Modeling of Thomson Scattering Spectra in High-Z, Laser Produced Plasmas

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Modelling of Thomson scattering spectra in high-Z, laser produced plasmas

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Abstract

Theoretical calculations of a Thomson scattering cross section and dynamical form factors are presented for high-Z laser produced inhomogeneous plasmas. Relevance of these results to astrophysical plasmas is pointed out. Comparisons with recent experimental observations are discussed with emphasis on the effects of plasma inhomogeneity, ion-ion collisions and non-Maxwellian distribution functions.

I. INTRODUCTION

Thomson scattering (TS) is an important radiative process involving incoherent scattering of electromagnetic waves by free electrons in plasmas. TS has several applications in astrophysics. In stellar interiors scattering from free electrons is one of the basic processes contributing to the radiative opacity [1]. It has been shown recently that by properly accounting for the plasma response in the calculations of the TS cross-section one can reduce the solar opacity [2] and hence change predictions for the core parameters. This has been done by including broadening of the Langmuir wave resonance due to the Doppler effect and particle collisions.

With the advent of laser plasma interaction experiments in the regime of parameters of interest to astrophysics, TS becomes extremely important as the accurate plasma diagnostics. One of the main applications of the laboratory astrophysics is benchmarking of numerical codes used in modelling of astrophysical processes. This can be only accomplished with the help of reliable laboratory plasma diagnostic techniques such as the TS.

Since early days of laser plasma interaction physics, TS has been perfected as the diagnostic of enhanced levels of fluctuations due to parametric instabilities. Larger scattering powers from non-thermal fluctuations are easier to detect and many of these measurements have fundamental importance elucidating nonlinear saturation processes which are valid in the wide range of laser produced plasmas. Scattering from stable plasmas for the purpose of defining electron density, temperatures of electrons and ions, flow velocities, ionization levels, particle distribution functions and thermal transport has been also conducted. Several interesting experiments have characterized low temperature laser produced plasmas [3,4], x-ray laser plasmas [5] and the inertial confinement fusion targets [6].

In this paper we discuss some theoretical results related to the calculations of the TS cross section in inhomogeneous plasmas and dynamical form factors which are valid under different physical conditions. Such theories and improved experimental techniques lead to further development of the TS as the most versatile and accurate diagnostic of laser produced

plasmas.

In Section II we derive an expression for the TS cross section in inhomogeneous plasmas. Section III summarizes different results for the dynamical form factor: in collisionless inhomogeneous plasmas, plasmas with frequent ion-ion collisions and in plasmas described by the non-Maxwellian distribution functions. Conclusions and discussion are presented in Sec. IV.

II. THOMSON SCATTERING FROM INHOMOGENEOUS PLASMAS

Hot and dense plasmas produced by powerful lasers are almost always inhomogeneous. This is also a desirable feature in the laboratory astrophysics applications, where expanding corona or blow-off plasma is close to a "real" astrophysical conditions [cf. Ref. [7]]. Following the standard procedure (cf. e.g. [8]), we first derive the scattering cross section for the incoherent Thomson scattering in such plasmas. The incident pump is modelled by,

$$\mathbf{E}_0(\mathbf{r}, t) = E_0(\mathbf{r})\hat{\eta}_0 \cos(\omega_0 t - \mathbf{k}_0 \cdot \mathbf{r}). \quad (1)$$

We have assumed that within the scattering volume the probe is well approximated by the cylindrical laser beam with the prescribed polarization $\hat{\eta}_0$. The scattered radiation field is described by the vector potential $\mathbf{A}(\mathbf{r}, t)$ which at the large distance, r , from the plasma is approximated by the following expression,

$$\mathbf{A}(\mathbf{r}, t) = \frac{1}{cr} \int_V d^3r' \int dt' \delta\mathbf{J}(\mathbf{r}', t') \delta(t' - t + r/c - \hat{n} \cdot \mathbf{r}'/c), \quad (2)$$

