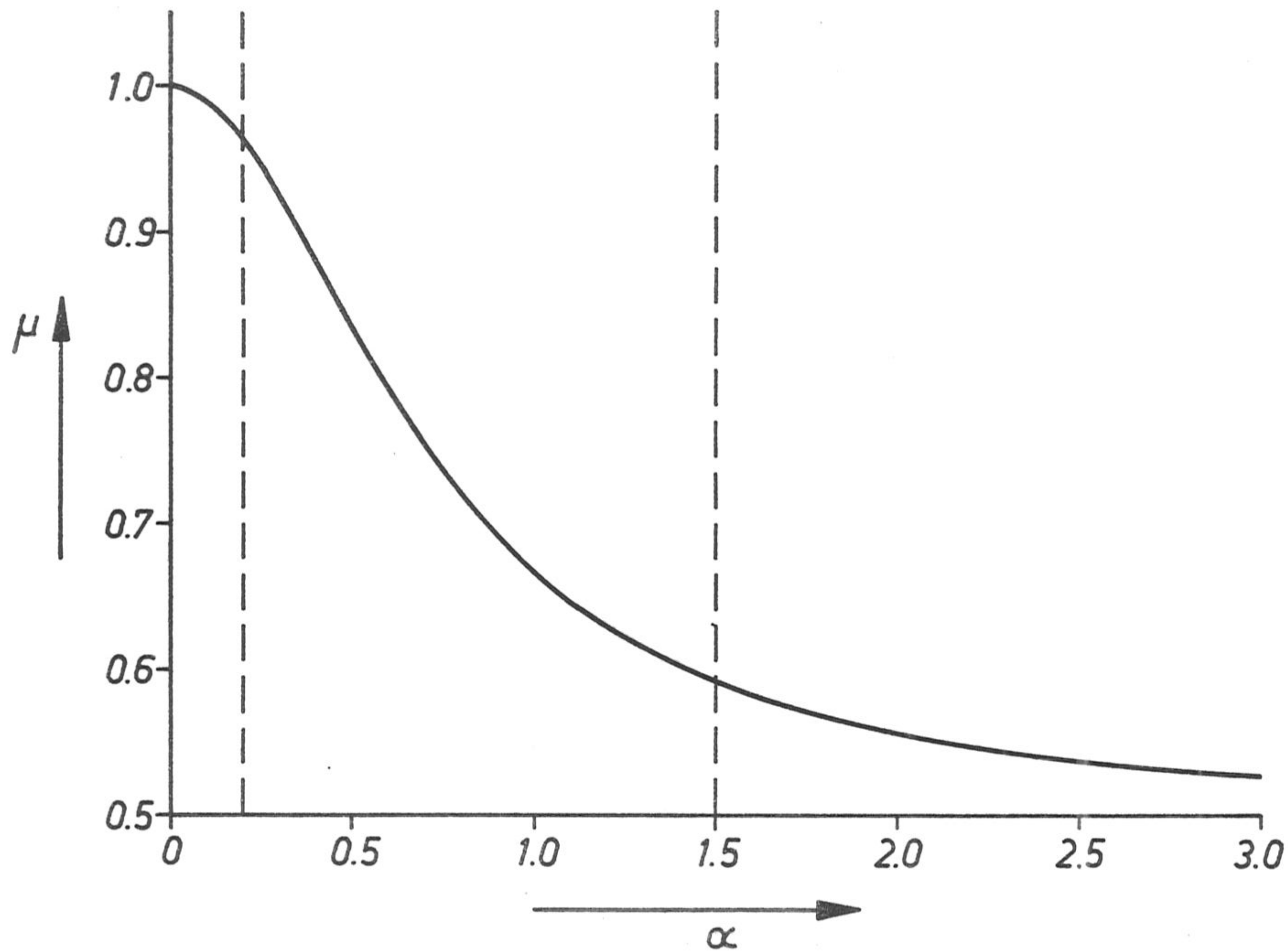

Fig. 1 *Reduction of the total scattering crossection as function of the parameter* α.

$$G_i(\omega) = -Z\alpha^2 [1 - f(y) + i\sqrt{\pi}\, y\, e^{-y^2}]$$

with $y = \omega/\omega_i$; $\omega_i = \sqrt{2k^2 \kappa T/M}$ (9b)

and

$$f(x) = 2x\, e^{-x^2} \int_0^x e^{t^2}\, dt \qquad (10)$$

Part I of the right hand side of eq. (8) represents the light scattered from electrons moving freely, while part II accounts for those electrons that move with the ions. For $\alpha \gg 1$ I describes the satellite peak while it describes the whole spectrum for $\alpha \ll 1$. The total energy contained in this part of the spectrum is proportional to $1/(1+\alpha^2)$.

Fig. 2-4 illustrate the situation. They show the calculated spectral distributions of light scattered from the beam of a ruby laser ($\lambda = 6943 Å$) for a scattering angle of 90^o and a temperature of 5 eV. In Fig. 2 the electron density is 10^{16} cm^{-3} corresponding to an α of 0.5. The profile already shows some deviations from a pure Gaussian profile (dashed curve), essentially in the line center. In Fig. 3 $n = 10^{17}$, $\alpha = 1.5$ and the type of the spectral distribution has changed completely. Fig. 4 shows an intermediate case with $n = 5 \cdot 10^{16}$ and $\alpha = 1.1$.

We now ask which parameters of the plasma may be derived from measurements of the spectrum of scattered light. In the case $\alpha \ll 1$ we get the electron temperature from the half width of the Gaussian profile and if reliable <u>absolute</u> measurements are performed, we can calculate the electron density. - In the case where α is about unity we can get more information already from the <u>relative</u> intensity distribution. By comparison with detailed numerical calculations one can deduce the temperature and density of the electrons and - if the central peak can be measured and resolved - the ion temperature. - In the case