where $\mathbf{r}'$ varies within the scattering volume V and the separation of scales, $r \gg r'$, has been used in the approximate definition of the retarded time. The fluctuating nonlinear current $\delta\mathbf{J}$ is the result of a coupling between the Thomson probe (1) and the electron density fluctuations. From the solution to kinetic equation in the quasilinear approximation one obtains

$$\delta\mathbf{J}(\mathbf{r}, t) = \frac{e^2}{m_e} \int \frac{dz}{2\pi} \exp(-izt) \frac{1}{z} \delta\mathbf{S}(\mathbf{r}, z), \quad (3)$$

where the fluctuating source $\delta\mathbf{S}$ is the Laplace transformed product of the fluctuating electron density, δN^e , and the laser pump (1),

$$\delta\mathbf{S}(\mathbf{r}, t) = \delta N^e(\mathbf{r}, t) \mathbf{E}_0(\mathbf{r}, t). \quad (4)$$

In the derivation of Eq. (3) we have assumed that the plasma is almost transparent to the probe laser, i.e. $\omega_0 \gg \omega_p$ and $\omega_0 \gg \mathbf{k} \cdot \mathbf{v}$, where $\mathbf{k}$ is the wave vector of a collective plasma response and $\mathbf{v}$ is the characteristic particle velocity. The fluctuating current (3) provides a source of the radiation at the frequency ω_1 , which is close to the probe frequency ω_0 . A large value of ω_1 simplifies description of a plasma response to scattered radiation and the propagation of a laser light inside the plasma. In the radiation zone, fields are approximated by the following expressions [9]

$$\mathbf{B}_1 = \frac{1}{c} \hat{n} \times \dot{\mathbf{A}}, \quad \mathbf{E}_1 = \frac{1}{c} \hat{n} \times (\hat{n} \times \dot{\mathbf{A}}), \quad (5)$$

where $\hat{n}$ is the unit vector along the direction of $\mathbf{r}$. Suppose that the instrument measures the scattered electric field along the direction of a unit vector $\hat{\eta}_1$. To calculate energy flux of the scattered radiation at the distance r from the plasma we first derive the statistically averaged value of a Poynting vector

$$\langle \mathcal{P} \rangle = \frac{c}{4\pi} (\hat{\eta}_0 \cdot \hat{\eta}_1)^2 \frac{1}{c^2} \langle \dot{\mathbf{A}}^2(\mathbf{r}, t) \rangle. \quad (6)$$

Using formulas (2), (3) and (5) in Eq. (6) we find the following expression

$$\begin{aligned} \langle \mathcal{P} \rangle = & \frac{c}{4\pi} (\hat{\eta}_0 \cdot \hat{\eta}_1)^2 \frac{1}{4r^2} r_0^2 \int_V d^3 r' \int_V d^3 r'_1 \int \frac{d\omega}{2\pi} E_0(\mathbf{r}'_1) E_0(\mathbf{r}') \langle \delta N^e(\mathbf{r}') \delta N^e(\mathbf{r}'_1) \rangle(\omega) \\ & \left\{ \exp \left[-i \frac{\omega + \omega_0}{c} \hat{n} \cdot (\mathbf{r}' - \mathbf{r}'_1) + i \mathbf{k}_0 \cdot (\mathbf{r}' - \mathbf{r}'_1) \right] + c.c. \right\}, \end{aligned} \quad (7)$$

where $r_0 = e^2/m_e c^2$. The symbol $\langle \delta N^e(\mathbf{r}') \delta N^e(\mathbf{r}'_1) \rangle(\omega)$ defines Fourier transformed in time density correlation function. We have used the relation

$$\langle \delta N^e(\omega) \delta N^e(\omega') \rangle = 2\pi \langle \delta N^e \delta N^e \rangle(\omega) \delta(\omega + \omega'),$$

which is valid in the stationary plasmas.