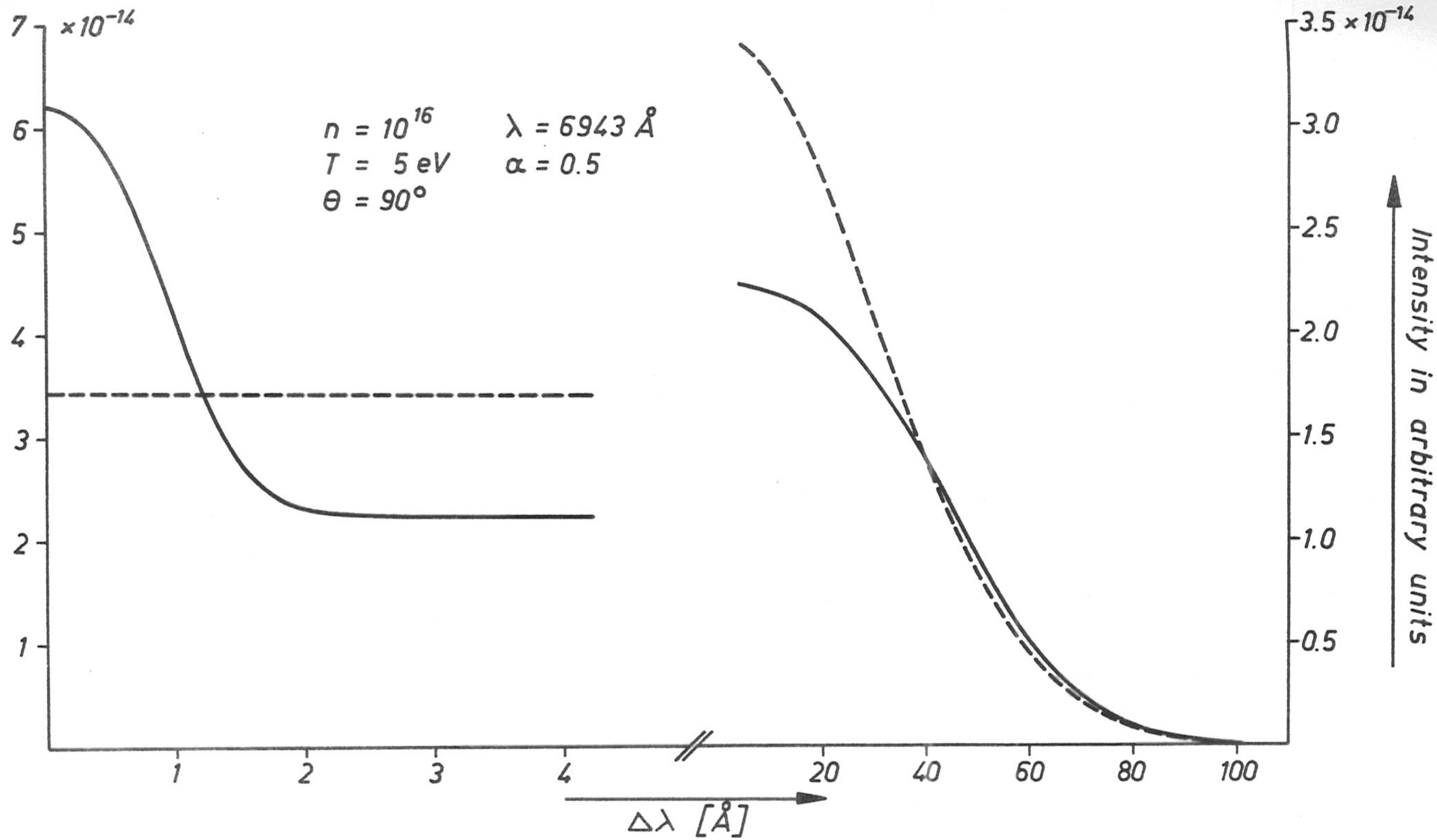

Fig. 2 —— Spectrum of scattered light
--- Gaussian profile (electrons assumed to be statistically independent)

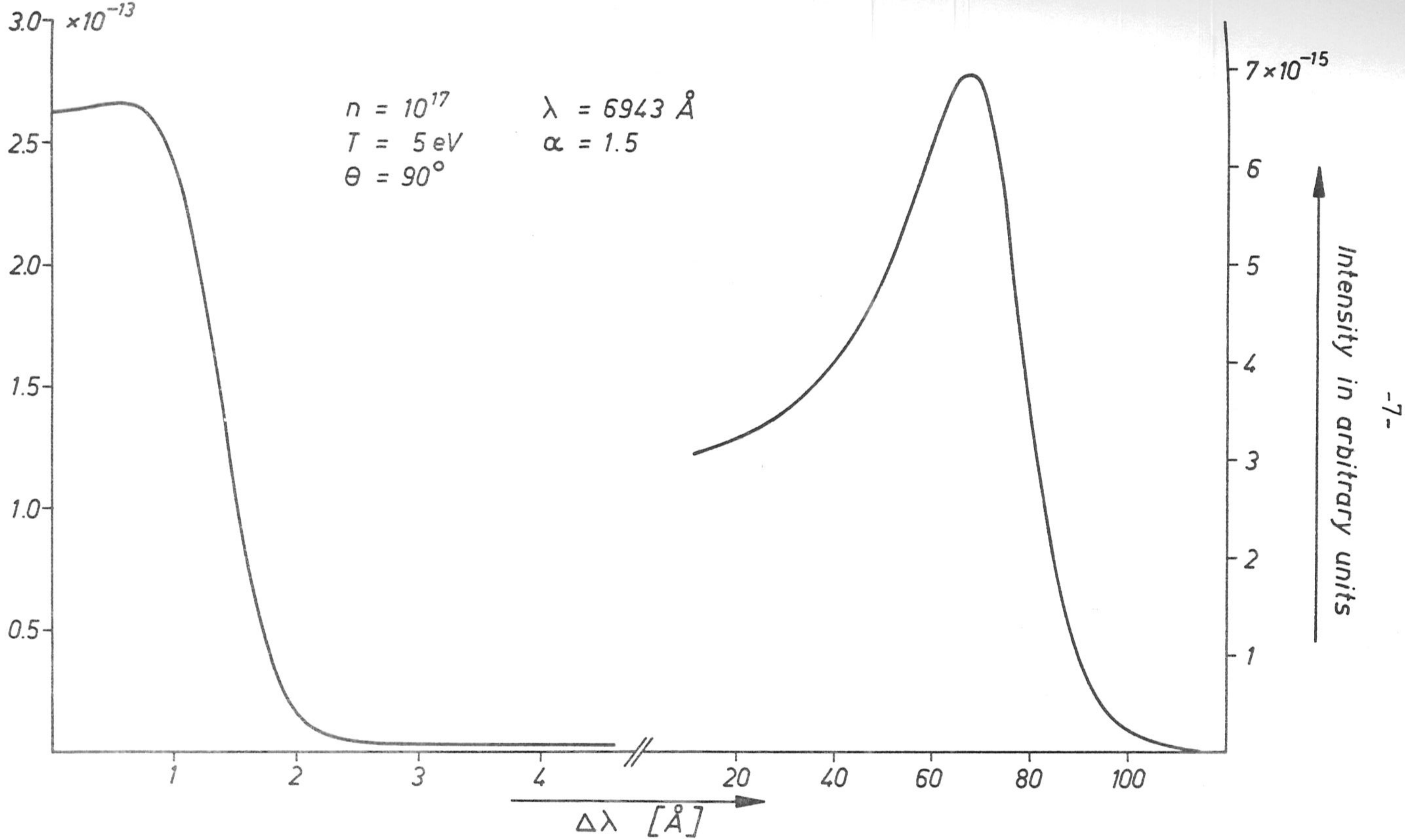

Fig. 3 Spectrum of scattered light.