Details of the derivation presented above are necessary to define limits of validity of the Thomson scattering cross-section in inhomogeneous plasmas. We will consider an example of the experimental data [10,12] where the inhomogeneity of the plasma flow and density play an important role in the calculation of a cross-section. One can easily account for this effect if the spatial scale of density and flow variations is much longer than the correlation length of electron fluctuations. In stable plasmas the correlation length of a density correlation function $\langle \delta N^e(\mathbf{r}') \delta N^e(\mathbf{r}_1') \rangle$ is of the order of a Debye length, λ_D . This is a much shorter distance than the hydrodynamical scale of plasma parameter variations or the spatial variations of a probe intensity in the plasma. Exploring this separation of scales we introduce in Eq. (7) a Fourier transform with respect to a difference between spatial variable, $\mathbf{r}' - \mathbf{r}_1'$, and retain a long scale spatial variations of the time averaged probe energy flux $cE_0^2(\mathbf{r}')/8\pi$ and of the dynamical form factor $S(\mathbf{k}, \omega; \mathbf{r}') = \langle \delta N^e \delta N^e \rangle(\mathbf{k}, \omega)/n_{0e}$. The argument $\mathbf{r}'$ in the dynamical form factor, S , corresponds to long scale variations described by the hydrodynamical variables. We have introduced $S(\mathbf{k}, \omega)$, which is calculated locally at different points in an inhomogeneous plasma. Thus, the expression for the scattered power, P_s , into a solid angle $d\Omega$ in the frequency range $d\omega$ reads

$$P_s = (\hat{\eta}_0 \cdot \hat{\eta}_1)^2 r_0^2 n_{0e} \frac{d\Omega}{2\pi} \int_V d^3r' \frac{cE_0^2(\mathbf{r}')}{8\pi} S(\mathbf{k}, \omega; \mathbf{r}'), \quad (8)$$

where the wave-vector $\mathbf{k}$ and the frequency ω are related by the condition

$$\mathbf{k} = \mathbf{k}_0 - \frac{\omega_0 - \omega}{c} \hat{n}. \quad (9)$$

The expression for the power scattered by electrons (8) at the frequency $\omega_1 = \omega_0 - \omega$ depends on the electron density, scattering volume V , and the dynamical form factor $S(\mathbf{k}, \omega)$ describing a plasma response. Expression (9) shows a relation between the direction of observation $\hat{n}$ and the wave-vector $\mathbf{k}$ which is probed in the plasma. For $k < k_D$ the main contributions to the scattering cross section are provided by the resonances at frequencies ω corresponding to the ion-acoustic and Langmuir fluctuations at the particular k .

III. DYNAMICAL FORM FACTOR

The usefulness of the TS as the diagnostic technique is greatly enhanced when applied with the relevant theoretical expressions for the $S(\mathbf{k}, \omega)$. Equally important is a proper modelling of the TS cross-section in the radiation transport codes.

Consider three different expressions for a dynamical form factor which are applicable in plasmas: without collisions, with high ion-ion collisionality and described by the non-Maxwellian electron distribution function produced during an intense inverse bremsstrahlung heating. All three cases illustrate the potential of the TS as the plasma diagnostic, and are of relevance to the TS measurements from the high-Z, gold plasmas [10].

A. Collisionless approximation

In collisionless plasmas the power spectrum of density fluctuations in the frequency and wave vector space is a well known quantity (cf. e.g. [11]), which reads

$$S(\mathbf{k}, \omega) = \frac{2\pi}{k} \left| 1 - \frac{\chi^e(\mathbf{k}, \omega)}{\epsilon(\mathbf{k}, \omega)} \right|^2 \bar{f}^e(v = \omega/k) + \left| \frac{\chi^e(\mathbf{k}, \omega)}{\epsilon(\mathbf{k}, \omega)} \right|^2 \sum_{\beta(\neq e)} Z_\beta^2 \frac{n_{0\beta}}{n_{0e}} \bar{f}^\beta(v = \omega/k), \quad (10)$$

where $\bar{f}^\beta(v = \omega/k)$ stands for the one dimensional distribution function of a β -type particles evaluated at the velocity ω/k . A dielectric response function

$$\epsilon(\mathbf{k}, \omega) = 1 + \sum_{\beta} \chi^\beta(\mathbf{k}, \omega), \quad (11)$$

is defined by the susceptibility function,

$$\chi^\alpha(\mathbf{k}, \omega) = \frac{\omega_{p\alpha}^2}{k^2} \int d^3v \frac{1}{\omega - \mathbf{k} \cdot \mathbf{v} + i\epsilon} \mathbf{k} \cdot \frac{\partial f^\alpha(\mathbf{v})}{\partial \mathbf{v}}. \quad (12)$$