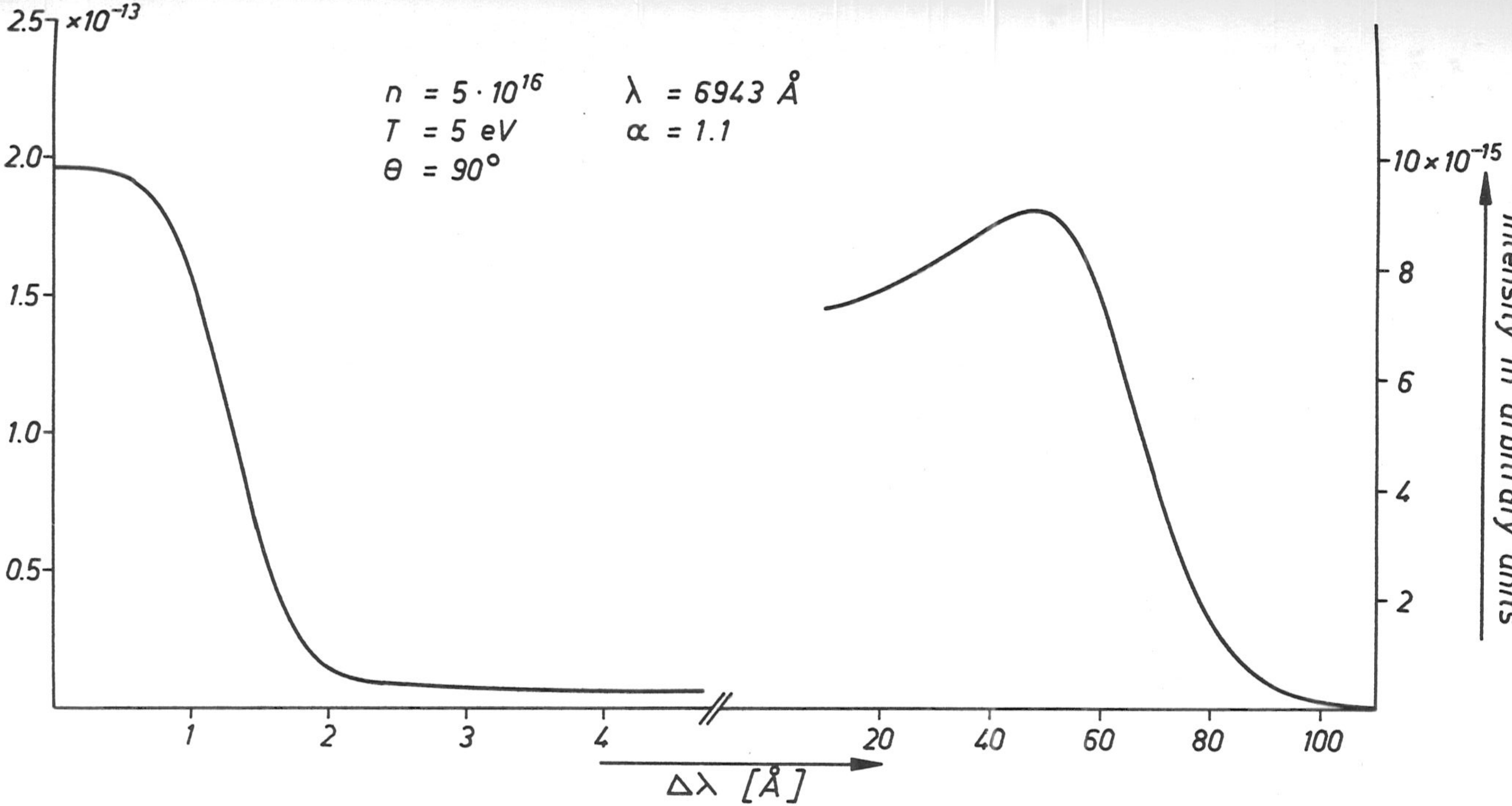

Fig. 4 Spectrum of scattered light

$\alpha \gg 1$ the situation becomes more unfavourable, because the total energy in the satellite peaks decreases rapidly with increasing α. In the framework of this theory the peaks become very narrow and as shown by V. GILINSKY and D. DUBOIS (1963) one has to take into account the broadening by collisions in order to get a more reliable theoretical profile. But principially one can deduce the electron density and temperature from the position of the peak and the ratio of the total energy contained in it to that of the central peak.

b. Experiments.

To my knowledge there are to date two groups which have succeeded in observing light scattered from a laboratory plasma.

One group is that of THOMPSON and FIOCCO at the M.I.T. In their first experiments they measured the light of a laser beam that was scattered by an electron beam. The electron density was about $5 \cdot 10^{9}$ cm^{-3} and the intensity ratio of the primary light to the scattered light was about 10^{18}. In order to get results under these extreme conditions they took advantage of the Doppler shift due to the high velocity of the electrons (i.e. they did not observe at right angles to the electron beam). In a further experiment they applied their techniques to a plasma of low electron density. In this case the measured signal was of about 10^{-15} of the input power, $n \approx 10^{13}$ cm^{-3}, $T \approx 1$ eV. This corresponds to an α of about 0.03, i.e. one would expect the spectrum of the scattered light to be Gaussian. Indeed one part of the measured spectrum had this shape, but there are still other features that are not yet explained.

The other group is that of E. FÜNFER, B. KRONAST and H.-J. KUNZE. They measured the light scattered from a high density Θ-pinch plasma (as order of magnitude, $n \approx 10^{17}$, $T \approx 5$ eV. In this case the intensity ratio of the scattered light to that of the primary laser beam was of the order of 10^{-11}. At these high densities an essential difficulty arises from the Bremsstrahlung, which increases proportional to n^2. In these experiments the self-radiation of the plasma was about 20 times larger than the intensity of the scattered light. This problem was solved by taking advantage of the fact, that the light scattered at right angles is always totally linearly polarized while the Bremsstrahlung is not. So one can eliminate the latter by a differential method.

Fig. 5 shows the experimental arrangement. The light of the ruby laser was focused on the centre of the Θ -pinch. The scattered light was observed at right angles. This provides a good localization of the plasma volume in which the measured light is scattered. - By means of a double refracting plate the emitted light was split into two beams, one polarized parallel to the polarization of the scattered light, the other perpendicular. After passing through the monochrometer the two components reached seperate photomultipliers, the signals of which were fed into a differential amplifier. The difference between these two signals was then observed on an oscilloscope. The differential amplifier was compensated for plasma light without incident laser light.

Owing to irregularities, there is always stray light produced by the primary laser beam at the windows and walls of the discharge tube entering the measuring apparatus. Thus, as in the experiment of THOMPSON and FIOCCO, it was not possible to measure in the line centre, but one had to take advantage of the fact that the light properly scattered by the plasma is broadened while the irregular "stray" light is not. Therefore the exit slit of the monochromator was located on one wing of the broadened line.

Fig. 6 shows the results of the measurements. Trace a is dI/dt of the Θ -pinch and trace b the primary laser output. Trace c is the differential signal when there is