Since the original derivation of an expression for the dynamical form factor (cf. Ref. [11] and references therein), Eq. (10) has been applied in the interpretation of numerous experiments. Here we would like to illustrate two features of this theory: (i) $S(\mathbf{k}, \omega)$ (10) can be calculated

locally with the f^α given, for example, by the local Maxwellian in inhomogeneous plasmas with a scale separation between the correlation length and hydrodynamical gradients (as discussed in Sec. II); (ii) the most straightforward and reliable measurement of plasma parameters can be obtained from the simultaneous scattering spectra in frequency ranges corresponding to ion acoustic and Langmuir resonances.

Figure 1 shows $S(\mathbf{k}, \omega)$ (10) for the parameters of a gold homogeneous plasma ($Z = 42$, $T_e = 2keV$, $n_e = 2.1 \times 10^{20} cm^{-3}$) plotted as the function of a scattered light wavelength for $(k\lambda_D)^{-1} = 2.3$. The probe wavelength is 5266 Å. From the separation of two ion acoustic peaks, Fig. 1, one can estimate the product ZT_e , which defines in high- Z plasmas the value of an ion-acoustic frequency. The Langmuir wave resonance (Fig. 1) is then used to calculate electron density, n_{0e} , and temperature, T_e . This can be done with a reasonable accuracy because for the small value of $(k\lambda_D)^{-1}$, corresponding to strongly damped Langmuir fluctuations the overall shape of the $S(\mathbf{k}, \omega)$ is sensitive to both temperature and density variations. After finding T_e one can improve an estimate of Z by matching again the separation between two ion-acoustic peaks with a more accurate estimate of T_e . Figure 1 illustrates also the difficulties in practical realization of this straightforward procedure. They are related to very different sensitivities required in measurements of low and high frequency parts of the spectra both in terms of spectral resolutions and sensitivities of the instruments.

The scattering cross-section from Eq. (8) which is calculated using $S(\mathbf{k}, \omega)$ (10) and local Maxwellian distribution functions

$$f^\alpha(\mathbf{v}, \mathbf{r}) = \left(\frac{m_\alpha}{2\pi T_\alpha(\mathbf{r})} \right)^{3/2} \exp \left[-\frac{m_\alpha(\mathbf{v} - \mathbf{u}_\alpha(\mathbf{r}))^2}{2T_\alpha(\mathbf{r})} \right], \quad (\alpha = e, i), \quad (13)$$

is shown in Fig. 2. Calculations of the spatial integral in Eq. (8) is simplified by the geometry of the experiment, where the direction of observation, $\hat{n}$, is approximately in the plane of $E_0(\mathbf{r})$ variations, i.e. in the transverse direction to $\mathbf{k}_0$. Figure 2 shows an ion-acoustic resonance which is Doppler shifted due to the flows, $\mathbf{u}_\alpha$ and exhibits small asymmetry between the peaks due to the drift of electrons with respect to ions, $u_e \neq u_i$. Note how much broader are the resonances as compared to simple theory in Fig. 1. This is almost entirely

due to the inhomogeneity of the flows, which are approximated by the linear functions of position, $u_\alpha = u_{0\alpha}(1 + x/L_u)$, where x is along the direction of the expanding plasma and $L_u = 400 \mu\text{m}$ is also in the very good agreement with hydrodynamical simulations.

B. Frequent ion-ion collisions

One of the characteristics of highly ionized plasmas is a small Landau damping of ion acoustic waves due to the large ratio, ZT_e/T_i . Small damping is responsible for sharpness of the ion acoustic resonances shown in Fig. 1. At the same time a large Z value makes ion-ion collision frequency, which is proportional to Z^2 , comparable to ion-acoustic frequency. Therefore, plasma response at the ion-acoustic resonance, described by $S(\mathbf{k}, \omega)$ should include collisional effects.