Arrangement for Studies of Light - Scattering from a Plasma

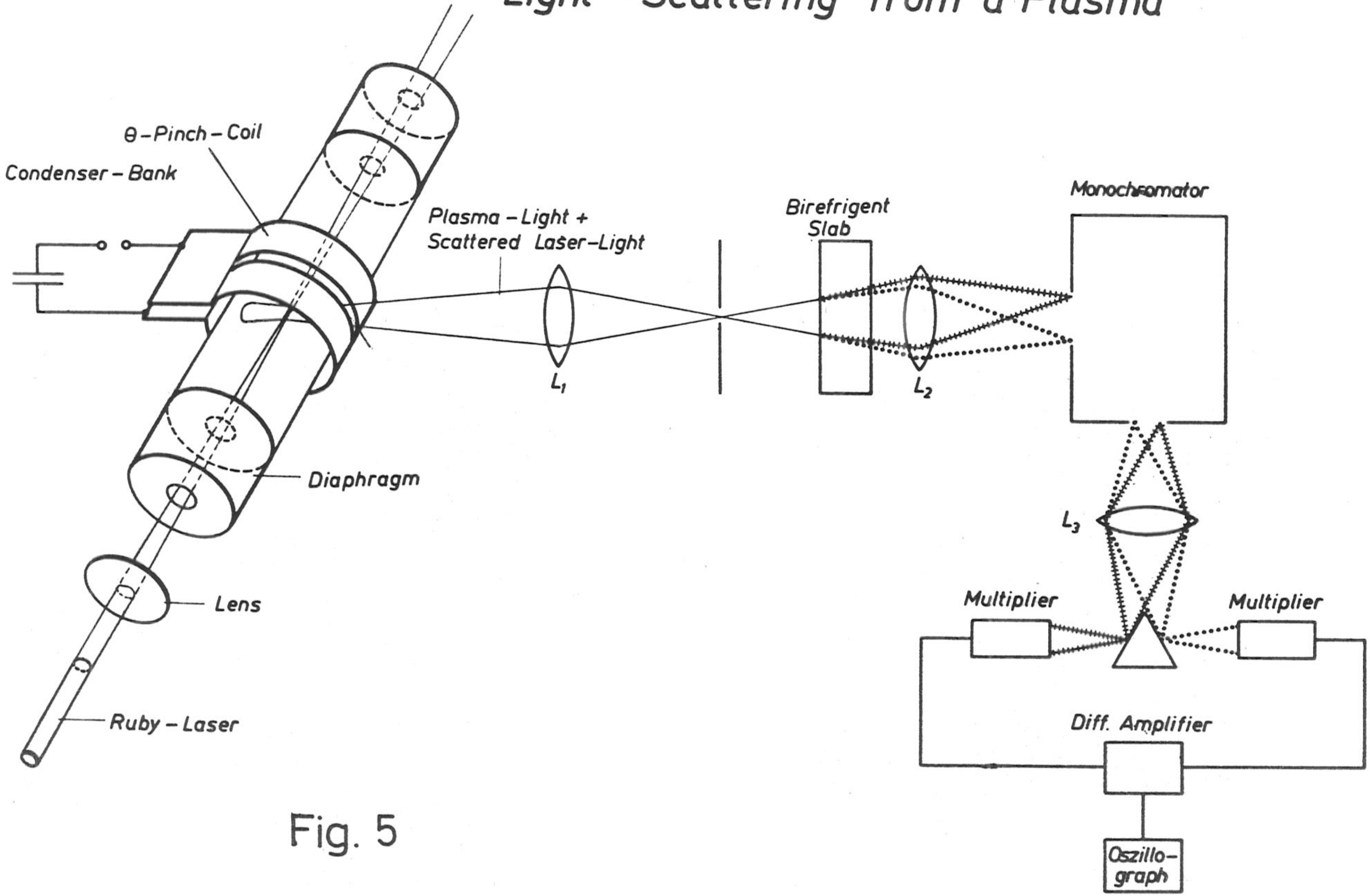

Fig. 5

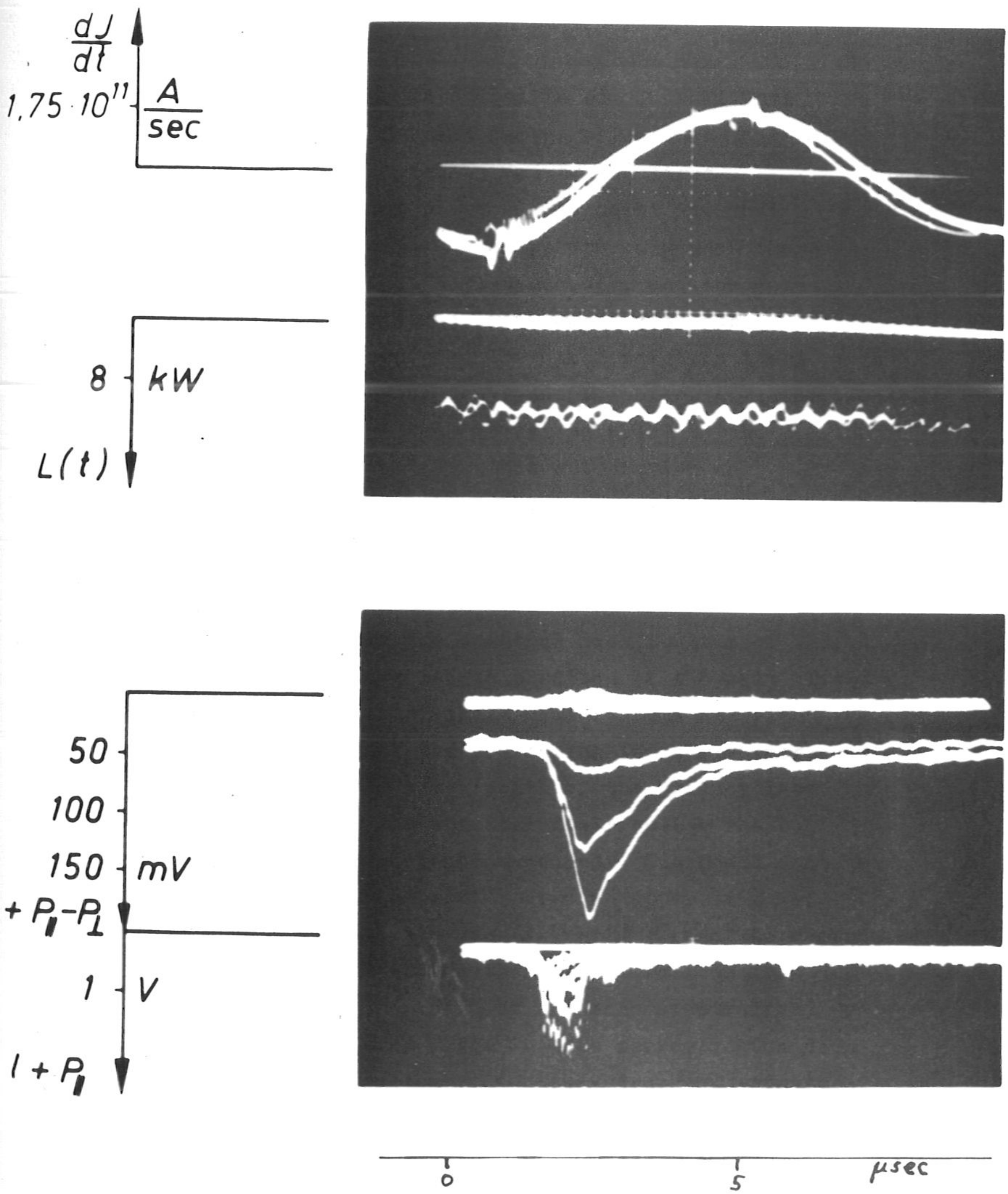

Fig. 6 Experimental results

only internal plasma radiation, and when the laser light is incident, trace d is registered. Finally, trace e corresponds to the total radiation contained in one beam.

I do not want to go into details but I shall just mention the two most important facts, which show that the measured signal d is really due to the light properly scattered by the plasma. One point is, that the time-varying part of trace d appears only when there are plasma and laser light simultaneously, and the onset and progress of the scattered light agree with our knowledge of the Θ -pinch. - The other point is the following: The E-vector of the linearly polarized laser radiation was rotated so that it was parallel to the direction of observation. For this orientation no scattered light should be observable. Indeed, in this case the scattering signal was reduced by more than a factor of ten.

In order to measure the spectral distribution of the scattered light, the wave length position of the exit slit was varied. The results of these measurements are given in Fig. 7. One sees that the scattering-intensity increases symmetrically toward the wave length of the primary laser light (6936.5 Å). But the spectral distribution was so narrow that no details could be measured. One could just estimate the half width or rather give an upper limit of this quantity. The estimate was 3.5 ± 3.5 Å and the upper limit consequently 7 Å . Furthermore it follows from the measurements that the ratio of the mean spectral intensity of the scattered light in this narrow region to the intensity in the spectrum 20 Å away from the primary wave length is larger than 75 .

If one compares these experimental results with the numerical calculations discussed before, one concludes that in this case $\alpha \geq 2$ which corresponds to a Debye length $D \leq 400$ Å

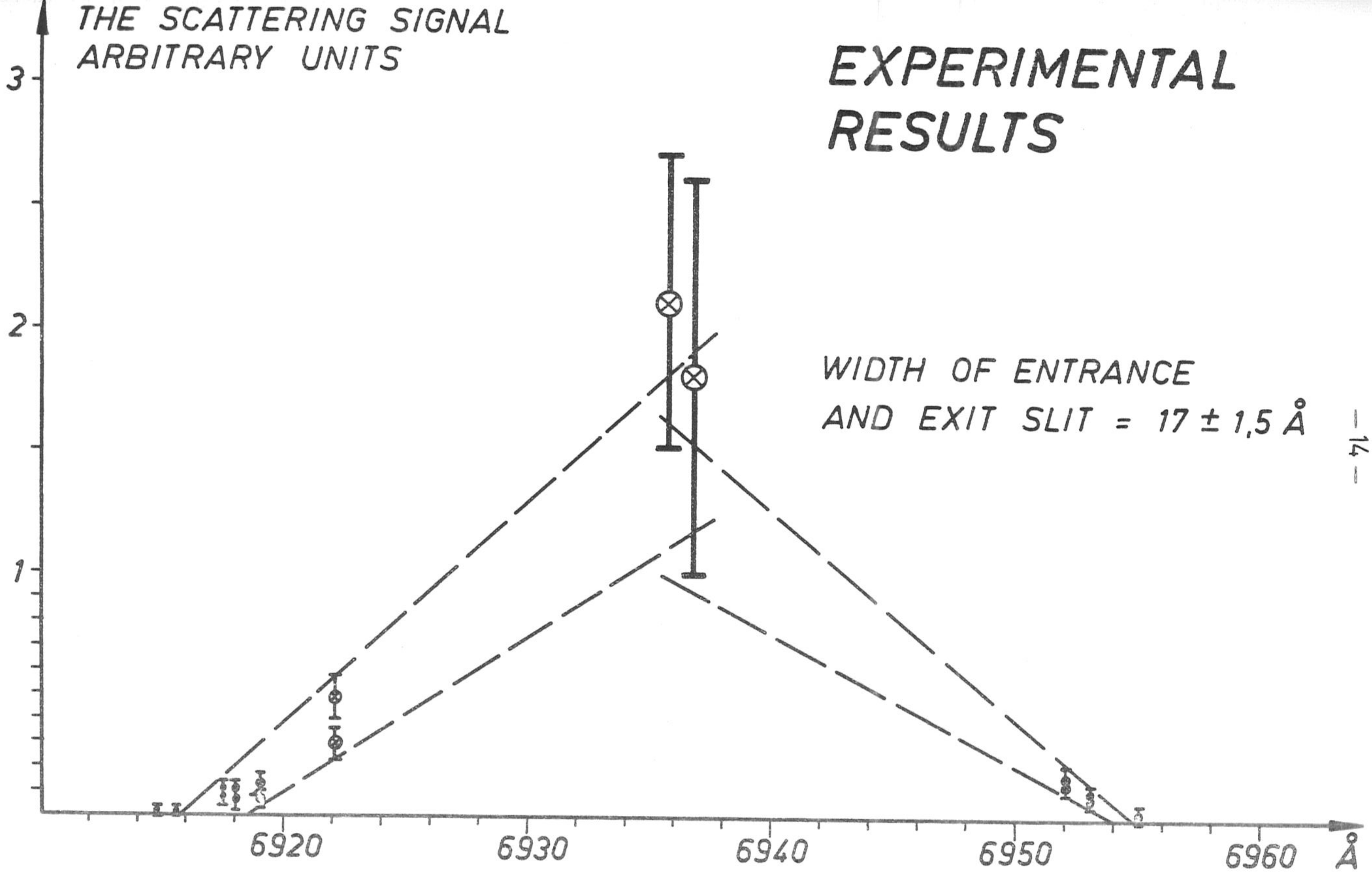
THE SCATTERING SIGNAL
ARBITRARY UNITS
EXPERIMENTAL
RESULTS
WIDTH OF ENTRANCE
AND EXIT SLIT = 17 ± 1,5 Å
3
2
1
6920
6930
6940
6950
6960
Å
POSITION OF THE SLIT CENTER

Fig. 7

In order to get more detailed results the experimental arrangement is going to be improved.

2. Light Mixing in a Plasma and Generation of the Second Harmonic.

Next, I want to make a few remarks on the problem of light mixing in a plasma, that is, the generation of the sum and difference frequency out of two fundamental frequencies.

One way to treat this problem is to extend the theory of light scattering. As pointed out before, the spectral distribution of the scattered light corresponds to the spectrum of the density fluctuations, and therefore the problem of calculating the spectrum of scattered light is essentially the problem of calculating the density fluctuations. In the previous section we were concerned with thermal fluctuations. We now ask for density fluctuations forced by an external electric field, e.g. the field of a laser beam.

The problem of calculating these fluctuations may be treated by means of the linearized Vlasov equation. One shows easily that within this approximation there arise no density fluctuations in a homogeneous plasma without magnetic field (one gets these only in a higher order approximation), because light waves are transverse. But if there is a magnetic field which is neither parallel to the driving electric field $\vec{E}_{ex}$ nor to its wave vector $\vec{k}_0$, the motion of the electrons oscillating in the electric field is influenced such that it obtains a longitudinal component (see Fig. 8), which do give rise to density fluctuations. One would expect the largest effect for the case $\vec{k}_0 \perp \vec{B} \perp \vec{E}_{ex}$, since the effect is due to the force $\vec{v} \times \vec{B}$. This case was treated in an earlier paper (W.H. KEGEL 1963) under the assumption that $\vec{B}$ is time-constant and homogeneous. In the case that the frequency of the forcing field $\omega_0 \gg \Omega_e$, Ω_e being the cyclotron frequency of the electrons, the result was:

$$\langle |n(\vec{k}_0, \omega_0)|^2 \rangle = E_{ex}^{0\,2} B^2 \frac{n_0^2}{\omega_0^4} \left(\frac{e}{mc}\right)^4 \cdot 4 \left(\frac{(2\pi)^2 - 6}{6}\right) \tag{11}$$