The formulation of a dynamical form factor which is valid in the entire regime of particle collisionality has been recently discussed in Ref. [13]. This theory is based on the non-local fluctuating hydrodynamical model. $S(\mathbf{k}, \omega)$ takes the following form in the limit of collisionless electrons and collisional ions,

$$S(k, \omega) = \frac{4k^2(c_s^2 + v_{Ti}^2)\gamma_a}{(\omega^2 - k^2v_s^2)^2 + 4\omega^2\gamma_a^2}, \quad (14)$$

where we have introduced the definitions,

$$v_s = \sqrt{c_s^2 + \Gamma_i v_i^2}, \quad \text{and} \quad \Gamma_i = \frac{5}{3} + \frac{4}{3} \frac{\omega}{v_i} \text{Im} \tilde{\eta} = \frac{9\omega^4 + 29.7\omega^2\nu_i^2 + 11.7\nu_i^4}{3(\omega^4 + 4.05\omega^2\nu_i^2 + 2.33\nu_i^4)}, \quad (15)$$

for the ion acoustic group velocity and ion specific heat ratio. The damping of an ion-acoustic wave, γ_a involves electron Landau damping and the frequency dependent ion viscosity, $\tilde{\eta}$,

$$\gamma_a = \sqrt{\frac{\pi}{8}} \frac{c_s}{v_{Te}} k v_s + \frac{2}{3} \frac{k^2 v_{Ti}^2}{\nu_i} \text{Re} \tilde{\eta}, \quad (16)$$

where,

$$\text{Re} \tilde{\eta} = \frac{3}{2} k^2 v_{Ti}^2 \frac{\nu_i^2 (1.49\nu_i^2 + 0.80\omega^2)}{\omega^4 + 4.05\omega^2\nu_i^2 + 2.33\nu_i^4}, \quad (17)$$

and $\nu_i = 4\sqrt{\pi}Z^4e^4n_i\Lambda_i/3\sqrt{m_i}T_i^{3/2}$ is the ion-ion collision frequency (Λ_i is a Coulomb logarithm). In spite of high values of ν_i in gold plasmas the effect of collisions on the shape of an ion-acoustic resonance is moderate. As expected Eq. (14) produces broadening of the resonance as compared to collisionless expression (10). However, as the values of $S(k, \omega)$ differ by only 10 - 30 %, this improvement has negligible effect on the interpretation of experimental results [10,12]. There, a $0.5\text{\AA}$ resolution of the spectrometer and an inhomogeneity of the plasma are far more important in determining the shape of the resonance.

C. Super-Gaussian distribution function

An important characteristic of a dense high- Z plasma is a strong inverse bremsstrahlung (IB) absorption. Langdon [14] has shown that due to IB the electron distribution function would take the form proportional to $\exp[-(v/u(t))^\mu]$, $2 \leq \mu \leq 5$, for the parameter $Zv_0^2/v_e^2 \gg 1$, where v_0 is the oscillatory velocity of an electron in the electric field of the electromagnetic wave and $u(t)$ increases in time during IB plasma heating. Next, Matte et al. [15] introduced a stationary version of this distribution function

$$f^e(v) = \frac{\mu}{4\pi v_e^3 \Gamma(3/\mu)} \exp \left[- \left(\frac{v}{v_*} \right)^\mu \right], \quad (18)$$

where $v_*^2 = 3\Gamma(3/\mu)v_e^2/\Gamma(5/\mu)$ [$v_e^2 = (1/3) \int d^3v v^2 f^e(v)$] and the index μ is related to the laser intensity by the fitting expression

$$\mu = 2 + 3/(1 + 1.66/\alpha^{0.724}), \quad \alpha = Zv_0^2/v_e^2. \quad (19)$$

In the experimental measurements [10,12], which motivate our study, the parameter μ (19) corresponds approximately to $3 \leq \mu \leq 3.5$. The Thomson probe in these measurements has an intensity $I \approx 2.5 \times 10^{14} \text{ W/cm}^2$ and gives rise to IB heating.

Zheng et al [16] suggested that a Thomson scattering could be applied to verify theoretical predictions (18) for the electron distribution function. They pointed out that only simultaneous measurements of the scattered light in the frequency regimes corresponding to

ion-acoustic and Langmuir wave resonances could provide necessary data to infer electron temperature and density in plasmas described by the non-Maxwellian distribution functions. These distribution functions change the ion-acoustic dispersion relation and therefore could complicate identification of the electron temperature. The Langmuir wave resonance for the parameters of the experiment [12] corresponds to very strongly damped fluctuations (cf. Fig. 1) and therefore even small variations in the slope of the distribution function results in dramatic changes to the scattering cross-section, as it is illustrated in Fig. 3. We have plotted in Fig. 3 part of the dynamical form factor (10) corresponding to the Langmuir wave resonance using modified distribution functions (18) with different exponents. Dramatic changes in the profile of $S(\mathbf{k}, \omega)$ with moderate variations of μ are not observed in the experimental results. As the TS experiments by Glenzer et al. [10] are the first accurate measurements in plasmas which should display features revealing non-Maxwellian distribution functions, their apparent absence must prompt more studies on the limits of applicability of the solutions (18) and (19).

IV. CONCLUSIONS

We have reported on the theories relevant to the TS experiments in highly ionized laser produced plasmas which have been applied in the interpretation of experimental measurements [12]. The full account of these studies, including results of experiments and detailed comparisons with Lasnex simulations will be published elsewhere [12].

The potential of the TS can be fully explored only by involving the relevant theoretical results for the scattering cross section and dynamical form factor. We have derived and applied to the experimental conditions a Thomson scattering cross section (8) in inhomogeneous plasmas. Using local Maxwellian distribution functions we have found values of important plasma parameters by accurately matching spectra of Thomson scattered radiation, Fig. 2. A simultaneous analysis of the scattered light spectra in the frequency range corresponding to ion acoustic and Langmuir fluctuations in a hot and dense plasma

has demonstrated means of accurate measurements of the ionization level Z . This could established TS as an important diagnostic in opacity studies of dense plasmas.

Finally we have compared theoretical predictions for the TS cross section for electrons described by the non-Maxwellian distribution functions with experimental results. Super-Gaussian distributions are expected in plasmas with strong IB heating where $Zv_0^2/v_e^2 > 1$. Changes in the electron distribution function involve flattening of f^e at small particle velocities and reduction of the number of fast electrons. Our comparisons show convincingly no changes in the tail of the distribution function over wide range of flow velocities, densities and electron temperature accessible in the experiment. Further studies are necessary in order to improve theoretical basis for non-Maxwellian distributions. It has been shown, for example, by theory and Fokker-Planck simulations [17] that localised heating of the plasma due to IB absorption leads to distribution functions which are linearly unstable to the return current instability. Signatures of this instability has been also observed in the experiment [12].

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FIGURES

FIG. 1. $S(\mathbf{k}, \omega)$ (10) for the parameters of a gold homogeneous plasma ($Z = 42$, $T_e = 2keV$, $n_e = 2.1 \times 10^{20} cm^{-3}$) plotted as the function of a scattered light wavelength for $(k\lambda_D)^{-1} = 2.3$. The probe wavelength is 526.6 nm. The insert shows two ion acoustic resonances centered around the probe wavelength. They are shown in the main plot as two unresolved lines.

FIG. 2. Comparison between experimental measurements (dots) and theoretical predictions based on the scattering cross-section for the inhomogeneous plasma (8). The plasma parameters are: $Z = 42$, $T_e = 2.1keV$, $n_e = 2.1 \times 10^{20} cm^{-3}$, $u_e = 1.47c_s$, $u_e - u_i = -0.15c_s$.

FIG. 3. Dynamical form factor for the non-Maxwellian distribution functions (18) for the down-shifted Langmuir resonance ($n_e = 7 \times 10^{19} cm^{-3}$, $T_e = 1.8keV$). Four plots correspond to $\mu = 2, 2.2, 2.4, 3$, from the curve extending furthest to the right to the curve ending at the shortest wavelength.

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