$$= E_{ex}^{0\,2} B^2 \frac{n_0^2}{\omega_0^4} \cdot 1.11 \cdot 10^{31} ,$$

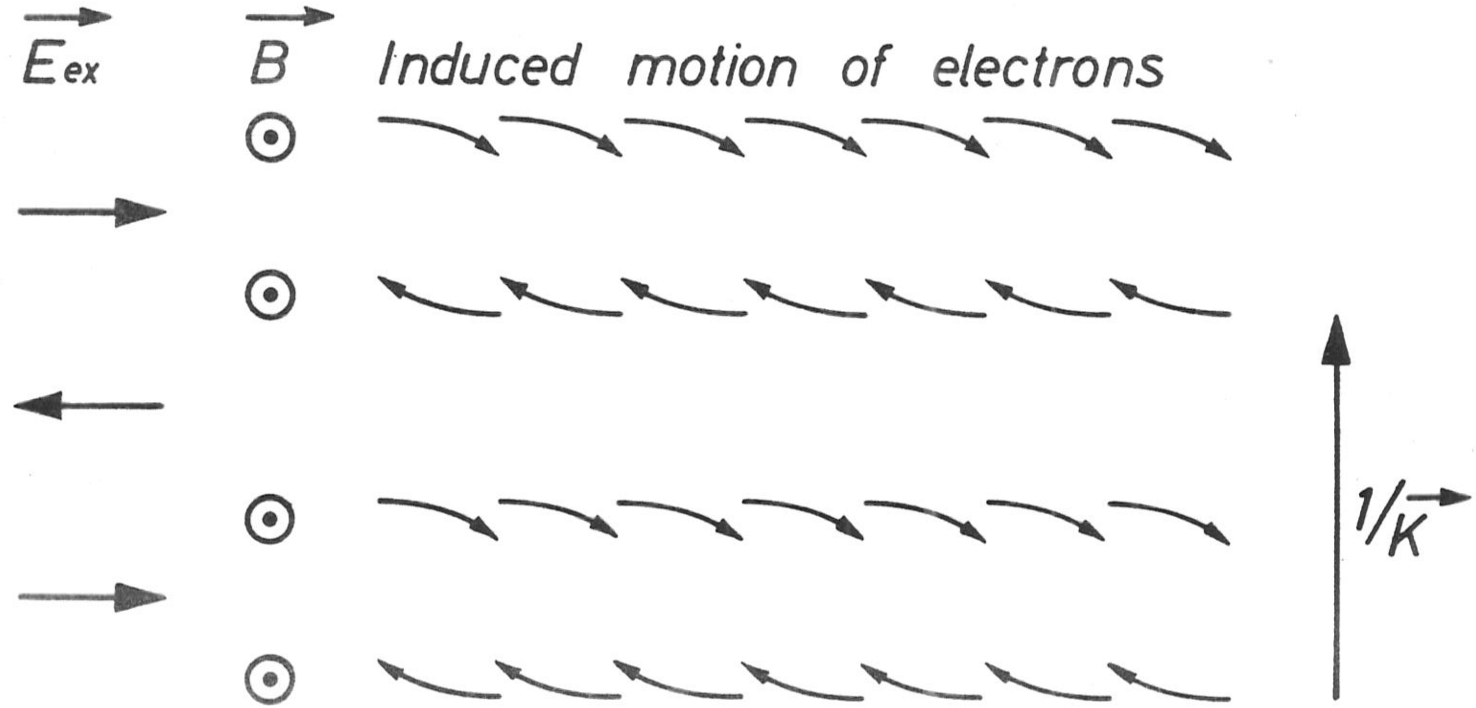

Fig.8 Qualitative explanation of induced density fluctuations

where all quantities are to be taken in c.g.s. units.

These fluctuations have the frequency and wave number of the electric field by which they are forced. If now the light of a second beam is scattered by these fluctuations, we see from eq. (2) that the scattered light has the sum and difference frequency of the two beams. That is to say light mixing occurs, if the condition (2) for the K's is also fulfilled.

It should be emphasized that eq. (1) combined with (11) also describes the generation of the second harmonic (take $\omega_o = \Omega_1$).

For an electron density of $n = 10^{17}$ cm^{-3}, a volume of 10^{-2} cm^3, a frequency of 2.7×10^{15} (ruby laser), an electric field of 100 e.s.u. or 3×10^4 V/cm, and a magnetic field of 10^4 Gauss, one calculates from eqs (11) and (1) that the intensity of the light that is scattered from a second beam by the forced fluctuations is of the same order of magnitude as the total light scattered by the thermal fluctuations into a solid angle of about 10^{-2}. That is to say it should be possible to detect the light mixing effect experimentally.

One essential property of this "scattered" light may be seen directly from eq. (2). If one assumes that the refraction index is unity (i.e. plasma frequency $\omega_p \ll \Omega$) and that there is no dispersion, the following relation must be fulfilled:

$$\frac{\Omega_1}{|\vec{K}_1|} = \frac{\Omega_2}{|\vec{K}_2|} = c_o \tag{12}$$

c_o is the velocity of light in vacuum. Since this relation holds also for the forced fluctuations (i.e. $\omega_o / |\vec{k}_o| = c_o$), it follows from the condition (2) that the mixing effect occurs only if the light beams are parallel * for the K's are vectors while the Ω's are scalar quantities.

---

* This is a correction to the earlier statement (W.H. KEGEL 1963) that the difference frequency were generated in the case of antiparallel beams.

The case becomes more difficult if one takes the dispersion into account. In the region where the dispersion relation is given by

$$c_0^2 K^2 = \Omega^2 - \omega_p^2 \qquad (13)$$

(i.e. $\Omega \gg \Omega_e$) it can easily be shown that the condition (2) cannot be satisfied. The situation is similar to that in the case of the generation of harmonics in crystals, where the intensity is a periodic function of the crystal thickness, for an essential part of the electric field cancels because of interference. If one takes into account this effect, it turns out that the right-hand side of eq. (1), which gives the intensity of the generated frequency is to be multiplied by a factor of

$$\frac{\sin^2(\Delta K \cdot L)}{(\Delta K \cdot L)^2} \qquad (14)$$

where L is the path length of the laser beam through the plasma and

$$\Delta K = K_1 \pm k_0 - \frac{\Omega_1 \pm \omega_0}{c(\Omega_1 \pm \omega_0)} \qquad (15)$$

The two light beams are still assumed to be parallel, $c(\Omega)$ being the phase velocity of a light wave in the plasma. In the special case of the generation of the second harmonic it follows from (13) and (15) that

$$\Delta K = \frac{\Omega_1}{c_0}\left(\sqrt{1-\frac{\omega_p^2}{\Omega_1^2}} - \sqrt{1-\frac{\omega_p^2}{4\Omega_1^2}}\right) \qquad (16a)$$

and for $\omega_p \ll \Omega_1$

$$\Delta K = \frac{\Omega_1}{c_0}\,\frac{3}{8}\,\frac{\omega_p^2}{\Omega_1^2} \qquad (16b)$$

being the fundamental frequency. (For $\Omega_1 = 2.7 \cdot 10^{15}$ and $n = 10^{16}$ ΔK is of the order 0.1 cm^{-1}).

In the case ΔK or $L \rightarrow 0$ the factor (14) goes to unity, i.e. for small dimensions one may neglect the influence of the dispersion.

The problem of light mixing and generation of harmonics has been treated in a more rigorous way by A. SCHLÜTER (1963). He treats the interaction of light with the plasma by means of the non linear plasma equations. In this treatment he takes the magnetic field of the light wave into account. I merely mention some of his results.

a) In the case of a homogeneous plasma with external magnetic field the following relation holds for the intensity of the second harmonic generated by a laser beam:

$$I(2\Omega_1) \sim B_0^2\, I^2(\Omega_1)\, \Omega_1^{-4} \left(\frac{e}{mc}\right)^4 \qquad (17)$$

in agreement with eq. (11) combined with (1) and (14).

b) The second harmonic generated by a laser beam of finite extent propagates in the surrounding undisturbed plasma in a direction which is inclined by an angle ρ to the primary laser beam. For $\omega_p \ll \Omega_1$, ρ is given by

$$\sin \rho = \sqrt{3}\, \frac{\omega_p}{2\Omega_1} \qquad (18)$$

This provides a method of determining the density by measuring the angle ρ.

c) A. SCHLÜTER (1963) gets a very surprising result when he treats the case of two parallel light beams traversing a plasma without external magnetic field assuming the frequency difference of the two light beams to be the plasma frequency. He treats the non linear plasma equations by making the ansatz that the electric field within the plasma consists essentially of three parts with the frequencies Ω_1, $\Omega_2 = \Omega_1 + \omega_p$, and ω_p with no restrictions to the amplitudes. He derives the "non linear dispersion relation"

$$(c_0^2 K_1^2 - \Omega_1^2) K_2 = (c_0^2 K_2^2 - \Omega_2^2) K_1 \qquad (19)$$

and gets the final result that in the collision free case longitudinal plasma oscillations are excited with an amplitude which is independent of the amplitudes of the primary fields. This amplitude is given by

$$\vec{E}_z(\omega_p) = \frac{8\pi n e}{\vec{K}} \tag{20}$$

If collisions are taken into account and it is assumed that the amplitude of the excited plasma oscillations is small compared to the value given by (20), the result is for arbitrary Ω_1 and $\Omega_2 \gg \omega_p$

$$|E_z(\Omega_2-\Omega_1)|^2 = \frac{e^2}{m^2}\,\frac{\omega_p^4}{\Omega^6}\cdot\frac{(\Omega_2-\Omega_1)^2}{4}\,K^2\cdot\frac{|E_1|^2\cdot|E_2|^2}{[(\Omega_1-\Omega_2)^2-\gamma^2-\omega_p^2]^2+4(\Omega_1-\Omega_2)^2\gamma^2} \tag{21}$$

γ being the collision frequency, $\Omega = (\Omega_1+\Omega_2)/2$; $\vec{K} = (\vec{K}_1+\vec{K}_2)/2$.

3. Measurement of Electron Densities by Means of the Faraday Effect.

If linear polarized light is transmitted through a plasma parallel to a magnetic field, the plane of polarization is rotated by an angle α. In the case $\Omega \gg \omega_p, \Omega_e$ (Ω = frequency of the light, ω_p = plasma frequency, Ω_e = electron cyclotron frequency) α is given for a homogeneous plasma by

$$\alpha = \frac{1}{2}\,\frac{\omega_p^2}{\Omega^2}\cdot\frac{\Omega_e}{c}\,Z = \frac{2\pi e^3}{m^2c^2}\,\frac{n}{\Omega^2}\,B\cdot Z \tag{22}$$

(B = magnetic field strength, Z = length of the plasma). If the length Z of the plasma is known, one determines $n\cdot B$ in measuring the angle α. Assuming $\Omega = 2.7 \times 10^{15}$ sec^{-1} (ruby laser), $n = 5\cdot 10^{16}$ cm^{-3}, $B = 40\ k\Gamma$ and $L = 20$ cm one calculates from (22) $\alpha = 0.3^\circ$. The rotation of the plane of polarization can be measured easily down to an angle of 0.1° with an accuracy of about 20%.

The corresponding experiments were performed by P.H. GRASSMANN and H. WULFF. Their experimental arrangement is shown in Fig. 9. The plasma is that of an Helium discharge, which in detail is described by J. DURAND (1963).

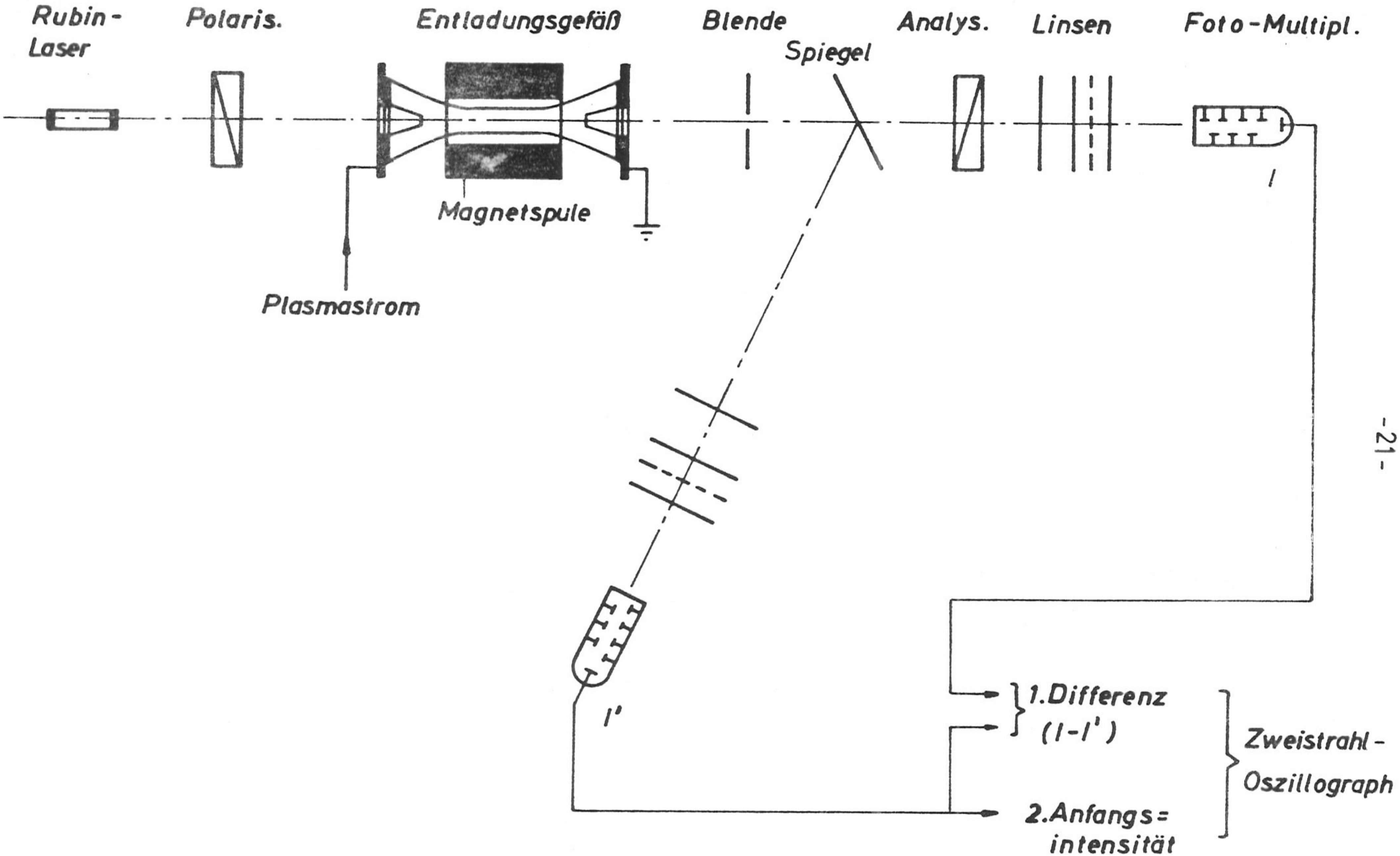
Rubin-
Laser
Polaris.
Entladungsgefäß
Magnetspule
Plasmastrom
Blende
Spiegel
Analys.
Linsen
Foto-Multipl.
I
I'
1.Differenz
(I-I')
2.Anfangs=
intensität
Zweistrahl-
Oszillograph

Fig.9

During a time of the order 1 m sec the plasma is stationary and homogeneous in Z - direction (except in a small region at the windows). This is also true for the magnetic field. The polarization of the laser light is improved by a Glan - Thompson - prisma. Then it passes the plasma and the analyser, which is another Glan - Thompson - prisma and finally is measured by a photomultiplier, which gives rise to the signal I. One part of the total light is branched off before entering the analyser and is measured separately by a second multiplier, giving rise to the signal I', which is a measure for the primary intensity I_0. The self-radiation of the plasma is eliminated from the laser light by means of interference filters. This possibility is the essential advantage of the laser compared to normal light sources.

The signals I and I' are fed into a differential amplifier the signal of which is registered simultaneously with I' on a two beam oscilloscope.

The intensity of light that passed two polarizers the transmission direction of which differ by an angle $(\pi/2 + \alpha)$ is given by

$$I = I_0 \cos^2(\pi/2 + \alpha) + I_R = I_0 \sin^2\alpha + I_R \qquad (23)$$

where I_0 is the intensity for $\alpha = \pi/2$ (the transmission direction of the two polarizers are parallel) and I_R is the intensity for $\alpha = 0$ (direction of extinction).
For $\alpha << 1$ we may write

$$\alpha^2 = \frac{I - I_R}{I_0} \qquad (24)$$

Because I_R is proportional to I_0 and independent of α, it may be eliminated experimentally. For this purpose the signals I and I' are fed into the differential amplifier, which is compensated for $\alpha = 0$ (no plasma). Then

$$\alpha^2 = \frac{I - I'}{I^0} \qquad (25)$$

and the accuracy of the measurements is limited by the precision of the compensation for $\alpha = 0$.

On the oscilloscope the quantities $I - I'$ and I' are measured. The ratio of these two quantities is, according to (25), proportional to α^2 . Calibration was performed by rotating the analyser by known angles α and measuring the quantities $I - I'$ and I'.

In this experiment the electron density n was deduced from the measured angles α to be

$$n = 5 \cdot 10^{16} cm^{-3} \pm 20\% \tag{26}$$

the magnetic field B being known. This density n was compared with that which one deduces from the line profile of the He I - line $\lambda = 4471$ Å . The latter was $4.5 \times 10^{16} cm^{-3}$; i.e. in good agreement with the first.

The method of determining the electron density by means of the Faraday effect is most suitable for high electron densities. A disadvantage of the method is, that only the integral of $n \cdot B$ along the path of the laser beam is measured. I. e. in order to determine the electron density itself, n and B have to be homogeneous along the path of the laser beam and B has to be known. But the method is a rather simple one.

4. Further Applications of the Laser.

I wish to mention some further applications of the laser in plasma diagnostics, without going into details.

The method of determining the electron density of a plasma by means of a Mach-Zender interferometer is well known. It can essentially be improved by using a laser as light source, as was done by G.A. SAWYER, W.E. QUINN and F.L. RIBE (1963).
The essential advantages are:
a) The high intensity easily overcomes the self-luminosity of the plasma. b) Due to the short duration of the laser pulse, there is a good time resolution (~ 80 nsec) . c) The high coherence gives sharp fringes to very high order.

d) In the case of high density gradients the fringes are still sharp and errors in the evaluation are small, due to the small divergence of the expanded beam.

Another method that is improved by the application of the laser, is the shadowgraph technique. This was also done at Los Alamos by F.C. JAHODA, W.E. QUINN and F.L. RIBE (1963) and at Frascati by U. ASCOLI-BARTOLI and coworkers. The essential advantage of the laser is again the small divergence of the expanded beam and the possibility of eliminating the self-luminosity of the plasma by means of an interference filter.

5. Limitations.

Finally I want to make a few remarks on the limitations of these methods.

In all diagnostic methods one wants to have a tool by which the plasma parameters are measured without disturbing the plasma itself. In the case of optical methods one measures the influences the plasma has on a light beam passing through the plasma. At the same time one makes the basic assumption that the state of the plasma is not altered seriously by the light. On the other hand one actually takes the most intense light source one knows, focusing the energy in some cases to the smallest area that is possible. That is to say, if one is going to determine quantitatively plasma parameters from the interaction of a laser beam with the plasma, one has to bear in mind the above assumption and check whether it is fulfilled or not.

As an illustration I just want to give an example. One can calculate the energy that is absorbed by a plasma from a laser beam. If one takes a 1 MW laser with a pulse duration of 10^{-7} sec which is focused to an area of 1 mm^2 into a plasma with $n = 10^{17}$ and $T = 5$ eV, the absorbed energy is about 27% of the thermal energy.

The amplitude of the kinetic energy which an electron receives from oscillating in the laser field is given by:

$$E_{KIN} = \frac{e^2 E^2}{2 m \Omega^2} = \frac{1}{n} \frac{\omega_p^2}{\Omega^2} \frac{E^2}{8\pi} \qquad (27)$$

For a field of 10^6 e.s.u. or 3×10^8 V/cm this quantity is of the order of 6 eV.

This shows that there might be cases in which one has to take into account such influences. Principially for the reduction of the measured data one has to use a self consistent theory which takes into account the influences of the light on the plasma. - In some extreme cases, when the duration of interaction is short compared to the mean time between two collisions - which might be the case in the application of very high power Q - spoiled lasers - these influences might be considerably non isotropic.

References

ASCOLI-BARTOLI, U., S. MARTELLUCCI and E. MAZZUCATO
Proc. VI Conf. Ion. Phen. Gas. Paris 1963

BOWLES, K.L.: Phys. Rev. Lett. 1, 454 (1958)

DURAND, J.: Zs. f. Naturf. 18a, 281 (1963)

FIOCCO, G. and E. THOMPSON: Phys. Rev. Lett. 10, 89 (1963)

FÜNFER, E., B. KRONAST and H.-J. KUNZE: Phys. Lett. 5, 125 (1963)

FÜNFER, E., W.H. KEGEL, B. KRONAST and H.-J. KUNZE:
Proc. VI Conf. Ion. Phen. Gas. Paris 1963

GILINSKY, V. u. D. DUBOIS: ibid

GRASSMANN, P.H. u. H. WULFF: ibid

JAHODA, F.C., W.E. QUINN and F.L. RIBE: LAMS - 2916 (1963)
also this meeting

KEGEL, W.H.: Proc. VI Conf. Ion. Phen. Gas. Paris 1963

SALPETER, E.E.: Phys. Rev. 120, 1528 (1960)

SAWYER, G.A., W.E. QUINN and F.L. RIBE: LAMS - 2916 (1963)
also this meeting

SCHLÜTER, A.: private communication

THOMPSON, E. and G. FIOCCO: Proc. VI Conf. Ion. Phen. Gas.
Paris 1963